New $D = 4$ gauged supergravities from $\mathcal{N} = 4$ orientifolds with fluxes

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Abstract

We consider classes of $T_6$–orientifolds, where the orientifold projection contains an inversion $I_{9-p}$ on $9 - p$ coordinates, transverse to a D$p$–brane. In absence of fluxes, the massless sector of these models corresponds to diverse forms of $\mathcal{N} = 4$ supergravity, with six bulk vector multiplets coupled to $\mathcal{N} = 4$ Yang–Mills theory on the branes. They all differ in the choice of the duality symmetry corresponding to different embeddings of $SU(1,1) \times SO(6,6+n)$ in $Sp(24+2n,\mathbb{R})$, the latter being the full group of duality rotations. Hence, these Lagrangians are not related by local field redefinitions. When fluxes are turned on one can construct new gaugings of $\mathcal{N} = 4$ supergravity, where the twelve bulk vectors gauge some nilpotent algebra which, in turn, depends on the choice of fluxes.
1. Introduction

New string or M–theory models are obtained turning on \( n \)–form fluxes, which allow, in general, the lifting of vacua, supersymmetry breaking and moduli stabilisation [1]–[24]. Examples of such new solutions are IIB and IIA orientifolds [25, 26, 27, 28, 29], where the orientifold projection (in absence of fluxes) preserves \( \mathcal{N} = 4 \) or \( \mathcal{N} = 2 \) supersymmetries.

Recently, the \( T_6/\mathbb{Z}_2 \) orientifold with \( \mathcal{N} = 4 \) supersymmetry [8, 9] and \( K_3 \times T_2/\mathbb{Z}_2 \) orientifold [24] with \( \mathcal{N} = 2 \) supersymmetry have been the subject of an extensive study. In these cases, turning on NS–NS and R–R three–form fluxes allows to obtain new string vacua with vanishing vacuum energy, reduced supersymmetry and moduli stabilisation [7, 8, 9, 21, 24]. These features can all be understood in terms of a \( n \) effective gauged su-
pergravity, where certain axion symmetries are gauged [31, 30, 3 2]. These are generalised no–scale models [33, 34].

In the present investigation, we consider more general four–dimensional orientifolds with fluxes (both in type IIB and IIA) where the orientifold projection involves an inversion \( I_{9−p} \) on \( 9−p \) coordinates, transverse to the Dp–brane world–volume, thus generalising the \( T_6/\mathbb{Z}_2 \) orientifold (with \( p = 3 \)) constructed by Frey–Polchinski [8] and Kachru–Shulz–
Trivedi [9] (see also [10] for a derivation of the complete low–energy supergravity from T–dialysed Type I theory in ten dimensions). Interestingly, their low–energy descriptions are all given in terms of an effective gauged supergravity, where certain axion symmetries are gauged [31, 30, 32]. These are generalised no–scale models [33, 34].

However, despite the uniqueness of \( \mathcal{N} = 4 \) supersymmetry, the low–energy actions crucially differ in the choice of the manifest “duality symmetries” of the Lagrangian, since different sets of fields survive the orientifold projection, and therefore different symme-
tries are manifestly preserved. Leaving the brane degrees of freedom aside, these duality symmetries are specified by their action on the (twelve) bulk vectors. Actually, \( \mathcal{N} = 4 \) supergravity demands that such symmetries be contained in \( SU(1, 1) \times SO(6, 6) \) [35, 36] and act on the vector field strengths and their duals as symplectic \( Sp(24, \mathbb{R}) \) transformations [37]. On the other hand, the symmetries of the Lagrangian correspond to block–lower–
triangular symplectic matrices, whose block–diagonal components have a definite action on the vector potentials [38, 39, 40]. For instance, in the orientifold models containing an \( I_{9−p} \) inversion, the block–diagonal symmetries always include \( GL(9−p, \mathbb{R}) \times GL(p−3, \mathbb{R}) \), as maximal symmetry of the \( GL(6, \mathbb{R}) \) associated to the moduli space of the six–torus met-
rics. The lower–triangular block contains the axion symmetries of the R–R scalars and of the NS–NS ones originating from the \( B \)–field, whenever present\(^1\).

\(^1\)For example, the latter is not present in the \( p = 3 \) case, i.e. the \( T_6/\mathbb{Z}_2 \) orientifold.
In the sequel, we describe all nilpotent algebras $N_p$, corresponding to axion symmetries of the R–R and NS–NS scalars for all orientifold models. All $N_p$’s are nilpotent subalgebras of $so(6,6)$, are generically non–abelian and contain central charges. There are four of them in type IIB ($p = 3, 5, 7, 9$) with dimensions 15, 23, 23, 15 respectively, while there are only three of them in type IIA ($p = 4, 6, 8$) of dimensions 20, 24, 20, respectively. A common feature of these algebras is that they always contain fifteen R–R axionic symmetries, while the extra symmetries correspond to NS–NS $B$–field axions in the bi–fundamental of $GL(9 − p, \mathbb{R}) \times GL(p − 3, \mathbb{R})$.

A further R–R axion symmetry originates from the $SU(1,1)$, which acts as electric–magnetic duality on the gauge fields living on the brane world–volume. The corresponding axion field can be identified with the $C_{p−3}$ R–R field, as dictated by the coupling

$$\int_{\Sigma_{p+1}} C_{p−3} \wedge F \wedge F,$$

where $F$ is the two–form field strength of gauge fields living on the branes.

Turning on fluxes in the orientifold models (three- and five–form fluxes in type IIB, two- and four–form fluxes in IIA) corresponds to a “gauging” in the corresponding supergravity Lagrangian, whose couplings are dictated by the particular choice of fluxes. Non–abelian gaugings may also occur corresponding to subalgebras of $N_p$, or quotient algebras $N_p/Z$, where $Z$ are some of the central generators of $N_p$.

As an illustrative example, let us consider the $p = 7$ type IIB orientifold defined in section 2, where the non–vanishing NS–NS and R–R fluxes are $H_{aij}$, $F_{aij}$, $G^i = \epsilon^{ijk} G_{abjkl}$ ($a, b = 5, 6$ and $i, j = 1, \ldots , 4$), and let us look at terms involving the axions coming from the $B$ and four–form fields, $B_{ia}$ and $C_{ijab} = C_{ij} \epsilon_{ab}$. Inspection of the three–form kinetic term reveals a non–abelian gauge coupling proportional to

$$\sqrt{−g} H_{aij} H_{\mu\nu b} g^{ab} g^{ij} \mu, \nu,$$

as well as axion gauge couplings proportional to

$$\sqrt{−g} H_{aij} H_{\mu\nu b} g^{ab} g^{ij} \mu, \nu,$$

together with similar expressions for the $F$–three form. Such terms come also from the reduction of type IIB four–form field. In addition, when a five–form flux $G^\ell$ is turned on an axion gauge coupling emerges of the type

$$\partial_{\mu} C_{ij} + \epsilon_{ijk\ell} G^k G^{\mu\ell},$$

where $G^{\mu\ell} = g^{\ell} g_{\mu}$ are the Kaluza–Klein vectors. We report here only a preliminary analysis of the deformation of the $\mathcal{N} = 4$ supergravity due to these new gaugings.
In the present paper we do not address either the question of unbroken supersymmetries or the question of moduli stabilisation, which would require the knowledge of the scalar potential and a study of the fermionic sector. However, we can anticipate that certain moduli are indeed stabilised in all these models, since a Higgs effect is taking place as suggested by the presence of charged axion couplings.

The paper is organised as follows: in section 2, we review the four-dimensional $T_6/\mathbb{Z}_2$ orientifold models, their spectra and their allowed fluxes. In section 3, the $\mathcal{N} = 4$ supergravity interpretation is given for the ungauged case (absence of fluxes) and the duality symmetries exposed. The $N_p$ algebras are exhibited as well as their action on the vector fields. In section 4, we give a preliminary description of gauged supergravity, for the particular case of type IIB orientifolds with some three–form fluxes turned on. In section 5 some conclusions are drawn. Finally, in appendix some useful formulae needed to compute the quadratic part of the vector field strengths in the Lagrangian, are given.

2. $\mathcal{N} = 4$ orientifolds: spectra and fluxes

In this section we review the construction of orientifold models preserving $\mathcal{N} = 4$ supersymmetries in $D = 4$ [29]. This is the simplest setting for orientifold constructions, and consists of modding out type II superstrings by the world–sheet parity $\Omega$ [25]. Following [28, 29], the orientifold projection can be given a suggestive geometrical interpretation in terms of non–dynamical defects, the orientifold $\mathcal{O}$–planes, that reflect the left–handed and right–handed modes of the closed string. Actually, one can combine world–sheet parity with other (geometrical) operations. In general, this can affect the nature of the orientifold planes, that, in the simplest instance of a bare $\Omega$ have negative tension and R–R charge, and are $(9 + 1)$–dimensional ($\mathcal{O}9$ planes) since they have to respect the full Lorentz symmetry preserved by $\Omega$. In the present paper, we are interested in the class of models generated by the $\Omega I_{9–p}$ generator, where $I_{9–p}$ denotes the inversion on $9 – p$ coordinates. Of course, $\Omega I_{9–p}$ must be a symmetry of the parent theory, and this is the case of type IIB for $p$ odd, and of type IIA for $p$ even. Actually, $\Omega I_{9–p}$ reflects the action of T-duality in orientifold models. Indeed, T-duality itself can be thought of as a chiral parity transformation

$$X_L \to X_L, \quad X_R \to -X_R,$$

and conjugates $\Omega$ so to get

$$\mathcal{O}_{9–p} \Omega \mathcal{O}^{-1}_{9–p} = \Omega I_{9–p}.$$

As a result, the full ten–dimensional Lorentz symmetry is now broken to the subgroup $SO(1, p) \times SO(9–p)$, and the closed–string sector involves $\mathcal{O}_{9–p}$ planes sitting at the fixed points of the orbifold $T_{9–p}/I_{9–p}$. The associated open-string sector will then correspond
to open strings with Dirichlet boundary conditions along $T_{9-p}$, i.e. open strings ending on $D(9-p)$ branes. As usual, tadpole conditions will fix the rank of the Chan–Paton gauge group, i.e. the total number of D-branes. In the present paper, however, we shall not be concerned with open–string degrees of freedom and we shall concentrate our analysis solely on the closed-string degrees of freedom.

Before we turn to the description of specific models, a general comment is in order. An important requirement in the construction is that the orientifold group be $\mathbb{Z}_2$, i.e. its generator $\Omega_{9-p}$ must square to the identity. Although $\Omega$ has always $\pm 1$ eigenvalues, and thus $\Omega^2 = 1$, this is not the case for $I_{9-p}$. For example, for $p = 7$ $I_2$ would correspond to a $\pi$ rotation on a two–plane and, although its action on the bosonic degrees of freedom is real and assigns to them a plus or minus sign according to the number of indices along the two–plane, its eigenvalue on spinors is $e^{i\pi \Sigma}$, where $\Sigma = \pm \frac{1}{2}$ are the two helicities. Thus, it does not square to the identity, but rather to $(-1)^F$, with $F$ the (total) space–time fermion number. Therefore, in this case the orientifold projection needs be modified by the inclusion of $(-1)^F L$, with $F_L$ the left-handed space–time fermion number \cite{47}. We are thus dealing with the four-dimensional orientifolds

$$\big( T_{p-3} \times T_{9-p} \big) / \Omega I_{9-p} \big[ (-1)^F_L \big]\left[ \frac{9-p}{2} \right] , \quad (7)$$

where $\left[ \frac{9-p}{2} \right]$ denotes the integer part of $(9-p)/2$. Here we have decomposed the six-torus as

$$T_6 = T_{p-3} \times T_{9-p} , \quad (8)$$

since $I_{9-p}$ only acts on the coordinates of $T_{9-p}$, while leaves invariant those along $T_{p-3}$. As we shall see, this is a natural decomposition since, in the orientifold, we are left with the perturbative symmetry $GL(p-3) \times GL(9-p)$ of the compactification torus. To fix the notation, in this paper we shall label coordinates on the $T_6$ with a pair of indices $(i,a)$, where $i = 1,\ldots,p-3$ counts the coordinates not affected by the space parity (those coordinates that would be longitudinal to the branes), while $a = 1,\ldots,9-p$ runs over the coordinates of $T_{9-p}$ (orthogonal to the branes). As usual, Greek indices $\mu, \nu, \ldots$ will label coordinates on the four–dimensional Minkowski space–time.

At this point, it is better to consider the cases $p$ odd or $p$ even separately. In the first case, $\Omega I_{9-p} \big[ (-1)^F_L \big]\left[ \frac{9-p}{2} \right]$ is a symmetry in type IIB, while in the latter case it is properly defined within type IIA.

\textbf{2.1. IIB orientifolds}

In type IIB superstring we have to consider four cases, corresponding to the allowed choices $p = 9, 7, 5, 3$. The massless ten-dimensional fields have a well defined parity with
respect to $\Omega$:

\[
\text{even : } \quad G_{MN}, \quad \phi, \quad C_{MN}, \quad (9) \\
\text{odd : } \quad B_{MN}, \quad C, \quad C_{MNPQ}^{(+)}, \quad (10)
\]

where $G_{MN}$ is the metric tensor, $\phi$ the dilaton, $B_{MN}$ the Kalb–Ramond two-form, and $C_{p+1}$ are the R–R $(p+1)$-forms$^2$. Henceforth, it is straightforward to select the four-dimensional excitations that survive the orientifold projection. In fact, after splitting the ten-dimensional index $M$ in the triple $(\mu, i, a)$ labelling $\mathcal{M}_{1,3} \times T_{p-3} \times T_{9-p}$, it is evident that the fields with an odd (even) number of $a$–type indices are odd (even) under the action of $I_{9-p}$. On the other hand, when present, $(-1)^F_L$ assigns a plus sign to the NS-NS states (which originate from the decomposition of the product of two bosonic representations of $SO(8)$) and a minus sign to the R–R states (which originate from the decomposition of the product of two spinorial representations of $SO(8)$). At the end, aside from the four–dimensional metric tensor, one is left with the massless (bosonic) degrees of freedom listed in table 1.

Table 1: Massless degrees of freedom for the IIB orientifolds

| $p$ | scalars | vectors |
|-----|---------|---------|
| 9   | $g_{ij}, \phi, C_{\mu\nu}, C_{ij}$ | $G^{'\mu}, C_{\mu}$ |
| 7   | $g_{ij}, g_{ab}, \phi, B_{ia}, C, C_{ia}, C_{ijkl}, C_{ijab}$ | $G^i_{\mu}, B_{a\mu}, C_{a\mu}, C_{ijkl\mu}$ |
| 5   | $g_{ij}, g_{ab}, \phi, B_{ia}, C_{\mu\nu}, C_{ij}, C_{ab}, C_{iabc}$ | $G^i_{\mu}, B_{a\mu}, C_{i\mu}, C_{abc\mu}$ |
| 3   | $g_{ab}, \phi, C, C_{abcd}$ | $B_{a\mu}, C_{a\mu}$ |

However, in orientifold models it happens often that fields which are odd under the projection can be consistently assigned with a (quantised) background value for the fields themselves, or for their field strengths. For example, in the $p = 7$ case the NS–NS fields $B_{ij}$ and the R–R fields $C_{ij}$ are both odd with respect to the orientifold projection and, thus, their quantum excitations are projected out. However, acting on them with a $\partial_a$ derivative changes their parity, and thus (quantised) fluxes along the internal directions, $H_{aij}$ and $F_{aij}$, can be incorporated in the model. Repeating a similar analysis for the other cases yields the allowed fluxes listed in table 2.

2.2. IIA orientifolds

$^2$Actually, the four-form $C_{4}^{(+)}$ is constrained to have a self–dual field strength, a peculiarity of type IIB
Table 2: Allowed fluxes for the IIB orientifolds. $F$, $H$ and $G$ fluxes are associated to the $B$, $C_2$ and $C_4$ fields

| $p$ | fluxes       |
|-----|--------------|
| 9   | none         |
| 7   | $H_{ija}$, $F_{ija}$, $G_{ijkab}$ |
| 5   | $H_{abc}$, $F_{iab}$, $H_{ija}$, $G_{ijabc}$ |
| 3   | $H_{abc}$, $F_{abc}$ |

Type IIA superstring selects $p$ even, and thus leaves us with the three cases $p = 8, 6, 4$. Although a bare $\Omega$ is not a symmetry in type IIA, we can nevertheless assign a well defined parity to the massless ten-dimensional degrees of freedom:

$$
even: \quad G_{MN}, \quad \phi, \quad C_M, \quad (11)$$
$$
odd: \quad B_{MN}, \quad C_{MNP}. \quad (12)$$

As before, $G_{MN}$ is the metric tensor, $\phi$ the dilaton, $B_{MN}$ the Kalb–Ramond two-form, while in this case the R–R potentials $C_{p+1}$ carry an odd number of indices. The additional action of $I_{9-p}$ and, eventually, of $(-1)^F_L$ thus yields the massless degrees of freedom listed in table 3.

Also in this case one can allow for (quantised) fluxes along the compactification torus, as summarised in table 4.

Table 3: Massless degrees of freedom for the IIA orientifolds

| $p$ | scalars                     | vectors                  |
|-----|-----------------------------|--------------------------|
| 8   | $g_{ij}, \ g_{99}, \ \phi, \ B_{i9}, \ C_i, \ C_{9\mu\nu}, \ C_{ij9}$ | $G^i_\mu, \ C_\mu, \ C_{i9\mu}, \ B_{9\mu}$ |
| 6   | $g_{ij}, \ g_{ab}, \ \phi, \ B_{ia}, \ C_a, \ C_{ij\mu\nu}, \ C_{ijk}, \ C_{iab}$ | $G^i_\mu, \ B_{a\mu}, \ C_{ij\mu}, \ C_{ab\mu}$ |
| 4   | $g_{44}, \ g_{ab}, \ \phi, \ B_{4a}, \ C_4, \ C_{a\mu\nu}, \ C_{abc}$ | $C^4_{\mu}, \ B_{a\mu}, \ C_\mu, \ C_{4a\mu}$ |
Table 4: Allowed fluxes for the IIA orientifolds. $F$, $H$ and $G$ fluxes are associated to the $B$, $C_1$ and $C_3$ fields

| $p$ | fluxes               |
|-----|---------------------|
| 8   | $H_{ij9}$, $G_{ijk9}$ |
| 6   | $H_{aij}$, $H_{abc}$, $F_{ia}$, $G_{ijab}$ |
| 4   | $H_{abc}$, $F_{ab}$, $G_{4abc}$ |

3. $\mathcal{N} = 4$ supergravity interpretation of $T_6$ orientifolds: manifest duality transformations and Peccei–Quinn symmetries.

The four–dimensional low–energy supergravities of $\mathcal{N} = 4$ orientifolds (in the absence of fluxes) can be consistently constructed as truncations of the unique four–dimensional $\mathcal{N} = 8$ supergravity which describes the low–energy limit of dimensionally reduced type II superstrings. Its duality symmetry group $E_{7(7)}$ acts non linearly on the 70 scalar fields, and linearly, as a $Sp(56,\mathbb{R})$ symplectic transformation, on the 28 electric field strengths and their magnetic dual. In this framework an intrinsic group–theoretical characterisation of the ten–dimensional origin of the four–dimensional fields is indeed achieved. In the so–called solvable Lie algebra representation of the scalar sector \[41, 42\], the scalar manifold

$$\mathcal{M}_{\text{scal}} = \exp(Solv(e_{7(7)}))$$

is expressed as the group manifold generated by the solvable Lie algebra $Solv(e_{7(7)})$ defined through the Iwasawa decomposition of the $e_{7(7)}$ algebra:

$$e_{7(7)} = su(8) + Solv(e_{7(7)}).$$

In this framework, there is a natural one–to–one correspondence between the scalar fields and the generators of $Solv(e_{7(7)})$. The latter consists of the 7 generators $H_p$ of the $e_{7(7)}$ Cartan subalgebra, parametrised by the $T_6$ radii $R_a = e^{\sigma a}$ together with the dilaton $\phi$, and of the shift generators corresponding to the 63 positive roots $\alpha$ of $e_{7(7)}$, which are in one–to–one correspondence with the axionic scalars that parametrise them. This correspondence between Cartan generators and positive roots on one side and scalar fields on the other, can be pinpointed by decomposing $Solv(e_{7(7)})$ with respect to some relevant groups. For instance, the duality group of maximal supergravity in $D$ dimensions is $E_{11-D(11-D)}$ and therefore, in the solvable Lie algebra formalism, the scalar fields in the $D$–dimensional theory are parameters of $Solv(e_{11-D(11-D)})$. Since $e_{11-D(11-D)} \subset e_{7(7)}$, ...
decomposing $Solv(e_{7(7)})$ with respect to $Solv(e_{11-D(11-D)})$ it is possible to characterise the higher–dimensional origin of the four–dimensional scalars. Moreover, in four dimensions the group $SL(2,\mathbb{R}) \times SO(6,6)_T \subset E_{7(7)}$, $SO(6,6)_T$ being the isometry group of the $T_6$ moduli–space, acts transitively on the scalars originating from ten–dimensional NS–NS fields of type II theories. These scalars therefore parametrise $Solv(sl(2,\mathbb{R}) + so(6,6)_T)$. Henceforth, decomposing $Solv(e_{7(7)})$ with respect to $Solv(sl(2,\mathbb{R}) + so(6,6)_T)$ one can achieve an intrinsic characterisation of the NS–NS or R–R ten–dimensional origin of the four–dimensional scalar fields, the R–R scalars (and the corresponding solvable generators) transforming in the spinorial representation of $SO(6,6)_T$. Finally, depending on whether we interpret the four–dimensional maximal supergravity as tied to type II supergravities on $T_6$ or $D = 11$ supergravity on $T_7$, the metric moduli are acted on transitively by $GL(6,\mathbb{R})_g$ or $GL(7,\mathbb{R})_g$ subgroups of $E_{7(7)}$, respectively. Therefore, in the two cases the metric moduli parametrise $Solv(gl(6,\mathbb{R})_g)$ or $Solv(gl(7,\mathbb{R})_g)$ and thus, decomposing $Solv(e_{7(7)})$ with respect to these two solvable subalgebras, depending on the higher–dimensional interpretation of the four–dimensional theory, we may split the axions into metric moduli of the internal torus and into scalars deriving from dimensional reductions of ten- or eleven–dimensional tensor fields. The latter will parametrise nilpotent generators transforming in the corresponding tensor representations with respect to the adjoint action of $GL(6,\mathbb{R})_g$ or $GL(7,\mathbb{R})_g$. As a result of the above decompositions, we are able to characterise unambiguously each parameter of $Solv(e_{7(7)})$ as a dimensionally reduced field. Let us consider the dimensional reduction of type II supergravities. As far as the axionic scalars are concerned the correspondence with roots can be summarised in terms of an orthonormal basis $\{\epsilon_p\}$ of $\mathbb{R}^7$:

$$C_{n_1n_2\ldots n_k} \leftrightarrow a + \epsilon_{n_1} + \ldots + \epsilon_{n_k}, \quad (15)$$
$$C_{n_1n_2\ldots n_k\mu\nu} \leftrightarrow a + \epsilon_{m_1} + \ldots + \epsilon_{m_{6-k}}, \quad (\epsilon^{n_1\ldots n_km_1\ldots m_{6-k}} \neq 0), \quad (16)$$
$$B_{nm} \leftrightarrow \epsilon_n + \epsilon_m, \quad (17)$$
$$B_{\mu\nu} \leftrightarrow \sqrt{2}\epsilon_7, \quad (18)$$
$$G_{nm} \leftrightarrow \epsilon_n - \epsilon_m, \quad (n \neq m), \quad (19)$$

where

$$a = -\frac{1}{2} \sum_{n=1}^{6} \epsilon_n + \frac{1}{\sqrt{2}} \epsilon_7. \quad (20)$$

In our notation, the $so(6,6)_T$ roots have the form $\{\pm \epsilon_n \pm \epsilon_m\}$, where $1 \leq n < m \leq 6$. Notice indeed that the nilpotent generators corresponding to non–metric axions transform

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3Now and henceforth we shall always label by $n, m = 1, \ldots, 6$ the $T_6$ directions, by $i, j = 1, \ldots, p-3$ the directions of $T_{p-3}$ which are longitudinal to the D$p$–brane and by $a, b = p-2, \ldots, 9-p$ the directions of the transverse $T_{9-p}$. The four–dimensional space–time directions are generically denoted by Greek letters.
in tensor representations of $GL(6, \mathbb{R})_g$, and this, in turn, defines the $GL(6, \mathbb{R})_g$ representation of the corresponding scalar. For instance, the $C_{n_1...n_k}$ parametrises the generator $T^{n_1...n_k} = E_{a+n_1+...+n_k}$ whose transformation property under $GL(6, \mathbb{R})_g$ is

$$g \in GL(6, \mathbb{R})_g : \quad g \cdot T^{n_1...n_k} \cdot g^{-1} = g^{n_1}_{m_1} \cdots g^{n_k}_{m_k} T^{m_1...m_k}. \quad (21)$$

The roots corresponding to R–R fields are spinorial with respect to $SO(6,6)_T$ and, depending on whether the number of their indices is even or odd, they belong to the root system of two $e_{7(7)}$ algebras which are mapped into each other by the $SO(6,6)_T$ outer automorphism (T–duality) [43, 44]. These two systems naturally correspond to the reduction of IIB and IIA superstrings, that are indeed related by T-dualities. Hence, the $T_6$ metric moduli in the type IIA or B descriptions, are acted upon transitively by two inequivalent $GL(6, \mathbb{R})_g$ subgroups of $E_{7(7)}$: in the former case $GL(6, \mathbb{R})_g$ is contained in $SL(8, \mathbb{R}) \subset E_{7(7)}$, while in the latter case $GL(6, \mathbb{R})_g$ is contained in the maximal subgroup $SL(3, \mathbb{R}) \times SL(6, \mathbb{R})_g$ of $E_{7(7)}$. As far as the R–R scalars are concerned, the two representations differ in the $SO(6,6)_T$ chirality of the 32 spinorial positive roots

$$\text{IIA : } 32^- = \left\{ \frac{1}{2}(\pm \epsilon_1 \ldots \pm \epsilon_6) + \frac{1}{\sqrt{2}} \epsilon_7 \right\},$$

$$\text{IIB : } 32^+ = \left\{ \frac{1}{2}(\pm \epsilon_1 \ldots \pm \epsilon_6) + \frac{1}{\sqrt{2}} \epsilon_7 \right\}. \quad (22)$$

Similarly, vector potentials, and their corresponding duals, are in one–to–one correspondence with weights $W$ of the 56 of $E_{7(7)}$ in the two representations discussed above:

$$C_{n_1...n_k\mu} \leftrightarrow w + \epsilon_{n_1} + \ldots + \epsilon_{n_k}, \quad B_{m\nu} \leftrightarrow \epsilon_n - \frac{1}{\sqrt{2}} \epsilon_7,$$

$$G^m_\mu \leftrightarrow -\epsilon_n - \frac{1}{\sqrt{2}} \epsilon_7,$$

where

$$w = -\frac{1}{2} \sum_{n=1}^{6} \epsilon_n. \quad (23)$$

The dual potentials correspond to the opposite weights $-W$.

The above axion–root ($\Phi \leftrightarrow \alpha$) and vector–weight ($A_\mu \leftrightarrow W$) correspondences can be retrieved also from inspection of the scalar and vector kinetic terms in the dimensionally reduced type IIA or type IIB Lagrangians [43, 45, 46] on a straight torus, which have the form:

- **dilatonic scalars:** $-\partial_\mu \vec{h} \cdot \partial^\mu \vec{h}$,
- **axionic scalars:** $-\frac{1}{2} e^{-2\alpha} h \left( \partial_\mu \Phi \cdot \partial^\mu \Phi \right)$,
- **vector fields:** $-\frac{1}{4} e^{-2W} hF_{\mu\nu} F^{\mu\nu}$,
where
\[ \vec{h} = \sum_{n=1}^{6} \sigma_n (\epsilon_n + \frac{1}{\sqrt{2}} \epsilon_7) - \frac{1}{2} \phi a , \] (24)
and, as usual, \( F_{\mu \nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \).

A generic axion \( \Phi \) and its dilatonic partner \( e^{\alpha \cdot h} \) can be thought of as the real and imaginary parts of a complex field \( z \) spanning an \( SL(2, \mathbb{R})/SO(2) \) submanifold, where the \( SL(2, \mathbb{R}) \) group is defined by the root \( \alpha \). In the models describing type II strings on \( T_{p-3} \times T_{9-p} \) orientifolds, the real part of the complex scalar \( z \) spanning the \( SL(2, \mathbb{R})/SO(2) \) factor in the scalar manifold is \( C_{i_1 \ldots i_{p-3}} \), where \( i_1, \ldots i_k \) label the directions of \( T_{p-3} \), as dictated by the coupling in eq. (1). From eqs. (15) and (24) one can then verify that
\[ \text{Im}(z) = e^{\alpha \cdot h} = \text{Vol}_{p-3} e^{\frac{2\pi}{V_7}} \phi , \]
where \( \text{Vol}_{p-3} \) denotes the volume of \( T_{p-3} \). The scalar \( \text{Im}(z) \) defines the effective four–dimensional coupling constant of the super Yang–Mills theory on \( Dp \)–branes through the relation:
\[ \frac{1}{g_{YM}^2} = V_{p-3} e^{\frac{2\pi}{V_7}} \phi . \] (25)

The embedding of the \( \mathcal{N} = 4 \) orientifold models \( T_{p-3} \times T_{9-p} \) (in absence of fluxes) inside the \( \mathcal{N} = 8 \) theory (in its type IIA or IIB versions) is defined by specifying the embedding of the \( \mathcal{N} = 4 \) duality group \( SL(2, \mathbb{R}) \times SO(6, 6) \) inside the \( \mathcal{N} = 8 E_{7(7)} \) one. As far as the scalar sector is concerned, this embedding is fixed by the following group requirement:
\[ SO(6, 6) \cap GL(6, \mathbb{R})_g = O(1, 1) \times SL(p - 3, \mathbb{R}) \times SL(9 - p, \mathbb{R}) . \] (26)
Condition (26) fixes the ten–dimensional interpretation of the fields in the ungauged \( \mathcal{N} = 4 \) models (except for the cases \( p = 3 \) and \( p = 9 \)) which, for a given \( p \), is indeed consistent with the bosonic spectrum resulting from the orientifold reductions listed in the previous section. In the \( p = 3 \) and \( p = 9 \) cases, the two embeddings are characterised by a different interpretation of the scalar fields, consistent with the \( T_6/\mathbb{Z}_2 \) orientifold reduction in the presence of \( D3 \) or \( D9 \) branes. We shall denote these two models by \( T_6 \times T_6 \) and \( T_6 \times T_0 \), respectively. In these cases, equation (26) in the solvable Lie algebra language amounts to requiring that metric moduli are related either to the \( T_{p-3} \) metric \( g_{ij} \) or to the \( T_{9-p} \) metric \( g_{ab} \). The scalar field parameterising the Cartan generator of the external \( SL(2, \mathbb{R}) \) factor is given in eq. (25), while the metric modulus corresponding to the \( O(1, 1) \) in eq. (26) is (modulo an overall power)
\[ O(1, 1) \leftrightarrow (V_{p-3})^{9-p} (V_{9-p})^{11-p} . \] (27)

The axions not related to the \( T_6 \) metric moduli consist of \( C_{i_1 \ldots i_{p-3}} \) in the external \( SL(2, \mathbb{R})/SO(2) \) factor, \( (p-3) \) \((9-p)\) moduli \( B_{ia} \) in the bifundamental of \( SL(p-3, \mathbb{R}) \times SL(9-p, \mathbb{R}) \) and 15
R–R moduli which we shall generically denote by $C_I$ and which span the maximal abelian ideal $\{ T^I \}$ of $\text{Solv}(so(6,6))$. The scalars $B_{ia}$ and $C_I$ parametrise a $15 + (p - 3) (9 - p)$ dimensional subalgebra $N_p$ of $\text{Solv}(so(6,6))$ consisting of nilpotent generators only. In figure 4 the $so(6,6)$ Dynkin diagrams for the various models and the corresponding intersections with $gl(6,\mathbb{R})_g$, represented by $sl(p - 3,\mathbb{R}) + sl(9 - p,\mathbb{R})$ subdiagrams, are illustrated. As far as the scalar fields are concerned, the $T_{p-3} \times T_{9-p}$ models within the same type IIA or IIB framework are mapped into each other by $so(6,6)_T$ Weyl transformations, which can be interpreted as T–dualities on an even number of directions of $T_6$.

We can now turn to the detailed analysis of each (IIB or IIA) orientifold model.

![IIB Superstring Diagram](image)

![IIA Superstring Diagram](image)

Figure 1: $SO(6,6)$ Dynkin diagrams for the $T_{p-3} \times T_{9-p}$ models. The shaded subdiagrams define the groups $SL(p - 3,\mathbb{R}) \times SL(9 - p,\mathbb{R})$ acting transitively on the metric moduli. The empty circles define simple roots corresponding to the metric moduli $g_{ij}, g_{ab}$, the grey circle denotes a simple root corresponding to a Kalb–Ramond field $B_{ia}$ and the black circle corresponds to a R–R axion.

3.1. $T_4 \times T_2$ IIB orientifold with D7 branes. The ungauged version.

Solvable algebra of global symmetries. The following model (with $p = 7$) describes the bulk sector of IIB superstring compactified on a $(T_4 \times T_2)/\mathbb{Z}_2$ orientifold with D7 branes wrapped on the $T_4$.

To this end, we describe the embedding of the scalar sector of the corresponding $\mathcal{N} = 4$ model within the $\mathcal{N} = 8$ by expressing the $so(6,6)$ Dynkin diagram $\{ \beta_n \}$ in terms of the simple roots of $e_{7(7)}$ \(^4\)

\[
\beta_1 = \epsilon_1 - \epsilon_2,
\]

\(^4\)In our conventions $\beta_1$ is the end root of the long leg and $\beta_5, \beta_6$ the symmetric roots.
\[ \beta_2 = \epsilon_2 - \epsilon_3, \]
\[ \beta_3 = \epsilon_3 - \epsilon_4, \]
\[ \beta_4 = \epsilon_4 + \epsilon_5, \]
\[ \beta_5 = -\epsilon_5 + \epsilon_6, \]
\[ \beta_6 = -\frac{1}{2}(\sum_{n=1}^{6} \epsilon_n) + \frac{1}{\sqrt{2}}\epsilon_7 = a. \]

According to eq. (15), the root \( \beta_6 \) corresponds to the ten-dimensional R–R scalar \( C_0 \), and thus identifies the type IIB duality group \( SL(2, \mathbb{R})_{IIB} \). The Dynkin diagram of the external \( SL(2, \mathbb{R}) \) factor in the isometry group consists, instead, of the single root

\[ \beta = a + \epsilon_1 + \epsilon_2 + \epsilon_3 + \epsilon_4. \]  

(28)

It is useful to classify the positive roots according to their grading with respect to three relevant \( O(1,1) \) groups generated by the Cartan operators \( H_\beta, H_\lambda, H_\rho \) and parametrised by the moduli \( \beta \cdot h, h_4, h_6 \):

\[ O(1,1)_0 \rightarrow e^{\beta \cdot h} = V_4, \]
\[ O(1,1)_1 \rightarrow e^{h_4} = (V_4)^{1/2}(V_2)^{1/2}, \]
\[ O(1,1)_2 \rightarrow e^{h_6} = e^{-\phi}, \]

(29)

where we have denoted by \( \lambda^a \) the \( so(6,6) \) simple weights, \( \lambda^a \cdot \beta_m = \delta^a_m \). \( O(1,1)_0 \) is generated by the Cartan generator of the external \( SL(2, \mathbb{R}) \) and \( O(1,1)_1, O(1,1)_2 \) are in \( GL(4, \mathbb{R}) \times GL(2, \mathbb{R}) \), the former corresponding to the metric modulus given in eq. (27). In table 5 we list the axionic fields of the model together with the corresponding generator of \( Solv(sl(2, \mathbb{R})) + Solv(so(6,6)) \), for each of which the \( O(1,1)^3 \) grading and the \( SL(4, \mathbb{R}) \times SL(2, \mathbb{R}) \) representations are specified. The indices \( i, j \) and \( a, b \) label as usual the directions of the torus which are longitudinal \( (T_4) \) and transverse \( (T_2) \) to the D–branes.

The fields \( B_{ia} \) and \( C_{ia} \) transform in the representation \((4,4)\) of \( SL(4, \mathbb{R}) \times SO(2,2) \) where \( SO(2,2) = SL(2, \mathbb{R}) \times SL(2, \mathbb{R})_{IIB} \), and therefore will be collectively denoted by \( \Phi_i^\lambda \), where \( \lambda = (\alpha, a) = 1, 2, 3, 4 \) labels the \( 4 \) of \( SO(2,2) \), with a choice of basis corresponding to the invariant metric \( \eta_{\lambda\sigma} = \text{diag}(+1, +1, -1, -1) \). Its expression in terms of the fields \( B_{ia} \) and \( C_{ia} \) is

\[ \Phi_i^\lambda = \frac{1}{\sqrt{2}}\{C_{i2} - B_{i1}, B_{i2} + C_{i1}, B_{i1} + C_{i2}, -B_{i2} + C_{i1}\}. \]  

(30)

We shall use the same notation for the corresponding generators, \( \{T^i_{\lambda}\} \equiv \{T^{1ia}, T^{2ia}\} \).

From the assigned gradings one can conclude that the generators \( T_0, T^i_\lambda \) and \( T^{ij} \) close a 23–dimensional nilpotent solvable subalgebra \( N_7 \) of \( Solv(so(6,6)) \). The non–trivial
Table 5: Axionic fields for the $T_4 \times T_2$ IIB orientifold, generators of $\text{Solv}(so(6,6))$, $O(1,1)^3$ gradings, and $\text{SL}(4,\mathbb{R}) \times \text{SL}(2,\mathbb{R})$ representations.

| $GL(4) \times GL(2)$–rep. | generator | root | field | dim. |
|---------------------------|-----------|------|-------|------|
| $(1,1)_{(0,0,0)}$        | $T_{(0,0,0)}$               | $\{\epsilon_i - \epsilon_j, \epsilon_a - \epsilon_b\}$ $(i < j, a > b)$ | $\{g_{ij}, g_{ab}\}$ | 7 |
| $(1,1)_{(0,0,1)}$        | $T_0$ | $a$ | $C_0$ | 1 |
| $(4,2)_{(0,1,0)}$        | $T_{1ia}$ | $\epsilon_i + \epsilon_a$ | $B_{ia}$ | 8 |
| $(4,2)_{(0,1,1)}$        | $T_{2ia}$ | $a + \epsilon_i + \epsilon_a$ | $C_{ia}$ | 8 |
| $(6,1)_{(2,0,1)}$        | $T^{ij}$ | $a + (\epsilon_i + \epsilon_a) + (\epsilon_j + \epsilon_b)$ | $C_{ij\ ab} \equiv C_{ij} \epsilon_{ab}$ | 6 |
| $(1,1)_{(2,0,0)}$        | $T$ | $\beta = a + \epsilon_1 + \epsilon_2 + \epsilon_3 + \epsilon_4$ | $C_{ijkl} \equiv c$ | 1 |

Commutation relations are determined by the grading and the index structure of the generators, and read

$$
[T_0, T^i_j] = M_\lambda^\chi T^\chi_\lambda, \\
[T^i_j, T^j_\chi] = \eta_{\lambda\chi} T^i_\chi. 
$$

(31)

where $M_\lambda^\chi$ is a nilpotent generator acting on the 4 of $SO(2,2)$ which, for our choice of basis, can be cast in the form

$$
M_\lambda^\chi = \frac{1}{2} \begin{pmatrix}
0 & -1 & 0 & -1 \\
1 & 0 & 1 & 0 \\
0 & 1 & 0 & 1 \\
-1 & 0 & -1 & 0
\end{pmatrix}.
$$

(32)

**Infinitesimal transformations.** Let us consider now the infinitesimal transformations of the scalar fields generated by $T_0$, $T_{\lambda i}$ and $T_{ij}$. For simplicity we shall restrict our analysis to those points in the moduli space where the only non-vanishing scalars are $\Phi^\lambda_i$, $C^{ij}$ and $C$. The corresponding coset representative thus takes the simple form

$$
L = \exp \left( C_{ij} T^{ij} \right) \exp \left( \Phi^\lambda_i T^i_\lambda \right) \exp \left( C T_0 \right),
$$

(33)

and its associated left–invariant one–form is

$$
L^{-1} dL = (L^{-1} \partial_0 L) dC + (L^{-1} \partial_i L) d\Phi^\lambda_i + (L^{-1} \partial^{ij} L) dC_{ij} \\
= T_0 dC + d\Phi^\lambda_i (\delta^\lambda_\chi - C M_\lambda^\chi) T^i_\chi + \frac{1}{2} T^{ij} d\Phi^\lambda_\chi \Phi_\lambda j + T^{ij} dC_{ij}.
$$

(34)

In general, the action of an element $T_A$ on the coset representative can be expressed as:

$$
L^{-1} T_A L = k_\lambda^\chi L^{-1} \partial_\alpha L,
$$

(35)
where the $k_A$ are the corresponding Killing vectors. In the case at hand, from eq. (31), we can derive
\[
L^{-1}T_0L = T_0 + \Phi_i^\lambda M_\lambda^X T^i_X + \frac{1}{2} \Phi_i^\lambda \Phi_j^\mu M_\lambda^\mu T^{ij},
\]
\[
L^{-1}T^i_\lambda L = \left( \delta^\nu_\lambda - CM_\lambda^\nu \right) T^i_\nu + T^{ij} \Phi_j^\lambda,
\]
\[
L^{-1}T^{ij}L = T^{ij},
\]
and, thus, read the non–vanishing components of the Killing vectors
\[
k_0 = \partial_0 + \Phi_i^\lambda M_\lambda^X \partial^i_X,
k^i_\lambda = \partial^i_\lambda + \frac{1}{2} \Phi_j^\lambda \partial^{ij},
k^{ij} = \partial^{ij},
\]
where
\[
\partial_0 = \frac{\partial}{\partial C}, \quad \partial^{ij} = \frac{\partial}{\partial C_{ij}} \quad \text{and} \quad \partial^i_\lambda = \frac{\partial}{\partial \Phi_i^\lambda}.
\]
Therefore, under the infinitesimal diffeomorphism $\xi_0 k_0 + \xi_\lambda^i k^i_\lambda + \xi^{ij} k^{ij}$ the fields transform as follows:
\[
\delta C = \xi^0, \quad \delta \Phi_i^\lambda = \xi^i_\lambda + \xi^0 \Phi_i^\lambda M_\lambda^X, \quad \delta C_{ij} = \xi_{ij} + \frac{1}{2} \xi^i_\lambda \Phi_j^\lambda \mid \cdot \cdot.
\]
\textbf{Scalar kinetic terms.} Since all the quantities of our gauging are covariant with respect to $SO(2, 2) \times GL(4, \mathbb{R})$ it is useful to define the (full) coset representative in the following way
\[
L = \exp \left( C_{ij} T^{ij} \right) \exp \left( \Phi_i^\lambda T^i_\lambda \right) \exp (c T) \mathbb{E},
\]
where $\mathbb{E}$ is the coset representative of the submanifold
\[
\mathbb{E} \in O(1, 1)_0 \times \frac{SO(2, 2)}{SO(2) \times SO(2)} \times \frac{GL(4, \mathbb{R})}{SO(4)}.
\]
The scalar kinetic terms are computed by evaluating the components of the vielbein
\[
\mathcal{P} = L^{-1}dL|_{G/H}:
\]
\[
L^{-1}dL|_{G/H} = \mathcal{P}_{ij} \hat{T}^{ij} + \mathcal{P}_{i}^{\lambda} \hat{T}^{i}_\lambda + \mathcal{P} \hat{T} + \mathcal{P}_E,
\]
where the restriction to $G/H$ amounts to select the non–compact isometries of the scalar manifold, $\mathcal{P}_E$ is the algebra–valued vielbein of the submanifold (41). Finally, the hatted generators denote the non–compact component of the corresponding solvable generator. The kinetic Lagrangian for the scalar fields is then
\[
\mathcal{L}_{\text{scal}} = \frac{1}{2} \mathcal{P}_\mu \mathcal{P}^\mu + \frac{1}{2} \sum_{i\lambda} \mathcal{P}_{i\mu}^{\lambda} \mathcal{P}_{i}^{\lambda \mu} + \frac{1}{2} \sum_{ij} \mathcal{P}_{ij\mu} \mathcal{P}_{ij}^{\mu} + \text{Tr}(\mathcal{P}_E^2),
\]
where
\[
\mathcal{P}_\mu = \partial_\mu c, \quad \mathcal{P}^\lambda_{i\mu} = (\partial_\mu \Phi^\lambda_i) E^i_\lambda, \quad \mathcal{P}_{ij\mu} = \left[ \partial_\mu C_{ij} + \frac{1}{4} (\partial_\mu \Phi^\lambda_i \Phi^\lambda_j - \partial_\mu \Phi^\lambda_j \Phi^\lambda_i) \right] E^i_\lambda E^j_\lambda. \tag{44}
\]

Vector fields. The twelve vector potentials are $B_{a\mu}, C_{a\mu}, G^i_\mu, C_{ijk\mu}$. As before, we shall collectively denote by $A^\lambda_\mu$ the pair \{B_{a\mu}, C_{a\mu}\}, and by $F^\lambda = dA^\lambda$ the corresponding field strengths. To avoid confusion, we shall then adopt the following notation for the remaining field strengths: $\mathcal{F}^i = dG^i$ and $F^i = \epsilon^{ijkl} dC_{jkl}$. Moreover, $\tilde{F}_\lambda, \tilde{\mathcal{F}}_i$ and $\tilde{F}_i$ will denote the “dual” field strengths, obtained by varying the Lagrangian with respect to the electric ones, not to be confused with the four-dimensional Hodge duals $*F^\lambda, *\mathcal{F}^i$ and $*F^i$. Following [37], we can then collect the field strengths and their duals in a symplectic vector
\[
\{F^\lambda, \mathcal{F}^i, F^i, \tilde{F}_\lambda, \tilde{\mathcal{F}}_i, \tilde{F}_i\}. \tag{45}
\]

In table 6, we list the field strengths and their duals as they appear in the symplectic section, together with their $O(1,1)^3$ gradings and the corresponding weights of the 56 of $E_{7(7)}$.

| Sp–section | $O(1,1)^3$–grading | weight |
|------------|--------------------|--------|
| $F_{1a}$   | $(-1, 0, -\frac{1}{2})$ | $\epsilon_a - \frac{1}{\sqrt{2}} \epsilon_7$ |
| $F_{2a}$   | $(-1, 0, \frac{1}{2})$ | $w + \epsilon_a$ |
| $\mathcal{F}^i$ | $(-1, -1, -\frac{1}{2})$ | $-\epsilon_i - \frac{1}{\sqrt{2}} \epsilon_7$ |
| $F^i$      | $(1, -1, -\frac{1}{2})$ | $w + \epsilon_j + \epsilon_k + \epsilon_l$ |
| $\tilde{F}^{1a}$ | $(1, 0, \frac{1}{2})$ | $-\epsilon_a + \frac{1}{\sqrt{2}} \epsilon_7$ |
| $\tilde{F}^{2a}$ | $(1, 0, -\frac{1}{2})$ | $-w - \epsilon_a$ |
| $\tilde{\mathcal{F}}_i$ | $(1, 1, \frac{1}{2})$ | $\epsilon_i + \frac{1}{\sqrt{2}} \epsilon_7$ |
| $\tilde{F}_i$ | $(-1, 1, \frac{1}{2})$ | $-w - \epsilon_j - \epsilon_k - \epsilon_l$ |

Under a generic nilpotent transformation
\[
\xi T + \xi_0 T_0 + \xi^i T^i_\lambda + \xi_{ij} T^{ij}, \tag{46}
\]
the field strengths transform as
\[
\delta F^\lambda = -\xi^\lambda_i \mathcal{F}^i + \xi_0 F^{\lambda'} M_{\lambda'}^\lambda,
\]
\[
\begin{align*}
\delta \mathcal{F}^i &= 0, \\
\delta F^i &= \xi \mathcal{F}^i, \\
\delta \tilde{F}_\lambda &= \xi \eta_{\lambda\nu} F^\nu - \xi_0 M_\lambda^\nu \tilde{F}_\nu - \eta_{\lambda\nu} \xi_i^\nu F^i, \\
\delta \tilde{\mathcal{F}}_i &= -\xi \tilde{F}_i + \xi_\lambda \tilde{F}_\lambda - 2 \xi_{ij} F^j, \\
\delta \tilde{\mathcal{F}}_i &= -\xi_i^\nu F^\lambda \eta_{\lambda\nu} + 2 \xi_{ij} \mathcal{F}^j.
\end{align*}
\]

(47)

We then deduce that the electric subalgebra is

\[
g_e = o(1,1)_{(0,0)} + so(2,2)_{(0,0)} + gl(4,\mathbb{R})_{(0,0)} + (1,1)_{(2,0)} + (4,4)_{(0,1)} + (1,6)_{(0,2)},
\]

(48)

where \(o(1,1)_{(0,0)}\) is the generator of \(O(1,1)_0\), and the grading refers to \(O(1,1)_0 \times O(1,1)_1\). The group \(O(1,1)_2\) is now included inside \(SO(2,2)\) and, in what follows, we shall not consider its grading any longer. Furthermore, we identify \(T\) as the generator in \((1,1)_{(2,0)}\), \(T^i_\lambda\) and \(T^{ij}\) are associated to \((4,4)_{(0,1)}\) and \((6,1)_{(0,2)}\), respectively. The interested reader may find in appendix the explicit symplectic realisation of the generators of \(N_7\), as well as the computation of the vector kinetic matrix.

3.2. \(T_2 \times T_4\) IIB orientifold with D5 branes. The ungauged version

**Solvable algebra of global symmetries.** In this second model the relevant axions are \(B_{ia}, C_{ab}, C_{iabc} \equiv C^d_i, C_{\mu\nu} \equiv c\) and \(C_{ij} = \epsilon_{ij} c'\), and can be associated to the following choice of simple roots

\[
\begin{align*}
\beta_1 &= \epsilon_1 - \epsilon_2, \\
\beta_2 &= \epsilon_2 + \epsilon_3, \\
\beta_3 &= -\epsilon_3 + \epsilon_4, \\
\beta_4 &= -\epsilon_4 + \epsilon_5, \\
\beta_5 &= -\epsilon_5 + \epsilon_6, \\
\beta_6 &= a + \epsilon_3 + \epsilon_4,
\end{align*}
\]

for the subalgebra \(so(6,6) \subset e_{7(7)}\). The Dynkin diagram of the external \(SL(2,\mathbb{R})\) consists, instead, of the single root

\[
\beta = a + \epsilon_1 + \epsilon_2,
\]

(49)

whose corresponding axion is \(C_{ij}\), according to eq. (15).

The triple grading, this time, refers to the \(O(1,1)^3\) group generated by the three Cartan \(H_\beta, H_\lambda^2, H_\lambda^6\) and parametrised by the moduli \(\beta \cdot h, h_2, h_6\):

\[
\begin{align*}
O(1,1)_0 & \rightarrow e^{\beta h} = V_2 e^{-\phi_2}, \\
O(1,1)_1 & \rightarrow e^{h_2} = \left(V_2^\frac{1}{2} (V_4)^\frac{1}{4}\right)^{\frac{1}{4}} e^{\frac{\phi}{2}}, \\
O(1,1)_2 & \rightarrow e^{h_6} = (V_4)^\frac{1}{2} e^{-\frac{\phi}{2}},
\end{align*}
\]

(50)
where, as usual, $O(1,1)_0$ is in the external $SL(2,\mathbb{R})$, while $O(1,1)_1$ and $O(1,1)_2$ are contained in $GL(2,\mathbb{R}) \times GL(4,\mathbb{R})$.

In table 7 we list the axionic fields of this model, together with the corresponding generator of $Solv(so(6,6))$, for each of which the $O(1,1)^3$ grading is specified, as well as their $SL(2,\mathbb{R}) \times SL(4,\mathbb{R})$ representations.

Table 7: Axionic fields for the $T_2 \times T_4$ IIB orientifold, generators of $Solv(so(6,6))$, $O(1,1)^3$ gradings, and $SL(2,\mathbb{R}) \times SL(4,\mathbb{R})$ representations.

| $GL(2) \times GL(6)$–rep. | generator | root | field | dim. |
|----------------------------|-----------|------|-------|------|
| $-$                        | $T_{(0,0,0)}$ | $\{\epsilon_i - \epsilon_j, \epsilon_a - \epsilon_b\}$ $(i < j; a > b)$ | $\{g_{ij}, g_{ab}\}$ | 7 |
| $(1,6)_{(0,0,1)}$          | $T^{ab}$ | $\alpha_7 + \epsilon_a + \epsilon_b$ | $C_{ab}$ | 6 |
| $(2,4)_{(0,1,0)}$          | $T^{ia}$ | $\epsilon_i + \epsilon_a$ | $B_{ia}$ | 8 |
| $(2,4)_{(0,1,1)}$          | $T_d^{i}$ | $\alpha_7 + \epsilon_i + \epsilon_a + \epsilon_b + \epsilon_c$ | $C_d^i$ | 8 |
| $(1,1)_{(0,2,1)}$          | $T$ | $\alpha_7 + \epsilon_i + \epsilon_j + \epsilon_a + \epsilon_b + \epsilon_c + \epsilon_d$ | $C_{\mu\nu} = c$ | 1 |
| $(1,1)_{(2,0,0)}$          | $T'$ | $\beta$ | $C_{ij} = \epsilon'$ | 1 |

Also in this case, the generators $T$, $T^{ia}$, $T_d^i$ and $T^{ab}$ close a 23–dimensional solvable subalgebra of $SO(6,6)$

$$N_5 = c T + B_{ia} T^{ia} + C_i^a T_a^i + C_{ab} T^{ab},$$

whose algebraic structure is encoded in the non–vanishing commutators

$$[T^{ia}, T^{bc}] = \epsilon^{abcd} T_d^i,$$

$$[T^{ia}, T_d^j] = \epsilon^{ij} \delta_d^i T.$$  

The corresponding coset representative reads

$$L = e^{cT} e^{B_{ia} T^{ia}} e^{C_i^a T_a^i} e^{C_{ab} T^{ab}},$$

while its left–invariant one–form is

$$L^{-1} dL = T dc + T^{ab} dC_{ab} + T_d^i dC_i^d + (T^{ia} + \epsilon^{ij} C_j^a T + \epsilon^{abcd} T_d^i C_{bc} ) dB_{ia} .$$

The transformation properties of the axionic scalars can be deduced from

$$L^{-1} TL = T,$$

$$L^{-1} T_a^i L = T_a^i + \epsilon^{ij} B_{ja} T ,$$

$$L^{-1} T^{ia} L = T^{ia} + \epsilon^{ij} C_j^a T + \epsilon^{abcd} T_d^i C_{bc} ,$$

$$L^{-1} T^{ab} L = T^{ab} + \epsilon^{abcd} B_{id} T^i_c ,$$

17
which identify the Killing vectors

\[
\begin{align*}
k &= \partial, \\
k^i &= \partial^i + \epsilon^{ij} B_{ja} \partial, \\
k^{ia} &= \partial^{ia}, \\
k^{ab} &= \partial^{ab} + \epsilon^{abcd} B_{id} \partial^j, \\
\end{align*}
\]

(57)

where

\[
\begin{align*}
\partial &= \frac{\partial}{\partial c}, & \partial^j &= \frac{\partial}{\partial C^i}, & \partial^{ia} &= \frac{\partial}{\partial B_{ia}}, & \partial^{ab} &= \frac{\partial}{\partial C_{ab}}.
\end{align*}
\]

(58)

Hence, under the infinitesimal diffeomorphism \(\xi T + \xi_{ia} T^{ia} + \xi_{i} T^{i} + \xi_{ab} T^{ab}\), one has

\[
\begin{align*}
\delta c &= \epsilon^{ij} \xi_{i} B_{ja} + \xi, \\
\delta C^a_i &= \epsilon^{abcd} \xi_{bc} B_{id} + \xi^a_i, \\
\delta B_{ia} &= \xi_{ia}, \\
\delta C_{ab} &= \xi_{ab}.
\end{align*}
\]

(59)

For later convenience we shall define the generator \(T_{ab} = -\frac{1}{4} \epsilon_{abcd} T^{cd}\), and the corresponding parameter \(\xi_{ab} = -\frac{1}{4} \epsilon_{abcd} \xi^{cd}\), in terms of which the relation (52) reads

\[
[T_{ab}, T^{ic}] = \delta_{[a}^c T_{b]}^{i}.
\]

(60)

**Vector fields.** The vector fields of this model are \(G^{i}_{\mu}, C^{ia}_{i}, B_{a\mu}, C^{a}_{\mu}\), and we name the corresponding field strengths and their duals by

\[
\mathcal{F}^{i}_{\mu
\nu}, F_{i\mu
\nu}, \mathcal{H}^{a}_{\mu
\nu}, F^{a}_{\mu
\nu}, \tilde{\mathcal{F}}^{i}_{\mu
\nu}, \tilde{F}^{i}_{\mu
\nu}, \tilde{\mathcal{H}}^{a}_{\mu
\nu}, \tilde{F}^{a}_{\mu
\nu}.
\]

(61)

In the table 8 we list the field strengths and their duals as they appear in the symplectic section, together with their \(O(1, 1)^{3}\) gradings, and the corresponding \(E_7(7)\) weights.

The transformation laws under a generic nilpotent transformation \(\xi T + \xi_{ia} T^{ia} + \xi_{i} T^{i} + \xi_{ab} T^{ab}\) can be deduced from the grading and weight structures. One finds

\[
\begin{align*}
\delta \mathcal{F}^{i} &= 0, \\
\delta F_{i} &= \xi' \epsilon_{ij} \mathcal{F}^{j}, \\
\delta \mathcal{H}^{a} &= \xi_{ia} \mathcal{F}^{i}, \\
\delta F^{a} &= \xi^{ab} \mathcal{H}_{b} - \xi^{a}_{i} \mathcal{F}^{i}, \\
\delta \tilde{\mathcal{F}}^{i} &= \xi' \epsilon_{ij} \mathcal{F}^{j} + \xi_{i} \mathcal{F}_{a} - \xi_{ia} \mathcal{H}^{a} + \xi F_{i}, \\
\delta \tilde{F}^{i} &= \epsilon^{ij} \xi_{ja} F^{a} - \epsilon^{ij} \xi^{a}_{j} \mathcal{H}_{a} + \xi \mathcal{F}^{i}, \\
\delta \tilde{\mathcal{H}}^{a} &= \xi' F^{a} + \xi^{ab} \tilde{F}_{b} + \xi_{i} \epsilon^{ij} F_{j}, \\
\delta \tilde{F}^{a} &= \xi' \mathcal{H}^{a} - \xi_{ai} \epsilon^{ij} F_{j}.
\end{align*}
\]

(62)
Table 8: Field strengths, $O(1,1)^3$ gradings, and corresponding weights.

| Sp–section | $O(1,1)^3$–grading | weight          |
|------------|-------------------|----------------|
| $\mathcal{F}_{\mu\nu}$ | $(-1,-1,-\frac{1}{2})$ | $-\epsilon_i - \frac{1}{\sqrt{2}}\epsilon_7$ |
| $F_{\mu\nu}$ | $(1,-1,-\frac{1}{2})$ | $w + \epsilon_i$ |
| $\mathcal{H}_{a\mu\nu}$ | $(-1,0,-\frac{1}{2})$ | $\epsilon_a - \frac{1}{\sqrt{2}}\epsilon_7$ |
| $F_{a\mu\nu}$ | $(-1,0,\frac{1}{2})$ | $w + \epsilon_b + \epsilon_c + \epsilon_d$ |
| $\mathcal{F}_{\mu\nu}$ | $(1,1,\frac{1}{2})$ | $\epsilon_i + \frac{1}{\sqrt{2}}\epsilon_7$ |
| $\tilde{F}_{\mu\nu}$ | $(-1,1,\frac{1}{2})$ | $w + \epsilon_j + \epsilon_a + \epsilon_b + \epsilon_c + \epsilon_d$ |
| $\mathcal{H}_{\mu\nu}$ | $(1,0,\frac{1}{2})$ | $-\epsilon_a + \frac{1}{\sqrt{2}}\epsilon_7$ |
| $\tilde{F}_{a\mu\nu}$ | $(1,0,-\frac{1}{2})$ | $w + \epsilon_i + \epsilon_j + \epsilon_a$ |

The explicit symplectic representation of the $N_5$ generators together with the computation of the vector kinetic matrix $\mathcal{N}$ may be found in appendix.

3.3. $T_0 \times T_0$ and $T_6 \times T_0$ IIB orientifolds with D3 and D9 branes. The ungauged version.

The $T_0 \times T_6$ model in the presence of D3–branes, with and without fluxes was constructed in [32, 30, 31]. The structure of the $T_6 \times T_0$ model, on the other hand, is somewhat trivial, since there is no room for fluxes to be turned on. For completeness, here we shall confine ourselves to the description of their embeddings within the $\mathcal{N} = 8$ theory, and to the identification of the solvable algebras $N_3$ and $N_9$, together with their action on scalar and vector fields.

Solvable algebra of global symmetries: the $T_0 \times T_0$ model. The embedding of the $\mathfrak{sl}(2,\mathbb{R}) + \mathfrak{so}(6,6)$ algebra inside $e_{7(7)}$ is defined by the following identification of the simple roots:

\[
\begin{align*}
\beta_1 &= -\epsilon_1 + \epsilon_2, \\
\beta_2 &= -\epsilon_2 + \epsilon_3, \\
\beta_3 &= -\epsilon_3 + \epsilon_4, \\
\beta_4 &= -\epsilon_4 + \epsilon_5, \\
\beta_5 &= -\epsilon_5 + \epsilon_6, \\
\beta_6 &= a + \epsilon_1 + \epsilon_2 + \epsilon_3 + \epsilon_4,
\end{align*}
\]

for the $\mathfrak{so}(6,6)$ component, and

\[
\beta = a,
\]

(63)
for the $sl(2, \mathbb{R})$ one. The correspondence axion–root is quite simple and is summarised in table 9.

Table 9: Axionic fields for the $T_0 \times T_6$ IIB orientifold, generators of $Solv(so(6, 6))$, $O(1, 1)^2$ gradings and $GL(6, \mathbb{R})$ representations.

| $GL(6)$–rep. | generator | root | field | dim. |
|--------------|-----------|------|-------|------|
| $-15(0,0)$   | $T_{(0,0)}$ | $\{\epsilon_a - \epsilon_b\}$ ($a > b$) | $\{g_{ab}\}$ | 15   |
| $15(0,1)$    | $T_{ab}$   | $a + \epsilon_c + \epsilon_d + \epsilon_e + \epsilon_f$ | $C^{ab} \equiv \epsilon^{abcdef} C_{cdef}$ | 15   |
| $1(2,0)$     | $T$        | $\beta$ | $C_0 = c$ | 1    |

In this case, the grading is with respect to the pair of $O(1, 1)$ groups generated by $H_{\beta}, H_{\lambda}$ and corresponding to the following moduli:

$$
O(1, 1)_0 \rightarrow e^{\beta_h} = e^{-\phi},
O(1, 1)_1 \rightarrow e^{\lambda_6 h} = V_6.
$$

The nilpotent algebra $N_3$, generated by $T_{ab}$, acts as Peccei–Quinn translations on the R–R scalars $C^{ab}$

$$
\delta C^{ab} = \xi^{ab}.
$$

The vector fields are $C_{a\mu}$ and $B_{a\mu}$, and the symplectic section of the corresponding field strengths $F_{a\mu}$ and $H_{a\mu}$ and their magnetic duals $\tilde{F}^a_{\mu\nu}$, $\tilde{H}^a_{\mu\nu}$ is listed in table 10.

Table 10: Field strengths, $O(1, 1)^2$ gradings, and corresponding weights.

| Sp–section | $O(1, 1)^2$–grading | weight |
|------------|---------------------|--------|
| $F_{a\mu}$ | $(1, -\frac{1}{2})$ | $w + \epsilon_a$ |
| $H_{a\mu}$ | $(-1, -\frac{1}{2})$ | $\epsilon_a - \frac{1}{\sqrt{2}} \epsilon_7$ |
| $F^a_{\mu\nu}$ | $(-1, \frac{1}{2})$ | $-w - \epsilon_a$ |
| $H^a_{\mu\nu}$ | $(1, \frac{1}{2})$ | $-\epsilon_a + \frac{1}{\sqrt{2}} \epsilon_7$ |

The duality action of an infinitesimal transformation $\xi^{ab} T_{ab} + \xi T$ is then

$$
\delta F_a = \xi H_a,
$$

20
\[\begin{align*}
\delta \mathcal{H}_a &= 0, \\
\delta \tilde{F}^a &= \xi^{ab} \mathcal{H}_b, \\
\delta \tilde{\mathcal{H}}^a &= -\xi^{ab} F_b - \xi \tilde{F}^a. 
\end{align*}\] (67)

**Solvable algebra of global symmetries: the \(T_6 \times T_0\) model.** The embedding of the \(sl(2, \mathbb{R})^+\) \(so(6, 6)\) algebra inside \(e_{7(7)}\) is defined by the following identification of the simple roots

\[\begin{align*}
\beta_1 &= \epsilon_1 - \epsilon_2, \\
\beta_2 &= \epsilon_2 - \epsilon_3, \\
\beta_3 &= \epsilon_3 - \epsilon_4, \\
\beta_4 &= \epsilon_4 - \epsilon_5, \\
\beta_5 &= \epsilon_5 - \epsilon_6, \\
\beta_6 &= a + \epsilon_5 + \epsilon_6, 
\end{align*}\] (68)

for the \(so(6, 6)\) component, and

\[\beta = a + \sum_{n=1}^{6} \epsilon_n,\] (69)

for the \(sl(2, \mathbb{R})\) one. The correspondence axion–root is quite simple, and is summarised in table 11.

**Table 11: Axionic fields for the \(T_6 \times T_0\) IIB orientifold, generators of \(Solv(so(6, 6))\), \(O(1, 1)^2\) gradings, and \(GL(6, \mathbb{R})\) representations.**

| GL(6)-rep. | generator | root | field | dim. |
|-----------|-----------|------|-------|------|
| \(T_{(0,0)}\) | \(\{\epsilon_i - \epsilon_j\} \ (i < j)\) | \(\{g_{ij}\}\) | 15 |
| \(15_{(0,1)}\) | \(T^{ij}\) | \(a + \epsilon_i + \epsilon_j\) | \(C_{ij}\) | 15 |
| \(1_{(2,0)}\) | \(T\) | \(\beta\) | \(C_{\mu\nu} = c\) | 1 |

In this case, the grading is with respect to a pair of \(O(1, 1)\) groups generated by \(H_\beta, H_{\lambda}\) and corresponding to the following moduli:

\[\begin{align*}
O(1, 1)_0 \rightarrow e^{\beta \cdot h} &= V_0 e^{\frac{\phi}{2}}, \\
O(1, 1)_1 \rightarrow e^{\lambda \cdot h} &= (V_0)^{\frac{1}{2}} e^{-\frac{3}{4} \phi}. 
\end{align*}\] (70)

The nilpotent algebra \(N_9\), generated by \(T^{ij}\), acts as Peccei–Quinn translations on the R–R scalars \(C_{ij}\),

\[\delta C_{ij} = \xi_{ij}.\] (71)
The vector fields are $C_{i\mu}$ and $G_{\mu}^i$, and the symplectic sections of the corresponding field strengths $F_{i\mu\nu}$ and $F^i_{\mu\nu}$ and their magnetic duals $\tilde{F}_{i\mu\nu}$, $\tilde{F}^i_{\mu\nu}$ are listed in table 12.

Table 12: Field strengths, $O(1, 1)^2$ gradings, and corresponding weights.

| Sp−section | $O(1, 1)^2$−grading | weight          |
|------------|---------------------|-----------------|
| $F_{i\mu\nu}$ | $(-1, \frac{1}{2})$ | $w + \epsilon_i$ |
| $F^i_{\mu\nu}$ | $(-1, -\frac{1}{2})$ | $-\epsilon_i - \frac{1}{\sqrt{2}}\epsilon_7$ |
| $\tilde{F}_{i\mu\nu}$ | $(1, -\frac{1}{2})$ | $-w - \epsilon_i$ |
| $\tilde{F}^i_{\mu\nu}$ | $(1, \frac{1}{2})$ | $\epsilon_i + \frac{1}{\sqrt{2}}\epsilon_7$ |

The duality action of an infinitesimal transformation $\xi_{ij} T^{ij} + \xi T$ is then

$$\delta F_i = \xi_{ij} F^j,$$
$$\delta F^i = 0,$$
$$\delta \tilde{F}^i = \xi F^i,$$
$$\delta \tilde{F}_i = \xi_{ij} F^j + \xi F_i.$$

(72)

As a result, the electric group contains the whole $SO(6, 6)$, as for the heterotic string on $T_6$. In other words, there are no Peccei–Quinn isometries in $SO(6, 6)$ which could be gauged. This feature is consistent with the fact that this model does not allow fluxes, and usually fluxes translate into local Peccei–Quinn invariances in the low–energy supergravity description.

3.4. $T_1 \times T_5$ IIA orientifold with D4–branes.

**Solvable algebra of global symmetries.** The embedding of the $sl(2, \mathbb{R}) + so(6, 6)$ algebra inside $e_{7(7)}$ is defined by the following identifications of simple roots:

$$\beta_1 = \epsilon_1 + \epsilon_2,$$
$$\beta_2 = -\epsilon_2 + \epsilon_3,$$
$$\beta_3 = -\epsilon_3 + \epsilon_4,$$
$$\beta_4 = -\epsilon_4 + \epsilon_5,$$
$$\beta_5 = -\epsilon_5 + \epsilon_6,$$
$$\beta_6 = a + \epsilon_2 + \epsilon_3 + \epsilon_4.$$  

(73)
for the $so(6,6)$ factor, and
\[ \beta = \alpha + \epsilon_1, \quad (74) \]
for the $sl(2,\mathbb{R})$ one. The correspondence axion–root is quite simple, and is summarised in table 13.

Table 13: Axionic fields for the $T_1 \times T_5$ IIA orientifold, generators of $Solv(so(6,6))$, $O(1,1)^3$ gradings, and $GL(5,\mathbb{R})$ representations.

| $GL(5)$–rep. | generator | root | field | dim. |
|--------------|-----------|------|-------|------|
| $-$          | $T_{(0,0,0)}$ | $\{\epsilon_a - \epsilon_b\}$ ($a > b$) | $\{g_{ab}\}$ | 10   |
| $\overline{10}_{(0,0,1)}$ | $T_{ab}$ | $a + \epsilon_c + \epsilon_d + \epsilon_e$ | $C_{cde} \equiv C^{ab}$ | 10   |
| $5_{(0,1,0)}$ | $T^a$ | $\epsilon_1 + \epsilon_a$ | $B_{1a} \equiv B_a$ | 5    |
| $\overline{5}_{(0,1,1)}$ | $T_e$ | $a + \epsilon_1 + \epsilon_a + \epsilon_b + \epsilon_e + \epsilon_d$ | $C_{\mu\nu a} \equiv C^e$ | 5    |
| $1_{(2,0,0)}$ | $T$ | $\beta$ | $C_1 \equiv c$ | 1    |

In this case the grading is with respect to the $O(1,1)^3$ group generated by $H_\beta$, $H_{\lambda 1}$, $H_{\lambda 6}$ and parametrised by the moduli $\beta \cdot h$, $h_1$, $h_6$:
\begin{align*}
O(1,1)_0 & \rightarrow e^{\beta \cdot h} = V_1 e^{-\frac{3}{4} \phi}, \\
O(1,1)_1 & \rightarrow e^{h_1} = V_1 (V_5)^\frac{1}{2} e^{\frac{\phi}{4}}, \\
O(1,1)_2 & \rightarrow e^{h_6} = (V_5)^\frac{3}{2} e^{-\frac{\phi}{4}}. \quad (75)
\end{align*}

The generators $T^a$, $T_a$ and $T_{ab}$ close a twenty–dimensional nilpotent subalgebra $N_4$ of $Solv(so(6,6))$:
\[ N_4 = B_a T^a + C^a T_a + C^{ab} T_{ab}, \quad (76) \]
whose algebraic structure is encoded in the non–vanishing commutator
\[ [T_{ab}, T^c] = T_{[a} \delta_{b]}^c. \quad (77) \]

The corresponding coset representative reads
\[ L = e^{C^a T_a} e^{B_a T^a} e^{C^{ab} T_{ab}} e^{cT} E, \quad (78) \]
where the $E$ factor parametrises the submanifold:
\[ O(1,1)_0 \times O(1,1)_1 \times O(1,1)_2 \times \frac{SL(5,\mathbb{R})}{SO(5)}. \quad (79) \]
A generic element $\xi_a T^a + \xi^a T_a + \xi^{ab} T_{ab}$ of $N_4$ then induces the following transformations on the axionic scalars

$$
\begin{align*}
\delta C^a &= \xi^a + \xi^{ab} B_b, \\
\delta B_a &= \xi_a, \\
\delta C^{ab} &= \xi^{ab}.
\end{align*}
$$

Vector fields. The vector fields of this model are $C_\mu, C^1_\mu, C_{1a\mu}, B_{a\mu}$, and we name the corresponding field strengths $F_{\mu\nu}$, $F^1_{\mu\nu}$, $F_{1a\mu\nu}$, $\mathcal{H}_{a\mu\nu}$. The symplectic section of the field strengths and their duals is

$$
\{F_{\mu\nu}, F^1_{\mu\nu}, F_{1a\mu\nu}, \mathcal{H}_{a\mu\nu}, \tilde{F}_{\mu\nu}, \tilde{F}^1_{\mu\nu}, \tilde{\mathcal{H}}_{a\mu\nu}\},
$$

and in table 14 we give their $O(1, 1)^3$ gradings and the corresponding $E_{7(7)}$ weights.

Table 14: Field strengths, $O(1, 1)^3$ gradings, and corresponding weights.

| vector         | $O(1, 1)^3$-grading | weight          |
|----------------|---------------------|-----------------|
| $F_{\mu\nu}$  | $(-1, -1, -\frac{1}{2})$ | $w$             |
| $F^1_{\mu\nu}$| $(-1, -1, -\frac{1}{2})$ | $-\epsilon_1 - \frac{1}{\sqrt{2}} \epsilon_7$ |
| $F_{1a\mu\nu}$| $(1, 0, -\frac{1}{2})$ | $w + \epsilon_1 + \epsilon_a$ |
| $\mathcal{H}_{a\mu\nu}$ | $(-1, 0, -\frac{1}{2})$ | $\epsilon_a - \frac{1}{\sqrt{2}} \epsilon_7$ |
| $\tilde{F}_{\mu\nu}$ | $(-1, 1, \frac{1}{2})$ | $-w$           |
| $\tilde{F}^1_{\mu\nu}$ | $(1, 1, \frac{1}{2})$ | $\epsilon_1 + \frac{1}{\sqrt{2}} \epsilon_7$ |
| $F^{1a}_{\mu\nu}$ | $(-1, 0, \frac{1}{2})$ | $-w - \epsilon_1 - \epsilon_a$ |
| $\tilde{\mathcal{H}}_{a\mu\nu}$ | $(1, 0, \frac{1}{2})$ | $-\epsilon_a + \frac{1}{\sqrt{2}} \epsilon_7$ |

The action of infinitesimal duality transformation $\xi_a T^a + \xi^a T_a + \xi^{ab} T_{ab} + \xi T$ on the symplectic section is

$$
\begin{align*}
\delta F &= \xi F^1, \\
\delta F^1 &= 0, \\
\delta F_{1a} &= \xi_a F + \xi \mathcal{H}_a, \\
\delta \mathcal{H}_a &= \xi_a F^1, \\
\delta \tilde{F} &= -\xi_a \tilde{F}^1 + \xi^a \mathcal{H}_a, \\
\delta \tilde{F}_1 &= -\xi_a \tilde{\mathcal{H}}^a - \xi^a F_{1a} - \xi \tilde{F},
\end{align*}
$$

24
\[ \delta \tilde{F}^{1a} = -\xi^a \mathcal{F}^1 - \xi^{ab} \mathcal{H}_b, \]
\[ \delta \tilde{\mathcal{H}}^a = \xi^a F + \xi^{ab} F_{1b} - \xi \tilde{F}^{1a}. \] (82)

The explicit symplectic realisation of the $N_4$ generators together with the computation of the vector kinetic matrix can be found in appendix.

3.5. $T_3 \times T_3$ IIA orientifold with $D6$–branes.

**Solvable algebra of global symmetries.** The embedding of the $sl(2, \mathbb{R}) + so(6,6)$ algebra inside $e_7(7)$ is defined by the following identification of the simple roots

\[ \beta_1 = \epsilon_1 - \epsilon_2, \]
\[ \beta_2 = \epsilon_2 - \epsilon_3, \]
\[ \beta_3 = \epsilon_3 + \epsilon_4, \]
\[ \beta_4 = -\epsilon_4 + \epsilon_5, \]
\[ \beta_5 = -\epsilon_5 + \epsilon_6, \]
\[ \beta_6 = a + \epsilon_4, \] (83)

for the $so(6,6)$ factor, and

\[ \beta = a + \epsilon_1 + \epsilon_2 + \epsilon_3, \] (84)

for the $sl(2, \mathbb{R})$ one. The correspondence axion–root is quite simple and is summarised in table 15.

**Table 15:** Axionic fields for the $T_3 \times T_3$ IIA orientifold, generators of $Solv(so(6,6))$, $O(1,1)^3$ gradings, and $GL(3, \mathbb{R}) \times GL(3, \mathbb{R})$ representations.

| $GL(3) \times GL(3)$–rep. | generator | root | field | dim. |
|--------------------------|-----------|------|-------|------|
| —                        | $T_{(0,0,0)}$ | $\{\epsilon_i - \epsilon_j, \epsilon_a - \epsilon_b\}$ ($i < j$, $a > b$) | $\{g_{ij}, g_{ab}\}$ | 6    |
| (1, 3) | $T_{(0,1,0)}$ | $T^{ia}$ | $\epsilon_i + \epsilon_a$ | $B_{ia}$ | 9    |
| (3, 3) | $T_{(0,1,1)}$ | $T^{i}$ | $a + \epsilon_i + \epsilon_b + \epsilon_c$ | $C_{ibc} \equiv C_i^a$ | 9    |
| (3, 3) | $T_{(0,2,1)}$ | $T^{ij}$ | $\epsilon_i + \epsilon_j + \epsilon_a + \epsilon_b + \epsilon_c$ | $C_{k\mu\nu} \equiv C_{ij}$ | 3    |
| (1, 1) | $T_{(2,0,0)}$ | $T$ | $\beta$ | $C_{ijk} \equiv c$ | 1    |

The triple grading refers to three $O(1,1)$ groups generated by $H_\beta, H_{\lambda^3}, H_{\lambda^6}$ and parametrised by the moduli $\beta \cdot h$, $h_3$, $h_6$:

\[ O(1,1)_0 \to e^{\beta \cdot h} = V_3 e^{-\frac{\phi}{4}}, \]

25
\[ O(1,1)_1 \to e^{h_3} = (V'_{3})^{1/3} V'_{3}^{1/3} e^{\frac{\phi}{2}} , \\
O(1,1)_2 \to e^{h_6} = (V'_{3})^{1/3} e^{-\frac{\phi}{4}} . \] (85)

The generators \( T^{ia} \), \( T_{ab} \), \( T^i_a \) and \( T^{ij} \) form now a 24–dimensional solvable subalgebra \( N_6 \) of \( \text{Solv}(so(6,6)) \):

\[ N_6 = B^{ia} T^{ia} + C^{ab} T_{ab} + C^a_i T^i_a + C_{ij} T^{ij} , \] (86)
whose algebraic structure is encoded in the non–vanishing commutators

\[ [T_{ab}, T^{ic}] = T^i_a \delta^c_b , \\
[T_{ia}, T^j_b] = T^{ij} \delta^a_b . \] (87)

A possible choice for the coset representative is then

\[ L = e^{C_{ij} T^{ij}} e^{C^i_a T^i_a} e^{B^{ia} T^{ia}} e^{C^{ab} T_{ab}} e^{E} E , \] (88)
with \( E \) parameterising the submanifold

\[ O(1,1)_0 \times \frac{GL(3, \mathbb{R})}{SO(3)} \times \frac{GL(3, \mathbb{R})}{SO(3)} . \] (89)

Under an infinitesimal transformation \( \xi_{ij} T^{ij} + \xi^a_i T^i_a + \xi_{ia} T^{ia} + \xi^{ab} T_{ab} \) of \( N_6 \) the variation of the axionic scalars is

\[ \delta C^a_i = \xi^a_i + \xi^{ab} B_{ab} , \\
\delta C_{ij} = \xi_{ij} + \xi_{a[i} C^a_{j]} , \\
\delta B_{ia} = \xi_{ia} , \\
\delta C^{ab} = \xi^{ab} . \] (90)

**Vector fields.** The vector fields of this model are \( G^i_{\mu} C^i_{\mu} = \epsilon^{ijk} C_{j\mu} \), \( B^{a\mu} \), \( C^a_{\mu} = \epsilon^{abc} C_{b\mu} \), and we name the corresponding field strengths \( \tilde{F}^{i\mu\nu}, F^{i\mu\nu}, \mathcal{H}_{a\mu\nu}, F^{a\mu\nu} \). The symplectic section of the field strengths and their duals is

\[ \{ \tilde{F}^{i\mu\nu}, F^{i\mu\nu}, \mathcal{H}_{a\mu\nu}, F^{a\mu\nu} \} , \] (91)
and in table 16 we give their \( O(1,1) \times \) gradings and the corresponding \( E_{7(7)} \) weights.

The action of an infinitesimal duality transformation \( \xi_{ij} T^{ij} + \xi^a_i T^i_a + \xi_{ia} T^{ia} + \xi^{ab} T_{ab} + \xi T \) on the symplectic section is

\[ \delta \tilde{F}^i = 0 , \\
\delta F^i = \xi \tilde{F}^i , \]
Table 16: Field strengths, $O(1, 1)^3$ gradings, and corresponding weights.

| Sp–section | $O(1, 1)^3$–grading | weight            |
|------------|----------------------|------------------|
| $\mathcal{F}^i_{\mu\nu}$ | $(-1, -1, -\frac{1}{2})$ | $-\epsilon_i - \frac{1}{\sqrt{2}}\epsilon_7$ |
| $F^i_{\mu\nu}$ | $(1, -1, -\frac{1}{2})$ | $w + \epsilon_j + \epsilon_k$ |
| $\mathcal{H}^a_{\alpha\mu\nu}$ | $(-1, 0, -\frac{1}{2})$ | $\epsilon_a - \frac{1}{\sqrt{2}}\epsilon_7$ |
| $F^a_{\mu\nu}$ | $(-1, 0, \frac{1}{2})$ | $w + \epsilon_b + \epsilon_c$ |
| $\mathcal{F}^a_{i\mu\nu}$ | $(1, 1, \frac{1}{2})$ | $\epsilon_i + \frac{1}{\sqrt{2}}\epsilon_7$ |
| $\tilde{F}^i_{\mu\nu}$ | $(-1, 1, \frac{1}{2})$ | $-w - \epsilon_j - \epsilon_k$ |
| $\tilde{\mathcal{H}}^a_{\mu\nu}$ | $(1, 0, \frac{1}{2})$ | $-\epsilon_a + \frac{1}{\sqrt{2}}\epsilon_7$ |
| $F^a_{\mu\nu}$ | $(1, 0, -\frac{1}{2})$ | $-w - \epsilon_b - \epsilon_c$ |

\[ \begin{align*}
\delta \mathcal{H}_a &= \xi_{ia} \mathcal{F}^i, \\
\delta F^a &= \xi_{ab} \mathcal{H}_b + \xi^a_i \mathcal{F}^i, \\
\delta \tilde{\mathcal{F}}_i &= -\xi_{ia} \tilde{\mathcal{H}}^a - \xi^a_i \tilde{F}_a - 2\xi_{ij} F^j - \xi \tilde{F}_i, \\
\delta \tilde{F}_i &= \xi_{ia} F^a + \xi^a_i \mathcal{H}_a + 2\xi_{ij} \mathcal{F}^j, \\
\delta \tilde{\mathcal{H}}^a &= \xi_{ab} \tilde{F}_a + \xi^a_i \tilde{F}^i + \xi F^a, \\
\delta \tilde{F}^a &= \xi_{ia} F^i + \xi \mathcal{H}_a. \end{align*} \]

The explicit symplectic realisation of the $N_6$ generators together with the computation of the vector kinetic matrix can be found in appendix.

3.6. $T_5 \times T_1$ IIA orientifold with $D8$–branes

**Solvable algebra of global symmetries.** The embedding of the $sl(2, \mathbb{R}) + so(6, 6)$ algebra inside $e_{7(7)}$ is defined by the following identification of the simple roots

\[ \begin{align*}
\beta_1 &= \epsilon_1 - \epsilon_2, \\
\beta_2 &= \epsilon_2 - \epsilon_3, \\
\beta_3 &= \epsilon_3 - \epsilon_4, \\
\beta_4 &= \epsilon_4 - \epsilon_5, \\
\beta_5 &= \epsilon_5 + \epsilon_6, \\
\beta_6 &= a + \epsilon_5. \end{align*} \]
for the \( so(6, 6) \) factor, and

\[
\beta = a + \epsilon_1 + \epsilon_2 + \epsilon_3 + \epsilon_4 + \epsilon_5, \tag{94}
\]

for the \( sl(2, \mathbb{R}) \) one. The correspondence axion–root is quite simple and is summarised in table 17.

Table 17: Axionic fields for the \( T_5 \times T_1 \) IIA orientifold, generators of \( Solv(so(6, 6)) \), \( O(1, 1)^3 \) gradings, and \( GL(5, \mathbb{R}) \) representations.

| \( GL(5) \)-rep. | generator | root          | field          | dim. |
|------------------|-----------|---------------|----------------|------|
| \( - \)          | \( T_{(0,0,0)} \) | \( \{\epsilon_i - \epsilon_j, \epsilon_a - \epsilon_b\} \) \( (i < j) \) | \( \{g_{ij}\} \) | 10   |
| \( 5_{(0,1,0)} \) | \( T_i \)   | \( a + \epsilon_i \) | \( C_i \) | 5    |
| \( 5_{(1,0,0)} \) | \( T^i \)   | \( \epsilon_i + \epsilon_6 \) | \( B_{i6} \equiv B_i \) | 5    |
| \( 10_{(0,1,1)} \) | \( T^{ij} \) | \( a + \epsilon_i + \epsilon_j + \epsilon_6 \) | \( C_{ij6} \equiv C_{ij} \) | 10   |
| \( 1_{(2,0,0)} \) | \( T \)      | \( \beta \)    | \( C_{\mu\nu 6} \equiv c \) | 1    |

The triple grading refers to three \( O(1, 1) \) groups generated by \( H_\beta, H_{\lambda^5}, H_{\lambda^6}, \) all commuting with \( SL(5, \mathbb{R}) \), and parametrised by the moduli \( \beta \cdot h, h_5, h_6 \):

\[
\begin{align*}
O(1, 1)_0 & \rightarrow e^{\beta \cdot h} = V_5 e^{\frac{\phi}{\sqrt{5}}}, \\
O(1, 1)_1 & \rightarrow e^{h_5} = (V_5)^\frac{1}{2} V_1 e^{\frac{\phi}{\sqrt{5}}}, \\
O(1, 1)_2 & \rightarrow e^{h_6} = (V_5)^\frac{1}{2} e^{-\frac{3}{2} \phi}.
\end{align*}
\tag{95}
\]

The generators \( T^{\alpha i} \), \( T^i \) and \( T^{ij} \) form now a twenty–dimensional solvable subalgebra \( N_8 \) of \( Solv(so(6, 6)) \):

\[
N_8 = B_{i6} T^{\alpha \alpha} + C_i T^i + C_{ij} T^{ij}, \tag{96}
\]

whose algebraic structure is encoded in the non–vanishing commutator

\[
[T^i, T^{ij}] = T^{ij}. \tag{97}
\]

A possible choice for the the coset representative is then

\[
L = e^{C_{ij} T^{ij}} e^{B_{\alpha i} T^{\alpha i}} e^{C_i T^i} e^{c \beta} E,
\tag{98}
\]

with the \( E \) parameterising the submanifold:

\[
O(1, 1)_0 \times O(1, 1)_1 \times O(1, 1)_2 \times \frac{SL(5, \mathbb{R})}{SO(5)}. \tag{99}
\]
Under an infinitesimal transformation $\xi_{ij} T^{ij} + \xi_i T^i + \xi'_i T^{ii} + \xi T$ of $N_8$ the variation of the axionic scalars is

$$
\delta C_{ij} = \xi_i [B_{j6}] + \xi_{ij}, \\
\delta B_{i6} = \xi'_i, \\
\delta C_i = \xi_i.
$$

(100)

**Vector fields.** The vector fields of this model are $G^i_\mu$, $C^i_{6\mu}$, $C_\mu$, $B_{a\mu}$, and we name the corresponding field strengths $\mathcal{F}_{i\mu}$, $F_{i6\mu}$, $F_{i\mu}$, $\mathcal{H}_{6\mu\nu}$. The symplectic section of the field strengths and their duals is

$$
\{\mathcal{F}_{\mu\nu}, F_{i6\mu\nu}, F_{i\mu\nu}, \mathcal{H}_{6\mu\nu}, \tilde{\mathcal{F}}_{i\mu\nu}, \tilde{F}_{i\mu\nu}, \tilde{\mathcal{H}}_{i\mu\nu}\},
$$

(101)

and in table 18 we give their $O(1,1)^3$ gradings and the corresponding $E_{7(7)}$ weights.

| Sp–section | $O(1,1)^3$–grading | weight |
|------------|-------------------|--------|
| $\mathcal{F}_{\mu\nu}$ | $(-1, -1, -\frac{1}{2})$ | $-\epsilon_i - \frac{1}{\sqrt{2}} \epsilon_7$ |
| $F_{i6\mu\nu}$ | $(-1, 1, \frac{1}{2})$ | $w + \epsilon_i + \epsilon_6$ |
| $F_{i\mu\nu}$ | $(-1, -1, \frac{1}{2})$ | $w$ |
| $\mathcal{H}_{6\mu\nu}$ | $(-1, 1, -\frac{1}{2})$ | $\epsilon_6 - \frac{1}{\sqrt{2}} \epsilon_7$ |
| $\tilde{\mathcal{F}}_{i\mu\nu}$ | $(1, 1, \frac{1}{2})$ | $\epsilon_i + \frac{1}{\sqrt{2}} \epsilon_7$ |
| $\tilde{F}_{i\mu\nu}$ | $(1, -1, -\frac{1}{2})$ | $-w - \epsilon_i - \epsilon_6$ |
| $\tilde{\mathcal{H}}_{i\mu\nu}$ | $(1, -1, \frac{1}{2})$ | $-w + \frac{1}{\sqrt{2}} \epsilon_7$ |

The action of an infinitesimal transformation $\xi_{ij} T^{ij} + \xi_i T^i + \xi'_i T^{ii} + \xi T$ on the symplectic section is

$$
\delta \mathcal{F}^i = 0, \\
\delta F_{i6} = \xi_i \mathcal{H}_6 - \xi'_i F + \xi_{ij} \mathcal{F}^j, \\
\delta F = -\xi_i \mathcal{F}^i, \\
\delta \mathcal{H}_6 = \xi'_i \mathcal{F}^i, \\
\delta \tilde{\mathcal{F}}_i = \xi_i \tilde{F} - \xi'_i \mathcal{H}^6 + \xi_{ij} \tilde{F}^j + \xi \tilde{F}_{i6}, \\
\delta \tilde{F}_{i6} = \xi \mathcal{F}^i, \\
\delta \tilde{\mathcal{H}}_6 = \xi \tilde{F}^i.
$$
\[ \delta \tilde{F} = \xi' \tilde{F} + \xi \mathcal{H}_6, \]
\[ \delta \mathcal{H}_6 = -\xi' \tilde{F} + \xi F. \] (102)

The explicit symplectic realisation of the $N_8$ generators, together with the computation of the vector kinetic matrix can be found in appendix.

4. Fluxes and gauged supergravity: local Peccei–Quinn symmetry as gauged duality transformations.

In the present section we consider the deformation of $\mathcal{N} = 4$ supergravity induced by the presence of fluxes. We shall restrict our analysis, here, only to IIB orientifolds with some (three–form) fluxes turned on, while we shall defer the study of more general fluxes and of the gauge structure of other models elsewhere.

Differently to what happened in the well–studied $T_6/Z_2$ orientifolds, non–abelian gauged supergravities (for the bulk sector) now emerge, due to the presence of gauge fields originating from the ten–dimensional metric, and of axionic scalars associated to the NS–NS two–form $B$.

4.1. The $T_4 \times T_2$ IIB orientifold model

In this model, the allowed three–form fluxes are $H^\lambda_{ij} = \{H_{aij}, F_{aij}\}$, and are in correspondence with the representation $(4, 6)_2$ of $SO(2, 2) \times GL(4, \mathbb{R})$. The grading simply counts the number of indices along the internal $T_4$ and, more specifically, is associated to the subgroup $O(1, 1)_1 \subset GL(4, \mathbb{R})$. As mentioned in the introduction, inspection of the dimensionally reduced three–form kinetic term indicates for the four–dimensional theory a gauge group $\mathcal{G}_g$ with connection $\Omega_g = X_i G^i_\mu + X_\lambda A_\mu^\lambda$ and the following structure:

\[ [X_i, X_j] = H^\lambda_{ij} X_\lambda. \] (103)

We may identify the gauge generators with isometries as follows:

\[ X_i = -H^\lambda_{ij} T_\lambda^j, \]
\[ X^\lambda = \frac{1}{2} H^{\lambda'}_{ij} T^{ij}. \] (104)

Using relations (31) and the property

\[ H^\lambda_{k\ell} H_{ij} = \frac{1}{2} H^{\lambda'}_{ij} H_{k\ell} - \frac{1}{4} H \epsilon_{ijkl}, \] (105)

where $H = H^\lambda_{ij} H^{\lambda'}_{ij}$, one can show that the generators defined in (104) fulfil the following algebraic relations

\[ [X_i, X_j] = H^\lambda_{ij} X_\lambda - \frac{1}{4} H T_{ij}. \] (106)
which coincide with (103) only if \( H = 0 \) which amounts to the condition that \( \int_{T^6} F_{(3)} \wedge H^{(3)} = 0 \) (this condition is consistent with a constraint found in [18] on the embedding matrix of a new gauge group in the \( \mathcal{N} = 8 \) theory, which seems to yield an \( \mathcal{N} = 8 \) “lifting” of the type IIB orientifold models \( T_{p-3} \times T_{3-p} \) discussed here). Under this condition the gauge group is indeed contained in the isometry group of the scalar manifold. Moreover it can be verified that under the duality action of the gauge generators defined in (104) the vector fields transform in the co–adjoint of the gauge group \( G \) and thus provide a consistent definition for the gauge connection \( \Omega_g \). The variation of the gauge potentials under an infinitesimal transformation with parameters \( \xi^\lambda, \xi^i \) reads

\[
\delta A_\mu^\lambda = \xi^i H^\lambda_{ij} G_{\mu}^j + \partial_\mu \xi^\lambda, \\
\delta G_{\mu}^i = \partial_\mu \xi^i, \\
\delta C_{ijk\mu} = 0,
\]

(107)

and is compatible with the following non–abelian field strengths

\[
F_{\mu\nu}^\lambda = \partial_\mu A_\nu^\lambda - \partial_\nu A_\mu^\lambda - H_{ij}^\lambda G_{\mu}^i G_{\nu}^j, \\
\mathcal{F}_{\mu\nu}^i = \partial_\mu G_{\nu}^i - \partial_\nu G_{\mu}^i, \\
F_{\mu\nu}^i = \epsilon^{ijkl} (\partial_\mu C_{jkl\nu} - \partial_\nu C_{jkl\mu}).
\]

(108)

The \( C_{ij} \) and \( \Phi_i^\lambda \) scalars are also charged and, up to rotations, subject to shifts

\[
\delta C_{ij} = \frac{1}{2} H^\lambda_{ij} \xi^\lambda - \frac{1}{2} \xi^k H^\lambda_{k[i} \Phi_{j]\lambda}, \\
\delta \Phi_i^\lambda = H^\lambda_{ij} \xi^j,
\]

(109)

and their kinetic terms are modified accordingly by covariantisations

\[
D_\mu C_{ij} = \partial_\mu C_{ij} - \frac{1}{2} H_{ij\lambda} A_\mu^\lambda + \frac{1}{2} G_{\mu}^k H_{k[i}^\lambda \Phi_{j]\lambda}, \\
D_\mu \Phi_i^\lambda = \partial_\mu \Phi_i^\lambda - H_{ij}^\lambda G_{\mu}^j.
\]

(110)

**Chern–Simons terms.** The gauge group consists of Peccei–Quinn transformations that shift the real part of the vector kinetic matrix \( \mathcal{N} \) (the generalised theta angle). In [19], [50], it was shown that such a local transformation is a symmetry of the Lagrangian provided suitable generalised Chern–Simons terms are introduced.

In the case at hand, the new contribution to the Lagrangian is

\[
\mathcal{L}_{c.s.} \propto \epsilon^{\mu\nu\rho\sigma} \left( H_{\lambda ij'} A_{\mu}^\lambda G_{\nu}^i \partial_\rho C_{j'}^\sigma + \frac{1}{8} H_{\lambda ij'} H_{kl}^\lambda G_{\mu}^i G_{\nu}^j G_{\rho}^k G_{\sigma}^l \right),
\]

(111)

corresponding to the non–vanishing entries

\[
C_{\lambda, ij'} = - H_{\lambda ij'} \quad \text{and} \quad C_{i, \lambda j'} = H_{\lambda ij'}
\]

(112)
where, in general, the coefficients $C_{\Gamma,\Lambda\Sigma}$ define the moduli–independent gauge variation of the real part of the kinetic matrix $\mathcal{N}$

$$\delta_\xi \text{Re} \mathcal{N}_{\Lambda\Sigma} = \xi^T C_{\Gamma,\Lambda\Sigma}.$$  \hfill (113)

4.2. Type $T_2 \times T_4$ IIB orientifold model.

Let us consider the $T_2 \times T_4$ model in presence of the fluxes $H_{ija} = \epsilon_{ij} H_a$ and $F_{iab}$. These fluxes appear as structure constants

$$[X_i, X_j] = \epsilon_{ij} H_a X^a,$$

$$[X_i, X^a] = F_{iab} X_b,$$  \hfill (114)

of the gauge algebra $\mathscr{G}_g \equiv \{X_i, X^a, X_a\}$ with connection $\Omega^a_\mu = G^i_\mu X_i + B_{a\mu} X^a + C_{\mu}^a X_a$, all other commutators vanishing.

The identification

$$X'_i = -F_{iab} T_{ab} + H_a T_i,$$

$$X'^a = F_{iab} T_i^b, \quad X'_a = -H_a T,$$  \hfill (115)

of the gauge generators with the isometries of the solvable algebra $N_5$, reproduces only a contracted version of the algebra (114) in which three of the central charges $X_a$ vanish and we are left with $X'_a = -H_a T$. If we denote by $\{X^\perp\} = \{X_a\}/\{X'_a\}$ these three central generators, we see that the subgroup $\mathscr{G}'_g = \{X'_i, X'^a, X'_a\}$ of the isometry group which is gauged coincides with the quotient:

$$\mathscr{G}'_g \equiv \mathscr{G}_g/\{X^\perp\},$$  \hfill (116)

that amounts to imposing the vanishing of the central terms on all fields.

On the other hand, transformations generated by the operators in (115) induce isometry transformations with parameters:

$$\xi_{ia} = -\xi_i H_a,$$

$$\xi^{ab} = -\xi^i F_{i}^{ab},$$

$$\xi^{a} = \xi_b F_{b}^{a},$$

$$\xi = -H_a \xi^a,$$  \hfill (117)

where $\xi_i = \epsilon_{ij} \xi^j$. Using eqs. (122) and (117), one can then verify that the vectors $G^i_\mu$, $B_{a\mu}$ and $C_{\mu}^a$ transform in the co–adjoint representation of $\mathscr{G}_g$ under the duality action
generated by \( \{X_i, X^a, X_a\} \), so that the above definition of the gauge connection \( \Omega^a_\mu \) is consistent:

\[
\begin{align*}
\delta B_{a\mu} &= \xi^i G^j_\mu \epsilon_{ij} H_a + \partial_\mu \xi_a = -\xi_i G^a_\mu H_a + \partial_\mu \xi_a, \\
\delta C^a_\mu &= \xi^i B_{b\mu} F^{ba}_i - G^a_\mu \xi_b F^{ba}_i + \partial_\mu \xi^a, \\
\delta G^a_\mu &= \partial_\mu \xi^a.
\end{align*}
\]

(118)

Notice that the action of the central charges \( X_a \) amounts just to a gauge transformation on \( C^a_\mu \). These ten vectors can therefore be used to gauge the group \( G_g \), and the non–abelian field strengths read

\[
\begin{align*}
\mathcal{H}_{a\mu\nu} &= \partial_\mu B_{a\nu} - \partial_\nu B_{a\mu} - \epsilon_{ij} H_a G^i_\mu G^j_\nu, \\
F^{a\mu}_\nu &= \partial_\mu C^a_\nu - \partial_\nu C^a_\mu + F^a_\mu B_{b\nu} - F^a_\nu B_{b\mu}, \\
\mathcal{F}^i_{\mu\nu} &= \partial_\mu G^i_\nu - \partial_\nu G^i_\mu.
\end{align*}
\]

(119)

Since \( G_g \) is not part of the global symmetries of the Lagrangian, we should restrict ourselves to the quotient \( G_g \), i.e. we demand that central charges \( \{T, X_a\} \) vanish on all physical fields. The gauge transformations of the scalar fields

\[
\begin{align*}
\delta c &= -H_a \xi^a + \xi_a F^{ab}_i B_{jb} \epsilon^{ij}, \\
\delta C^a_i &= \xi^b F^{ba}_i + \xi^j F^{ab}_j B_{bi}, \\
\delta B_{ia} &= -\xi_i H_a, \\
\delta C_{ab} &= -\xi^i F_{iab},
\end{align*}
\]

(120)

are then compatible with the covariant derivatives

\[
\begin{align*}
D_\mu c &= \partial_\mu c + H_a C^a_\mu - B_{a\mu} F^{ab}_i B_{jb} \epsilon^{ij}, \\
D_\mu C^a_i &= \partial_\mu C^a_i + B_{b\mu} F^{ba}_i - G^b_\mu F^{ab}_j B_{bi}, \\
D_\mu B_{ia} &= \partial_\mu B_{ia} + G^a_\mu H_a, \\
D_\mu C_{ab} &= \partial_\mu C_{ab} + G^a_\mu F_{iab}.
\end{align*}
\]

(121)

**Chern–Simons terms.** Also in this case local Peccei–Quinn transformations demand the inclusion in the Lagrangian of the Chern–Simons terms

\[
\mathcal{L}_{c.s.} = \epsilon^{\mu\nu\rho\sigma} \left( H_a G^i_\mu C_{iv} \partial_\rho C^a_\sigma - H_a C^a_\mu C_{iv} \partial_\rho G^i_\sigma - \epsilon^{ij} F^{ab}_j B_{b\mu} C_{iv} G^k_\rho G_{k\sigma} + \frac{1}{8} H_a F^{ab}_k G^i_\mu C_{iv} G^k_\rho B_{b\sigma} - \frac{1}{8} \epsilon^{ij} H_a F^{ab}_j B_{b\mu} C_{iv} G^k_\rho G_{k\sigma} \right),
\]

(122)

corresponding to the non–vanishing components

\[
\begin{align*}
C_{i, a} &= \delta^i_j H_a, \\
C_{a, i} &= -\delta^i_j H_a, \\
C^{a, ib} &= -\epsilon^{ij} F^{ab}_j,
\end{align*}
\]

(123)

of the \( C_{\Gamma, \Lambda \Sigma} \) coefficients.
5. Conclusions and outlooks

In the present paper, we have investigated the symmetries and the structure of several $T_6$ orientifolds which, in absence of fluxes, have $\mathcal{N}=4$ supersymmetries in four dimensions. We have not addressed here the question of vacua with some residual supersymmetry, that will be the subject of future investigations. All these models lead to different low–energy supergravity descriptions. When fluxes are turned on, the deformed Lagrangian is described by a gauged $\mathcal{N}=4$ supergravity and fermionic mass–terms and a scalar potential are developed.

The low–energy Lagrangians underlying these orientifolds are different versions of gauged $\mathcal{N}=4$ supergravity with six bulk vector multiplets and additional Yang–Mills multiplets living on the brane world–volume. The gaugings are based on quotients (with respect to some central charges) of nilpotent subalgebras of $so(6,6)$. These nilpotent subalgebras are basically generated by the axion symmetries associated to R–R scalars and to NS–NS scalars originating from the two–form $B$–field.

Along similar lines, one can also consider new examples of orientifolds with $\mathcal{N}=2,1$ four–dimensional supersymmetries, with and/or without fluxes.

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Appendix. Symplectic realisation of the solvable generators

In this appendix, we give the coset representatives of our models in the symplectic basis of vector fields. This is needed in order to compute the kinetic matrix $\mathcal{N}_{\Lambda \Sigma}$, which is a complex symmetric matrix in the space of vectors in the theory. Its imaginary and real parts describe the terms

$$\text{Im} \mathcal{N}_{\Lambda \Sigma} F^A_{\mu \nu} F^{\Sigma \mu \nu} + \frac{1}{2} \text{Re} \mathcal{N}_{\Lambda \Sigma} \epsilon^{\mu \nu \rho \sigma} F^A_{\mu \nu} F^{\Sigma \rho \sigma}.$$

(124)
**Model** $T_4 \times T_2$  The $Sp(24, \mathbb{R})$ representation of the solvable generators in model 1 in the basis (15) is:

\[
T = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 \\
0 & 1 & 0 & 0 & 0 & 0 \\
\eta & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & -1 \\
0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix},
\]

\[
T = \begin{pmatrix}
M^T & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & -M & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix},
\]

\[
T^i_\lambda = \begin{pmatrix}
0 & -(t^i_\lambda)^T & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 \\
0 & 0 & -t^i_\lambda n & 0 & 0 \\
-t^i_\lambda n & 0 & 0 & t^i_\lambda & 0 \\
0 & 0 & 0 & 0 & 0
\end{pmatrix},
\]

\[
T^{ij} = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0
\end{pmatrix},
\]

where each block is a $4 \times 4$ matrix, $\mathbb{I}$ denotes the identity matrix, $\eta \equiv \eta_{\lambda \lambda'}$ and

\[
(t^i_\lambda)^X = \delta^i_j \delta^X_{\lambda'}, \quad (t^{ij})_{kl} = \delta^i_k \delta^j_l - \delta^i_l \delta^j_k.
\]

The coset representative is

\[
L = \exp(C_{ij} T^{ij}) \exp(\Phi^i T^i_\lambda) \exp(cT) \mathbb{E} = \begin{pmatrix} A & 0 \\ C & D \end{pmatrix},
\]

where $\mathbb{E}$ parametrises the manifold

\[
\mathbb{E} \in O(1, 1)_0 \times \frac{SO(2, 2)}{SO(2) \times SO(2)} \times \frac{GL(4, \mathbb{R})}{SO(4)}.
\]
and can be written in the following general form:

\[
E = \begin{pmatrix}
e^{-\varphi} E_{(t)}^\lambda \hat{\sigma} & 0 & 0 & 0 & 0 & 0 \\
0 & e^{-\varphi} E & 0 & 0 & 0 & 0 \\
0 & 0 & e^\varphi E & 0 & 0 & 0 \\
0 & 0 & 0 & e^\varphi \eta E_{(t)} \eta & 0 & 0 \\
0 & 0 & 0 & 0 & e^\varphi E^{-1} & 0 \\
0 & 0 & 0 & 0 & 0 & e^{-\varphi} E^{-1} \\
\end{pmatrix},
\]  
(129)

with

\[
E_{(t)}^\lambda \hat{\sigma} \in \frac{SO(2,2)}{SO(2) \times SO(2)},
\]
\[
E^i_j \in \frac{GL(4,\mathbb{R})}{SO(4)},
\]
\[
e^\varphi \Phi^H \in O(1,1)_0,
\]  
(130)

the hatted indices being the rigid ones transforming under the isotropy group. The blocks is \( L \) read

\[
A = \begin{pmatrix}
e^{-\varphi} E_{(t)}^\lambda \hat{\sigma} & -e^{-\varphi} \Phi^\lambda_i E^i_j & 0 \\
0 & e^{-\varphi} E^i_j & 0 \\
0 & ce^{-\varphi} E^i_j & e^\varphi E^j_i \\
\end{pmatrix},
\]
\[
C = \begin{pmatrix}
ce^{-\varphi} E_{(t)}^\lambda \hat{\sigma} & -ce^{-\varphi} \Phi_{ij} E^i_j & -e^\varphi \Phi_{ij} E^j_i \\
ce^{-\varphi} \Phi_{ij} E_{(t)}^\delta & -ce^{-\varphi} 2 \hat{C}_{ij} E^i_j & -e^\varphi 2 \hat{C}_{ij} E^j_i \\
-e^\varphi \Phi_{ij} E_{(t)}^\delta & e^{-\varphi} 2 \hat{C}_{ij} E^i_j & 0 \\
\end{pmatrix},
\]
\[
D = \begin{pmatrix}
e^{\varphi} E_{(t)}^\lambda \hat{\sigma} & 0 & 0 \\
0 & e^{\varphi} E^{-1} \hat{\sigma} & -ce^{-\varphi} E^{-1} \hat{\sigma} \\
0 & 0 & e^{-\varphi} E^{-1} \hat{\sigma} \\
\end{pmatrix},
\]
\[
\tilde{C}_{ij} = C_{ij} + \frac{1}{4} \Phi^\lambda_i \Phi^\lambda_j.
\]  
(131)

In the sequel we shall need also the expression of \( A^{-1} \):

\[
A^{-1} = \begin{pmatrix}
e^{\varphi} E_{(t)}^\lambda \hat{\sigma} & e^\varphi E_{(t)}^\lambda \hat{\sigma} \Phi^\lambda_i & 0 \\
0 & e^{\varphi} E^{-1} \hat{\sigma} & 0 \\
0 & -e^{-\varphi} ce^{-\varphi} E^{-1} \hat{\sigma} & e^{-\varphi} E^{-1} \hat{\sigma} \\
\end{pmatrix}.
\]  
(132)

In terms of the matrices \( h, f \)

\[
f = \frac{1}{\sqrt{2}} A, \quad h = \frac{1}{\sqrt{2}} (C - iD),
\]  
(133)

the kinetic matrix is expressed as (see 51 and references therein)

\[
\mathcal{N} = h f^{-1} = \begin{pmatrix}
N_{\lambda\nu} & N_{\lambda i} & N_{\lambda i}' \\
* & N_{ij} & N_{ij}' \\
* & * & N_{ij}''
\end{pmatrix},
\]  
(134)

and is characterised by the following entries:

\[
N_{\lambda\nu} = -i e^{2\varphi} E_{(t)}^\lambda \hat{\sigma} E_{(t)}^\lambda \hat{\sigma} + c \eta_{\lambda\nu},
\]
\[
N_{\lambda i} = -i e^{2\varphi} E_{(t)}^\lambda \hat{\sigma} E_{(t)}^\lambda \hat{\sigma} \Phi^X_i + c \Phi_{\lambda i},
\]
\[
N_{\lambda i}' = -\Phi^X_i,
\]
\[
N_{ij} = -i \left( (e^{2\varphi} + e^{-2\varphi} c^2) E^{-1} i E^{-1} j + e^{2\varphi} \Phi^X_i E_{(t)}^\lambda \hat{\sigma} E_{(t)}^\lambda \hat{\sigma} \Phi^X_j \right) + c \Phi^X_i \Phi^X_j,
\]
\[
N_{ij}' = i c e^{-2\varphi} E^{-1} i E^{-1} j - 2 \tilde{C}_{ij},
\]
\[
N_{ij}'' = -i e^{-2\varphi} E^{-1} i E^{-1} j.
\]  
(135)
Model $T_2 \times T_4$ The $Sp(24, \mathbb{R})$ representation of the solvable generators in model 2 in the basis (61) is:

$$T' = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\epsilon_{ij} & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & \mathbb{1} & 0 & 0 & 0 & 0 \\
0 & 0 & \mathbb{1} & 0 & 0 & 0 & 0 & 0
\end{pmatrix},$$

$$T_{ab} = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & \delta^{cd} & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & \delta^{cd} & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix},$$

$$T^{ia} = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\delta^{a}_b \delta^{i}_j & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & -\epsilon^{ij} \delta^{a}_b & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & -\delta^{a}_b \epsilon^{ij} & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix},$$

$$T^i_a = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & \epsilon^{ij} \delta^{b}_a & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & -\delta^{b}_a \delta^{i}_j & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & \delta^{b}_a \epsilon^{ij} & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix}.$$

37
The coset representative has the form:

\[
L = e^c T' e^c T e^{B_{ta} T^a} e^{C^a T_a} e^{C_{ab} T^{ab}} \mathbb{E} = \begin{pmatrix} A & 0 \\ C & D \end{pmatrix},
\]

(137)

where this time the matrix \( \mathbb{E} \) describes the submanifold:

\[
\mathbb{E} = O(1, 1)_0 \times \frac{GL(2, \mathbb{R})}{SO(2)} \times \frac{GL(4, \mathbb{R})}{SO(4)},
\]

(138)

and has the following form:

\[
\mathbb{E} = \begin{pmatrix}
\begin{array}{cccccccc}
 e^{-\varphi} E_{2i} & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
 0 & e^{\varphi} E^{-1}_{2j} & 0 & 0 & 0 & 0 & 0 & 0 \\
 0 & 0 & e^{-\varphi} E^{-1}_{4i} & 0 & 0 & 0 & 0 & 0 \\
 0 & 0 & 0 & e^{\varphi} E_{4b} & 0 & 0 & 0 & 0 \\
 0 & 0 & 0 & 0 & e^{-\varphi} E_{2i} & 0 & 0 & 0 \\
 0 & 0 & 0 & 0 & 0 & e^{\varphi} E_{4b} & 0 & 0 \\
 0 & 0 & 0 & 0 & 0 & 0 & e^{\varphi} E^{-1}_{4i} & 0 \\
 0 & 0 & 0 & 0 & 0 & 0 & 0 & e^{-\varphi} E_{4b} \\
\end{array}
\end{pmatrix},
\]

(139)

The blocks \( A, C, D \) of \( L \) can be conveniently described in terms of the following matrices:

\[
(B)_{ta} = B_{ta}, \quad (C)^a_i = C^a_i, \quad (\mathcal{E})^{ab} = C^{ab}.
\]

(140)
it is also useful to compute $A^{-1}$:

$$
A^{-1} = \begin{pmatrix}
    e^{\varphi} E_1^{-1} & 0 & 0 & 0 \\
    -e^{-\varphi} c E_2 & e^{\varphi} E_2 & 0 & 0 \\
    -e^{\varphi} E_4 B^i & 0 & e^{\varphi} E_4 & 0 \\
    e^{\varphi} E_4^{-1} (B^i + C^i) & 0 & -e^{\varphi} E_4^{-1} c^i & e^{\varphi} E_4^{-1} \\
\end{pmatrix}.
$$

(142)

We then compute the kinetic matrix $\mathcal{M}$ whose independent components are:

$$
\mathcal{M} = \begin{pmatrix}
    N_{ij} & N_{ij}^a & N_{ia} & N_{ia}^a \\
    * & N_{ij} & N_{ia} & N_{ia}^a \\
    * & * & N_{ab} & N_{ab}^a \\
    * & * & * & N_{ab} \\
\end{pmatrix},
$$

(143)

where

$$
\begin{align*}
N_{ij} &= -i \left[ E_{2}^{-1;j} E_{2}^{-1;j} (e^{2\varphi} + e^{-2\varphi} c^i) + e^{2\varphi} B_{ia} E_{4}^{a \hat{a}} E_{4}^{\hat{a} \hat{b}} B_{j\hat{b}} +
    e^{2\varphi} (-B_{ia} C^{ca} + C^{ca}_{i}) E_{4}^{-1} \hat{a} E_{4}^{-1} \hat{b} (C^{bd} B_{j\hat{d}} + C_{j}^{\hat{d} \hat{b}}) - 2 B_{a(i} C^{ca}_{j)} c^i, \\
N_{ij}^i &= -i e^{-2\varphi} c^i \epsilon_{ik} E_{2}^{j \hat{k}} E_{2}^{j \hat{k}} + c \delta_{jk} - B_{ia} C^{ca}_{k} e^{k} \\
N_{ia} &= -i e^{2\varphi} B_{ib} E_{4}^{b \hat{a}} E_{4}^{a \hat{a}} + (-B_{ib} C^{cb} + C^{cb}_{i}) E_{4}^{-1} \hat{c} E_{4}^{-1} \hat{a} C^{da} + c^i C^{ca}_{i}, \\
N_{ia}^a &= -i e^{2\varphi} B_{ib} E_{4}^{b \hat{a}} E_{4}^{a \hat{a}} + (-B_{ib} C^{cb} + C^{cb}_{i}) E_{4}^{-1} \hat{c} E_{4}^{-1} \hat{a} C^{da} + c^i C^{ca}_{i}, \\
N_{ij}^i &= -i e^{-2\varphi} E_{2}^{i \hat{k}} E_{2}^{j \hat{k}}, \\
N_{ia}^i &= -i e^{2\varphi} B_{ja}, \\
N_{ia} &= -i e^{2\varphi} B_{ja}, \\
N_{ab} &= -i e^{2\varphi} \left( -C^{ad} E_{4}^{-1} \hat{d} E_{4}^{-1} \hat{c} C^{cb} + E_{4}^{a b} E_{4}^{b \hat{a}} \right), \\
N_{ab}^a &= -i e^{2\varphi} E_{4}^{-1} \hat{a} E_{4}^{-1} \hat{d} + c^i \delta_{ab}^i, \\
N_{ab} &= -i e^{2\varphi} E_{4}^{-1} \hat{a} E_{4}^{-1} \hat{d}.
\end{align*}
$$

(144)

**Model $T_1 \times T_5$** The $Sp(24, \mathbb{R})$ representation of the $N_4$ generators is the following:

$$
T^a = \begin{pmatrix}
    0 & 0 & 0 & 0 & 0 & 0 & 0 \\
    0 & 0 & 0 & 0 & 0 & 0 & 0 \\
    0 & 0 & 0 & 0 & 0 & 0 & 0 \\
    0 & 0 & 0 & 0 & 0 & -\delta_b^a & 0 \\
    0 & 0 & 0 & 0 & 0 & 0 & 0 \\
    0 & 0 & 0 & 0 & 0 & 0 & 0 \\
    0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\end{pmatrix},
$$

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We have chosen the coset representative to have the form given in eq. (78). We may choose for the matrix $E$ the following matrix form:

$$E = \begin{pmatrix}
    e^{\varphi} E & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
    e^{-\varphi} E & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
    0 & e^{\varphi} E^{-1} a_b & 0 & 0 & 0 & 0 & 0 & 0 \\
    0 & 0 & e^{-\varphi} E^{-1} a_b & 0 & 0 & 0 & 0 & 0 \\
    0 & 0 & 0 & e^{-\varphi} / E & 0 & 0 & 0 & 0 \\
    0 & 0 & 0 & 0 & e^{\varphi} / E & 0 & 0 & 0 \\
    0 & 0 & 0 & 0 & 0 & e^{-\varphi} E^a_b & 0 & 0 \\
    0 & 0 & 0 & 0 & 0 & 0 & e^{\varphi} E^a_b & 0 \\
\end{pmatrix},$$

where:

$$E^a_b, \quad E \in O(1,1) \times O(1,1)_2 \times SL(5, \mathbb{R}) / SO(5),$$

$$e^{H \varphi} \in O(1,1)_0.$$
\[ C = \begin{pmatrix}
0 & B_a C^a E e^{-\varphi} & 0 & e^{-\varphi} (B_b C^{ba} + C^a) E^{-1}_{a \ d} \\
-B_a C^a E e^\varphi & -c B_a C^a E e^{-\varphi} & -e^\varphi (B_b C^{ba} + C^a) E^{-1}_{a \ d} & -e^{-\varphi} e (B_b C^{ba} + C^a) E^{-1}_{a \ d} \\
0 & -e^\varphi C^a E & 0 & e^{-\varphi} C^{ab} E^{-1}_{b \ d} \\
C^a E e^\varphi & C^a E e^{-\varphi} & e^\varphi C^{ab} E^{-1}_{b \ d} & e^{-\varphi} C^{ab} E^{-1}_{b \ d}
\end{pmatrix},
\]

\[ D = \begin{pmatrix}
e^{-\varphi}/E & 0 & -B_a e^{-\varphi} E^a_d & 0 \\
e^{-\varphi}/E & e^\varphi/E & c B_a e^{-\varphi} E^a_d & -B_a e^\varphi E^a_d \\
0 & 0 & e^{-\varphi} E^a_d & 0 \\
0 & 0 & -c e^{-\varphi} E^a_d & e^\varphi E^a_d
\end{pmatrix},
\]

\[ A^{-1} = \begin{pmatrix}
e^{-\varphi} E^a_d B_a & e^{-\varphi} E^a_d B_a & e^{-\varphi} E^a_d & -c e^{-\varphi} E^a_d \\
e^{-\varphi} E^a_d B_a & e^{-\varphi} E^a_d B_a & e^{-\varphi} E^a_d & e^\varphi E^a_d \\
0 & 0 & e^\varphi E^a_d & 0 \\
0 & 0 & e^\varphi E^a_d & 0
\end{pmatrix}, \quad (148)
\]

from equations (133) and (134) we compute the matrix \( \mathcal{N} \):

\[ \mathcal{N} = \begin{pmatrix}
N & N_1 & N^{(1a)} & N^a \\
* & N_{11} & N^{(1a)}_1 & N^{(1a)}_1 \\
* & * & N^{(1a)}(1b) & N^{(1a)}b \\
* & * & * & N^{ab}
\end{pmatrix}, \quad (149)
\]

whose entries are

\[ N = -i e^{-2\varphi} \left( B_a B_b E^a_{\ a} E^b_{\ a} + \frac{1}{E^2} \right), \]

\[ N_1 = i c e^{-2\varphi} \left( B_a B_b E^a_{\ a} E^b_{\ a} + \frac{1}{E^2} \right), \]

\[ N^{(1a)} = i e^{-2\varphi} B_b E^a_{\ a} E^b_{\ a}, \]

\[ N^a = C^a + B_b C^{ba} - i c e^{-2\varphi} B_b E^a_{\ a} E^b_{\ a}, \]

\[ N_{11} = -i (e^{-2\varphi} e^2 + e^{2\varphi}) \left( B_a B_b E^a_{\ a} E^b_{\ a} + \frac{1}{E^2} \right), \]

\[ N^{(1a)}_1 = -(B_b C^{ba} + C^a) - i c e^{-2\varphi} B_b E^a_{\ a} E^b_{\ a}, \]

\[ N^{(1a)}_1 = i (e^{-2\varphi} e^2 + e^{2\varphi}) B_b E^a_{\ a} E^b_{\ a}, \]

\[ N^{(1a)}(1b) = -i e^{-2\varphi} E^a_{\ a} E^b_{\ a}, \]

\[ N^{(1a)b} = -C^{ab} + i e^{-2\varphi} c E^a_{\ a} E^b_{\ a}, \]

\[ N^{ab} = -i (e^{-2\varphi} e^2 + e^{2\varphi}) E^a_{\ a} E^b_{\ a}, \quad (150)\]

where the asterisks denote the symmetric entries.

**Model \( T_3 \times T_3 \)** The \( Sp(24, \mathbb{R}) \) representation of the \( N_6 \) generators is the following:

\[ T_{ab} = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & \delta_{ab}^{cd} & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix}, \]
\[ T^{ia} = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\delta^a_b \delta^i_j & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\delta^a_b & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\end{pmatrix}, \]

\[ T^i_a = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\delta^b_a \delta^i_j & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\delta^b_a & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\end{pmatrix}, \]

\[ T^{ij} = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
-2 \delta^{ij}_{kl} & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
2 \delta^{ij}_{kl} & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\end{pmatrix}, \]

\[ T = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
1 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & -1 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 1 & 0 & 0 & 0 & 0 \\
0 & 0 & 1 & 0 & 0 & 0 & 0 & 0 \\
\end{pmatrix}. \]
We have chosen the coset representative to have the form given in eq. (88). We may choose for the matrix $E$ the following matrix form:

$$
E = \begin{pmatrix}
  e^{-\varphi} E_{1i}^j & 0 & 0 & 0 & 0 & 0 \\
  0 & e^{\varphi} E_{1j}^i & 0 & 0 & 0 & 0 \\
  0 & 0 & e^{-\varphi} E_{2-1}^i & 0 & 0 & 0 \\
  0 & 0 & 0 & e^{\varphi} E_{2b}^i & 0 & 0 \\
  0 & 0 & 0 & 0 & e^{-\varphi} E_{1-1}^i & 0 \\
  0 & 0 & 0 & 0 & 0 & e^{\varphi} E_{2-1}^i \\
  0 & 0 & 0 & 0 & 0 & 0 \\
\end{pmatrix},
$$

(152)

where

$$
E_{1i}^j \in \left(\frac{GL(3, \mathbb{R})}{SO(3)}\right)_1,
E_{2b}^i \in \left(\frac{GL(3, \mathbb{R})}{SO(3)}\right)_2,
\varphi H \in SO(1, 1)_0.
$$

(153)

The blocks $\{A, C, D\}$ of $L$ and $A^{-1}$ have the following form:

$$
A = \begin{pmatrix}
  e^{-\varphi} E_{1i}^j & 0 & 0 & 0 \\
  c e^{-\varphi} E_{1i}^j & e^{\varphi} E_{1j}^i & 0 & 0 \\
  B_{1a} e^{-\varphi} E_{1i}^j & 0 & e^{-\varphi} E_{2-1}^a & 0 \\
  C_{1a} e^{-\varphi} E_{1i}^j & 0 & 0 & e^{\varphi} E_{2a}^i \\
\end{pmatrix},
$$

$$
C = \begin{pmatrix}
  -2 c e^{-\varphi} C_{1a} E_{1k} & -2 e^{\varphi} \tilde{C}_{ij} E_{1k} & -c e^{-\varphi} (C_{i}^{a} + B_{ib} C_{ba}) E_{2-1}^a & -e^{-\varphi} c B_{ib} E_{2-1}^b \\
  2 e^{-\varphi} \tilde{C}_{ij} E_{1k} & 0 & e^{-\varphi} (C_{i}^{a} + B_{ib} C_{ba}) E_{2-1}^a & -e^{-\varphi} c B_{ib} E_{2a}^i \\
  -c e^{-\varphi} C_{i}^{a} E_{1k} & e^{\varphi} C_{1a} E_{1k} & 0 & e^{\varphi} E_{2a}^i \\
  -c e^{-\varphi} C_{i}^{a} E_{1k} & e^{\varphi} C_{1a} E_{1k} & 0 & 0 \\
\end{pmatrix},
$$

$$
D = \begin{pmatrix}
  e^{\varphi} E_{1-1}^j & -c e^{-\varphi} E_{1-1}^j & -e^{\varphi} B_{ia} E_{2-1}^j & -e^{\varphi} (C_{i}^{a} + B_{ib} C_{ba}) E_{2-1}^a \\
  0 & e^{-\varphi} E_{1-1}^j & 0 & 0 \\
  0 & 0 & e^{\varphi} E_{2-1}^j & e^{\varphi} E_{2-1}^a \\
  0 & 0 & 0 & e^{\varphi} E_{2-1}^a \\
\end{pmatrix},
$$

$$
A^{-1} = \begin{pmatrix}
  e^{\varphi} E_{1-1}^j & 0 & 0 & 0 \\
  -c e^{\varphi} E_{1-1}^j & e^{-\varphi} E_{1-1}^j & 0 & 0 \\
  -e^{\varphi} E_{2a}^a B_{ia} & 0 & e^{\varphi} E_{2a}^a & 0 \\
  -e^{-\varphi} E_{2-1}^a B_{ia} & 0 & 0 & e^{-\varphi} E_{2-1}^a \\
\end{pmatrix},
$$

$$
\tilde{C}_{ij} = C_{ij} + \frac{1}{2} C_{ai} B_{ja}.
$$

(154)

The vector kinetic matrix can now be calculated and has the following form:

$$
\mathcal{N} = \begin{pmatrix}
  N_{ij} & N_{ij}' & N_{ia} & N_{ia}' \\
  * & N_{ij}'' & N_{ia}' & N_{ia}' \\
  * & * & N_{ab} & N_{ab}' \\
  * & * & * & N_{ab} \\
\end{pmatrix},
$$

(155)

whose entries are

$$
N_{ij} = 2 c B_{ai} C_{aj} - i \left[ (c^{2} e^{-2\varphi} + e^{2\varphi}) E_{1-1}^i E_{1-1}^j +
 e^{2\varphi} (C_{i}^{a} + B_{ib} C_{ba}) (C_{j}^{c} + B_{jc} C_{bc}) E_{2-1}^c E_{2-1}^c + e^{2\varphi} B_{ia} B_{jb} E_{2c}^a E_{2c}^b \right],
$$

43
\[
N_{ij}^\prime = -2 \bar{c}_{ij} + i c e^{-2\varphi} E_1^{-1} i E_1^{-1} j, \\
N_i^a = -c C_i^a + i e^{2\varphi} \left[ \left( C_i^{cb} + B_{ic} C^{cb} \right) C^{ad} E_2^{-1} b \hat{E}_2^{-1} \hat{c} + B_{ab} E_2^a E_2^b \right], \\
N_{ia} = -c B_{ia} + i e^{2\varphi} \left( C_i^{cb} + B_{ic} C^{cb} \right) E_2^{-1} b \hat{E}_2^{-1} \hat{a}, \\
N_{ij}^{\prime a} = -i e^{-2\varphi} E_1^{-1} i E_1^{-1} j, \\
N_i^{\prime a} = C_i^a, \\
N_{ia}^{\prime} = B_{ia}, \\
N^{ab} = -i e^{2\varphi} \left[ E_2^a \hat{E}_2^b + C^{ac} C^{bd} E_2^{-1} c \hat{E}_2^{-1} \hat{d} \right], \\
N_{ab}^a = c - i e^{2\varphi} C^{ac} E_2^{-1} c \hat{E}_2^{-1} \hat{b}, \\
N_{ab} = -i e^{2\varphi} E_2^{-1} a \hat{E}_2^{-1} \hat{b}. \\
\]

(156)

**Model** $T_5 \times S_1$  The $Sp(24, \mathbb{R})$ representation of the $N_8$ generators is the following:

\[
T^i = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
-\delta_i^j & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & -\delta_j^i & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix},
\]

\[
T^{\prime i} = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & -\delta_j^i & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix},
\]

\[
T^{ij} = \begin{pmatrix}
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0 \\
\delta_{kl}^i & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & 0
\end{pmatrix},
\]
We have chosen the coset representative to have the form given in eq. 98. We may choose for the matrix $E$ the following matrix form

$$E = \begin{pmatrix}
  e & 0 & 0 & 0 & 0 & 0 & 0 \\
  e^{-φ} E^{-1, j} & 0 & 0 & 0 & 0 & 0 & 0 \\
  0 & 0 & e^{-φ} E & 0 & 0 & 0 & 0 \\
  0 & 0 & 0 & e^{-φ} / E & 0 & 0 & 0 \\
  0 & 0 & 0 & 0 & e^{-φ} E^{-1, j} & 0 & 0 \\
  0 & 0 & 0 & 0 & 0 & e^{-φ} / E & 0 \\
  0 & 0 & 0 & 0 & 0 & 0 & e^{-φ} E
\end{pmatrix}, \quad (158)
$$

where:

$$E^i_j, E \in O(1, 1)_1 \times O(1, 1)_2 \times \frac{SL(5, \mathbb{R})}{SO(5)},$$

$$e^{φ} H \in SO(1, 1)_0. \quad (159)$$

The blocks $\{A, C, D\}$ of $L$ and $A^{-1}$ have the following form

$$A = e^{-φ} \begin{pmatrix}
  E^i_j & 0 & 0 & 0 \\
  (B_i C_j + C_{ij}) & E^i_j & E^{-1, i} & -B_i E & C_i / E \\
  -C_i E^i_j & 0 & E & 0 \\
  B_i E^i_j & 0 & 0 & 1 / E
\end{pmatrix},$$

$$C = e^{-φ} \begin{pmatrix}
  c B_i E^i_j & c E^{-1, j} & -c B_i E & c C_i / E \\
  c E^i_j & 0 & 0 & 0 \\
  c B_i E^i_j & 0 & 0 & c / E \\
  -c C_i E^i_j & 0 & c E & 0
\end{pmatrix},$$

$$D = e^{φ} \begin{pmatrix}
  E^{-1, i} & (B_i C_j + C_{ij}) & E^i_j & C_i / E & -B_i E \\
  0 & E^i_j & 0 & 0 \\
  0 & B_i E^i_j & 1 / E & 0 \\
  0 & -C_i E^i_j & 0 & E
\end{pmatrix},$$

$$A^{-1} = e^{φ} \begin{pmatrix}
  E^{i, j} & (C_j B_i - C_{ij}) & E^i_j & E^{i, j} B_j & -E^{i, j} C_j \\
  0 & E^i_j & 0 & 0 \\
  0 & 0 & 1 / E & 0 \\
  0 & -B_i E & 0 & E
\end{pmatrix}. \quad (159)$$

The vector kinetic matrix can now be calculated and has the following form

$$\mathcal{N} = \begin{pmatrix}
  N_{ij} & N_i^j & N_i & N_i' \\
  * & N^{ij} & N^i & N'^{ii} \\
  * & * & N & N' \\
  * & * & * & N''
\end{pmatrix}, \quad (160)$$
whose entries are

\[
N_{ij} = -ie^{2\varphi} \left[ E^{-1} E^{-1} + (B_i C_k + C_{ik}) (B_j C_n + C_{jn}) E^k E^l + \frac{1}{E^2} C_i C_j + B_i B_j E^2 \right],
\]

\[
N^i_j = c - ie^{2\varphi} \left[ (B_i C_k + C_{ik}) E^k E^j \right],
\]

\[
N_i = -ie^{2\varphi} \left[ (B_i C_k + C_{ik}) E^k B_j E^j + \frac{1}{E^2} C_i \right],
\]

\[
N'_i = ie^{2\varphi} \left[ (B_i C_k + C_{ik}) E^k E^j E^j + E^2 B_i \right],
\]

\[
N^{ij} = -ie^{2\varphi} E^i E^j,
\]

\[
N^i = -ie^{2\varphi} B_j E^i E^j,
\]

\[
N''^{ij} = ie^{2\varphi} C_j E^i E^j,
\]

\[
N = -ie^{2\varphi} \left[ B_i B_j E^i E^j \right],
\]

\[
N' = c + ie^{2\varphi} C_i B_j E^i E^j,
\]

\[
N'' = -ie^{2\varphi} \left[ C_i C_j E^i E^j + E^2 \right].
\]

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