Anomalies, Chern-Simons terms and the Standard Model

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Abstract. In orientifold vacua where usually there are several anomalous U(1)s, 4-d anomalies are cancelled via appropriate couplings to axions. However, If some linear combinations of U(1)s are non-anomalous, then there are additional Chern-Simons-like couplings necessary for the cancellation of the mixed abelian anomalies between anomalous and non-anomalous abelian factors. In the context of Low-Scale Orientifold Models the presence of such couplings meets dramatic experimental consequences. It provides new contributions to couplings that may be visible at LHC.

1. Introduction
Recently, many attempts have been made with partial success, in order to embed the Standard Model (SM) in open string theory [1]-[15]. In such a context the Standard Model particles are open string states attached on (different) stacks of D-branes. \( N \) coincident D-branes typically generate a unitary group \( U(N) \). Therefore, every U-factor in the gauge group supplies the model with extra abelian gauge fields.\(^1\)

Such U(1) fields have generically 4d anomalies. The anomalies are cancelled via the Green-Schwarz mechanism [16, 17, 33] where a scalar axionic field (zero-form, or its dual two-form) is responsible for the anomaly cancellation. This mechanism gives a mass to the anomalous U(1) fields and breaks the associated gauge symmetry. The masses of the anomalous U(1)s are typically of order of the string scale but in open string theory they can be also much lighter [18, 19]. If the string scale is around a few TeV, observation of such anomalous U(1) gauge bosons becomes a realistic possibility [21, 22].

As it has been shown in [19, 20], we can compute the general mass formul\(i\) of the anomalous U(1)s in supersymmetric and non-supersymmetric models by evaluating the ultraviolet tadpole of the one-loop open string diagram with the insertion of two gauge bosons on different boundaries. It turns out that U(1) gauge fields that are free of four-dimensional anomalies can still be massive [2, 18, 19, 20]. This is due to the presence of mass-generating six-dimensional anomalies. Since there are decompactification limits in the theory, six-dimensional anomalies affect four-dimensional masses. In six dimensions, two type of fields are necessary to cancel the anomalies, a scalar axion and a two-form. There is also a four-form field but it is dual to the anomalies, scalar axion and a two-form.

\(^1\) There are cases where we can also have SO(\(n\)) or Sp(\(n\)) gauge factors. However, SU(3) can be minimally embedded only in U(3) and in non-minimal cases (bigger gauge groups that are then broken by projections to those of the Standard Model), they leave also other potentially anomalous U(1)s.
scalar. Via the Green-Schwarz mechanism, the pseudoscalar axions give mass to the anomalous U(1) fields. However, the two-forms are not involved in mass generation.

However, Green-Schwarz mechanism is not enough to cancel all the anomalies. Mixed abelian anomalies between anomalous and non-anomalous factors need generalized Chern-Simons terms to be cancelled. Here, we present our results by using a toy-model. A more detailed study can be found in [25].

Next, we provide results of a systematic study of the Standard Model embeddings in type I string theory. We focus in configurations of three and four stacks of D-branes and we discuss the statistics of these models [15]. We focus in two representative models which apart from the Standard Model particles, they also contain three anomalous U(1), three axions and three pairs of Higgs doublets [8, 10]. If these anomalous U(1) gauge boson masses are in the TeV range, they behave like $Z'$ gauge bosons widely studied in the phenomenological literature [21]-[25].

As we mention above, the presence of anomalous and non-anomalous U(1)'s (the hypercharge) requires generalized Chern-Simons terms. These anomaly related couplings produce new signals that distinguish such models from other $Z'$ models. If the string scale is of order of a few TeV, such signals may be visible in LHC [24].

This paper is organized as follows: In section 2 we present the general analysis of the anomaly related effective actions and the Green-Schwarz mechanism. We provide the mass formuli of the anomalous U(1)'s in four- and six-dimensional models. We also argue with a toy-model, that cancellation of mixed anomalies of anomalous and non-anomalous U(1)'s require generalized Chern-Simons terms. In section 3 we explore some three and four stack D-brane models which can describe the Standard Model, and we give some phenomenological aspects of the generalized Chern-Simons terms.

2. Anomalies

Anomalies are generated when classical symmetries are broken at the quantum level [26, 27, 28]. There are Global and Local (Gauge) anomalies. Global anomalies contribute finitely to physical processes. As an example, we remind the decay rate $\pi^0 \rightarrow \gamma\gamma$ that receive contribution from the anomalies providing the correct experimental number for three colored quarks.

Gauge anomalies afflict symmetries necessary to normalize the theory and they must be avoided. The longitudinal polarization of a gauge field related to them does not decouple. The axial Ward-identities contain an anomaly (axial current is not conserved) leading to inconsistencies. Anomalies arise in Parity violating (chiral) theories. This means that left and right handed fermions do not transform in the same way under the gauge symmetry.

In 4d, the anomalous diagram is a triangle with three external bosons. In a theory of gauge group U(N), the group theory factor implies that the possible anomalous diagrams can be:

$$SU(N)^3, \quad U(1) \times SU(N)^2, \quad U(1)^3.$$  

(2.1)

In general, there can also be gravitational anomalies. However, we will not discuss them in the present study. The two last diagrams introduce the concept of the anomalous U(1)s. The Feynman diagrams which contribute to the U(1) anomalies are:

$$\begin{align*}
\begin{array}{c}
\text{U(1)}_j \\
\text{U(1)}_k
\end{array}
\end{align*}$$

$$\begin{align*}
\begin{array}{c}
\text{U(1)}_i \\
\text{G}^{\alpha} \\
\text{U(1)}_i \\
\text{G}^{\alpha} \\
\text{g}_{\mu\nu}
\end{array}
\end{align*}$$

(2.2)

Consider for simplicity only one anomalous U(1). In terms of a gauge transformation $A_{\mu} \rightarrow A_{\mu} + \partial_{\mu}\epsilon$, of the effective action, the anomalies are:

$$\delta S = \int d^4x \left\{ \epsilon \left( A_1 F \wedge F + A_2 \text{Tr}[G_a \wedge G_a] + A_3 R \wedge R \right) \right\},$$

(2.3)
where $F^A, G^a$ field strengths of the anomalous $A^\mu$ and a non-abelian field $G^a_\mu$. Also $A_1 = Tr[Q^3]$, $A_2 = Tr[QT^aT^a]$ and $A_3 = Tr[Q]$ the group theory factors. We suppress the indexes for simplicity. We also remind that $F \wedge H = \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} F_{\mu\nu} H_{\rho\sigma}$.

First, we will concentrate our study in the mixed (second) anomalous diagram and we will describe the Green-Schwarz mechanism that cancels the mixed anomaly between abelian and not-abelian factors [16, 17, 33]. This mechanism introduces a mass for the anomalous U(1). It also generates a Fayet-Iliopoulos potential. However, this mechanism is not enough for cancelling all the anomalies. Mixed abelian anomalies require generalized Chern-Simons terms.

2.1. Green-Schwarz mechanism

In this section we will explore the Green-Schwarz mechanism in 4d. The fields that contribute to the anomaly cancellation are antisymmetric tensors $B^k_{\mu\nu}$ and they are coming from the $k$th twisted closed string spectrum (they are RR fields). However, it is more convenient to use the Poincaré dual of $B_{\mu\nu}$ scalar field $\alpha$ (axion):

$$L_\alpha = \frac{1}{4g_A^2} F^A F^A - \frac{1}{4g_a^2} Tr[G^a G^a] - \frac{1}{2} (\partial^\mu \alpha - M A^\mu)^2 - \frac{1}{2} c_1 \alpha Tr[G^a \wedge G^a].$$

(2.4)

where $M, c_1$ constants. Notice that the third term in the lagrangian is not invariant under a U(1) gauge transformation unless the axion $\alpha$ also transforms like:

$$A^\mu \rightarrow A^\mu + \partial^\mu \epsilon, \quad \alpha \rightarrow \alpha + M \epsilon.$$  

(2.5)

However, this transformation of the axion generates a non-invariance coming from the fourth term in (2.4). This term will annihilate the anomalous term that is generated by the fermionic transformation, giving an anomaly free gauge theory. The total variation of the lagrangian under the above gauge transformation is:

$$\delta L_{\text{total}} = -\frac{1}{2} \left( M c_1 - \frac{A}{16\pi^2} \right) \epsilon Tr[G^a \wedge G^a],$$

(2.6)

where the first term is coming from the variation of $L_\alpha$ and the second are the mixed anomalies from the variation of the measure of the chiral fermions. The anomaly is cancelled for: $A = 16\pi^2 M c_1$.

The NSNS-twisted moduli $m$ (SUSY partner of $\alpha$ which form together a complex scalar field $\phi = m + i\alpha$) couple to the vector fields generating Fayet-Iliopoulos D-terms:

$$S_{FI} = \int d^4 x \frac{1}{g_A^2} \left( m + \sum_i q_i |\Phi_i|^2 \right)^2.$$  

(2.7)

where $\Phi_i$ denote various open strings with charge $q_i$ under the anomalous U(1)s. On the fixed points we have: $\langle m \rangle = 0$. The global U(1)$_A$ remains unbroken despite the fact that the gauge boson became massive [29]. Away from the fixed points we have: $\langle m \rangle \neq 0$. Restoration of SUSY (that is more economical state for the system) implies that the charged scalars will acquire a non-vanishing VEV. This breaks the global U(1)$_A$ symmetry.

2.2. Calculation of the bare mass of the anomalous U(1)s

In this section, we outline the evaluation the contribution to the anomalous U(1) mass for supersymmetric orientifolds. Detailed results can be found in [19].

Closer look to (2.4) shows that these terms are coming from different orders in string perturbation theory. The $(\partial\alpha^i)^2$ is a tree-level (sphere) term, the $A^i \partial\alpha^i$ comes in the disk and
the quadratic term in the gauge fields is a one-loop contribution. To clarify this, we mention that $g_\varphi^2$ is proportional to $g_s = e^{\phi}$ and every power of the axion absorbs a dilaton factor $e^{-\phi}$ because it is a RR filed. The string perturbation series are weighted by $g_s^\chi$ where $\chi = 2 - 2h - c - b$ is the Euler character and $h, c$ and $b$ denote the handle, the cross-cups and the boundaries of a closed orientable Riemann surface respectively.

The diagrams at one-loop that contribute to terms quadratic in the gauge bosons (anomalous U(1)s) are the genus-one surfaces with boundaries: the annulus and the Möbius strip. In the infrared (IR) region they diverge logarithmically and give the logarithmic running of the couplings. In the ultraviolet (UV) region the tadpoles of the annulus with both gauge bosons inserted in the same boundary, and the Möbius strip vanish due to the tadpole cancellation. In this UV limit the annulus amplitude with the gauge bosons inserted in opposite boundaries provides the mass-term of the anomalous U(1) [19]. Since we are interested in the anomalous gauge boson mass, we concentrate on the latter diagram. The gauge boson vertex operator is

$$\tilde{V}^a = \lambda^a \zeta^\mu (\partial X^\mu + i(p \cdot \psi)\psi^\mu)e^{ip \cdot X},$$

where $\lambda$ is the Chan-Paton matrix and $\zeta^\mu$ is the polarization vector. The 2-point annulus amplitude is given by

$$A^{ab} = -\frac{1}{4G} \int [d\tau][dz] \int \frac{d^dp}{(2\pi)^d} \sum_k \left\langle \tilde{V}^a(\zeta_1, p_1, z) \tilde{V}^b(\zeta_2, p_2, z_0) \right\rangle_k,$$

where we keep undetermined the number of non-compact dimensions $d$, and $G$ denotes the order of the orientifold group. The fundamental polygon of the annulus is $[0, t/2] \otimes [0, 1/2]$. The index $k$ denotes the various orbifold sectors that we may have. Using the translation symmetry of the annulus, we fix the position of one VO to $z_0 = 1/2$. The other VO is placed on the imaginary axis with $z \in [0, t/2]$.

Strictly speaking, the amplitude above is zero on-shell if we enforce the physical state conditions $\zeta \cdot p = p^2 = 0$ and momentum conservation $p_1 + p_2 = 0$. There is however a consistent off-shell extension, without imposing momentum conservation, that has given consistent results in other cases (see [35] for a discussion). We thus impose momentum conservation only at the end of the calculation.

It appears that the amplitude (2.9) has a kinematical multiplicative factor that is $O(p^2)$, thus would seem to provide a leading correction only to the anomalous gauge boson coupling.

Following the procedure of [19], we find that terms in the correlation functions which are spin-structure independent vanish. The only spin-dependant term lies in the fermionic correlation function. After integration over the position $z$ and the annulus modulus $t$, a term proportional to $1/p_1 \cdot p_2$ appears from the ultraviolet (UV) region (as a result of the quadratic UV divergence in the presence of anomalous U(1)s) that provide the mass-term.

**General formuli for the disk couplings of axions to gauge bosons**

We would like to illustrate the general results in some concrete examples such as type I compactifications on $T^6/Z_N$ orbifolds. Therefore, we would like to introduce some notation.

The resulting Chan-Paton group is typically non semi-simple and indeed contains one (for $N = 3, 7$) or more abelian factors which are all superficially anomalous. When $N$ is even, there are $Z_2$ elements $\Omega$ in the orbifold group. The $\Omega \overline{\Omega}$ involution where $\Omega$ is the (generalized) world-sheet parity, generates $O5$-planes in the configuration. D5-branes are then needed for tadpole cancellation and the gauge group comprises two different kinds of gauge groups.

Denoting by $\gamma$ the discrete Wilson lines, projectively embedding the orbifold group into the Chan-Paton group, we may parameterize them as follows:

$$\gamma_1^a = \exp(-2\pi i \otimes_r V_r^a \cdot H_r)$$

(2.10)
where $H_r$ are the Cartan generators of $SO(32)^\alpha$ with $\alpha = 9, 5$. They are normalized to $tr(H_r H_s) = 2\delta_{rs}$. Tadpole conditions fixes the values of $V_r^\alpha$.

In order to study the fate of the anomalous $U(1)$’s, it is convenient to introduce the combinations

$$\lambda_i = \frac{1}{2\sqrt{m_i}} \sum_{r=1}^{n_i} Q_i^r H_r$$

where $i$ denotes the brane and $Q_i^r = (0, 0, \ldots, 0, 1, \ldots, 1, 0, \ldots, 0)$ are 16-dimensional vectors, with $n$ one-entries at the position where the corresponding $U(n)$ lives. Notice that $\lambda$’s satisfy $tr[\lambda_i \lambda_j] = \frac{1}{2}\delta_{ij}$.

Finally, the masses of the anomalous $U(1)$’s for various orientifold sectors are:

- $\mathcal{N} = 1$ Sectors: The contribution to the masses for $\mathcal{N} = 1$ sectors of $Z_N$ orbifolds, labelled by $k$, are (we assume that the D5-branes are longitudinal to the $T_2^2$):

$$\frac{1}{2} M_{90, k}^2 = \frac{1}{2} M_{55, k}^2 = -\frac{1}{8\pi^3 N} N_k^1 N_k^2 N_k^3 tr[\gamma_k \lambda^a] tr[\gamma_k \lambda^b]$$

$$\frac{1}{2} M_{95, k}^2 = \frac{\tilde{\eta}_k}{8\pi^3 N} N_k^3 tr[\gamma_k \lambda^a] tr[\gamma_k \lambda^b]$$

where $\tilde{\eta}_k = \text{sign} \left( \prod_{\Lambda=1}^3 \sin[\pi k v_\Lambda] \right)$ when all tori are twisted and $(-1)$ when a perpendicular torus to the D5 brane remains untwisted by the orbifold action. Also $N_k^i = (2\sin[\pi k v_i])^2$ is the number of the effective fixed points of torus $T_i^2$.

- $\mathcal{N} = 2$ Sectors: For such sectors, one $v_i k$ is integer $i.e.$ one torus is untwisted by the orbifold action. This torus can be longitudinal or perpendicular to the D5 branes. Without loss of generality, we can assume that the longitudinal torus to the D5 brane is $T_i^2$ and the not untwisted-perpendicular one (if any) is $T_j^2$.

Therefore, the contribution to the masses for $\mathcal{N} = 2$, $k$ sectors of $Z_N$ orbifolds are:

$$\frac{1}{2} M_{90, k}^2 = \frac{1}{2} M_{55, k, ||}^2 = -\frac{2V_3}{4\pi^3 N} N_k^1 N_k^2 \sqrt{N_k^3} tr[\gamma_k \lambda^a] tr[\gamma_k \lambda^b]$$

$$\frac{1}{2} M_{55, k, \perp}^2 = -\frac{(2V_2)^{-1}}{4\pi^3 N} N_k^1 N_k^2 \sqrt{N_k^3} tr[\gamma_k \lambda^a] tr[\gamma_k \lambda^b]$$

$$\frac{1}{2} M_{95, k, ||}^2 = \frac{2V_3}{4\pi^3 N} \tilde{\eta}_k tr[\gamma_k \lambda^a] tr[\gamma_k \lambda^b]$$

where $\tilde{\eta}_k = (-1)^{k v_3}$ and $V_i$ denotes the volume of the internal torus $T_i^2$. Notice that $||$ and $\perp$ denote that the $k$th sector leaves invariant the longitudinal (third) or a perpendicular (second) torus to the D5 brane$^2$.

We should mention, that the above masses are unnormalized. To obtain the normalized mass matrix, we must also take into account the kinetic terms of the U(1) gauge bosons which are

$$S_{\text{kinetic}} = -\frac{1}{4g_s} \left[ V_1 V_2 V_3 \sum_i F_i^2 + V_3 \sum_j \tilde{F}_j^2 \right].$$

where $i$ and $j$ denote the gauge groups that are coming from different stacks of D9 and D5-branes. This implies $M_{90}^2 \rightarrow M_{90}^2/(V_1 V_2 V_3)$, $M_{55}^2 \rightarrow M_{55}^2/V_3$ and $M_{95}^2 \rightarrow M_{95}^2/(V_3 \sqrt{V_1 V_2})$.

$^2$ As an example consider the $Z_6$ orientifold which has vector $v = (1, -3, 2)/6$. Tadpole condition implies D9 branes and D5 branes which are longitudinal to the $T_3^2$. The $k = 2, 3$ are $\mathcal{N} = 2$ sectors and the contribution to $M^2_{90}$ is given by (2.15), (2.14) respectively.
2.3. The structure of six-dimensional mixed gauge anomalies

In the previous section we outlined the computation of the bare masses of the anomalous U(1)s by evaluating the ultraviolet tadpole of the one-loop open string diagram with the insertion of two gauge bosons on different boundaries. It turns out that U(1) gauge fields that are free of four-dimensional anomalies can still be massive. This is unexpected and we should study the contribution of higher anomalies in the mass-generation of the U(1)s. We will especially study the six-dimensional anomalies since we cannot have eight-dimensional anomalies in orientifold models that obey the condition $\sum_i v_i = 0$ (SUSY condition). If there are decompactification limits in the theory, six-dimensional anomalies affect four-dimensional masses.

In six dimensions, the leading diagram that can give a contribution to anomalies is the square diagram [17]. The mixed group theory factors that do not identically vanish are these with two or three external non-abelian gauge bosons. The Feynman diagrams that eventually contain anomalies are:

Therefore, in the presence of an anomalous U(1) field, the effective action is not invariant under a transformation $\delta A^i = d\epsilon^i$:

$$\delta \epsilon S|_{\text{gauge}} = \int d^6x \left\{ \epsilon^i \left( A_{QQTT} F^i \wedge Tr[G^2] + A_{QTTT} Tr[G^3] \right) \right\} , \quad (2.18)$$

where $A_{QQTT} = Tr[Q_i Q_j T^\alpha T^\alpha]$, $A_{QTTT} = Tr[Q_i T^\alpha \{ T^\beta T^\gamma \}]$ the group theory factors. Powers of forms are understood as wedge products. We denote by $G_{\mu\nu}$ the field strength of a non-abelian gauge field $G^\alpha$.

Gauge invariance is preserved by the six-dimensional Green-Schwarz mechanism. However, two inequivalent fields should contribute to this cancellation. The cancellation of the first anomalous term is arranged by a 2-form $B^i$ (RR twisted field) which transforms under the U(1) transformation like $\delta B^i = -\epsilon^i F^i$. The lagrangian of this field is:

$$S_{QQTT} = \int d^6x \left[ -\frac{1}{4g^2_i} F_{\mu\nu}^2 - \frac{1}{12} [dB^i + \Omega_{\alpha}]^2 + A_{QQTT} B^i \wedge Tr[G^2] \right] , \quad (2.19)$$

where the last term is proportional to the anomaly of the first diagram. The 3-form $\Omega_{\alpha} = A^i dA^i$ is the Chern-Simons term of the abelian gauge field $A^i_\mu$. This part of the action does not generate a mass for the gauge boson.

The second anomaly is cancelled by a pseudoscalar axion that transforms under the U(1) transformation as $\delta \alpha^i = -M\epsilon^i$:

$$S_{QTTT} = \int d^6x \left[ -\frac{1}{4g^2_i} F_{\mu\nu}^2 + \frac{1}{2} (d\alpha^i + MA^i)^2 + \frac{A_{QTTT}}{M} \alpha^i Tr[G^3] \right] . \quad (2.20)$$

This action supplies a mass term for the U(1) gauge field and breaks the gauge symmetry in six dimensions.

**Six-dimensional mass formulii**

The general mass formulii for the anomalous U(1) gauge fields in six-dimensional orientifolds can be easily evaluated in the same way that we did for the four-dimensional cases [20]. $\mathcal{N}=1$
six-dimensional orientifolds are created as $T^4/Z_N$ where $N = 2, 3, 4, 6$. The results for strings attached on the same kind of branes (untwisted states) are:

$$\frac{1}{2} M_{\alpha \alpha}^2 = -\frac{1}{\pi^2 N} \sum_k \left( \frac{2 \sin \frac{\pi k}{N} }{N} \right)^2 Tr[\gamma_k \lambda^\alpha] Tr[\gamma_k \lambda^\alpha] ,$$

(2.21)

where $a = 5, 9$ denotes the kind of D-branes on which the open string is attached. In the case where strings have one end on a $D5$ and the other on a $D9$-brane (twisted states) we have:

$$\frac{1}{2} M_{99}^2 = -\frac{1}{\pi^2 N} \sum_k Tr[\gamma_k \lambda^5] Tr[\gamma_k \lambda^9] .$$

(2.22)

We should mention, that the above masses are again unnormalized. To obtain the normalized mass matrix, we must also take into account the kinetic terms of the U(1) gauge bosons which are again (2.17). However, the volume of the torus that the D5-branes is longitudinal to, should be normalized to identity. This implies $M_{99}^2 \rightarrow M_{99}^2/(\sqrt{V_1 V_2})$, $M_{69}^2 \rightarrow M_{69}^2/(\sqrt{V_1 V_2})$, $M_{59}^2 \rightarrow M_{59}^2/(\sqrt{V_1 V_2})$.

2.4. Generalized Chern-Simons: A toy model

In this section, we will argue that the Green-Schwarz mechanism is not enough to cancel all the anomalies in a model. In particular, generalized Chern-Simons terms are necessary to cancel mixed anomalies between anomalous and non-anomalous U(1)s [25]. We will concentrate on a toy-model with two U(1)s, one anomalous with gauge field $A$, field strength $F^{A}_{\mu\nu}$ and charge operator $Q_A$, and the other non-anomalous with gauge field $Y_\mu$, field strength $F_{\mu\nu}^Y$ and charge operator $Q_Y$. $G_{\mu\nu}^a$ denote as before other non-abelian field strengths. In general:

$$Tr[Q_Y] = 0, \quad Tr[Q_Y^3] = 0, \quad Tr[Q_Y T^a T^a] = 0$$

$$Tr[Q_A Q_Y^2] = c_1, \quad Tr[Q_A^2 Q_Y] = c_2, \quad Tr[Q_A^3] = c_3, \quad Tr[Q_A T^a T^a] = \xi .$$

(2.23)

(2.24)

By definition non-anomalous U(1)’s obey conditions (2.23). However, mixed traces with the non-anomalous U(1) might be different from zero (2.24). This structure can emerge in type I theories after recombining some U(1) factors arising from the ‘center of masses’ of brane configurations. In particular we have in mind intersecting brane embeddings of (some supersymmetric extension of) the standard models where the non anomalous U(1) $Y$ is a combination of the CP U(1)’s with $Tr[Q_Y] = Tr[Q_Y^3] = 0$ but in general $Tr[Q_A Q_Y] \neq 0$.

The above traces imply the following anomalous transformations of the one-loop effective action. Under

$$A_\mu \rightarrow A_\mu + \partial_\mu \epsilon, \quad Y_\mu \rightarrow Y_\mu + \partial_\mu \zeta$$

(2.25)

the action transforms as:

$$\delta S_{1-loop} = \int d^4x \left\{ \epsilon \left[ \frac{c_2}{3} F^A \wedge F^A + c_2 F^A \wedge F_Y + c_1 F^Y \wedge F^Y + \xi Tr[G \wedge G] \right] \right.$$

$$+ \zeta \left[ c_2 F^A \wedge F^A + c_1 F^A \wedge F^Y \right] \}$$

(2.26)

Following Green-Schwarz mechanism, we add the action

$$S_{\text{axion}} = \int d^4x \left\{ - \frac{1}{4g_Y^2} (F^Y)^2 - \frac{1}{4g_A^2} (F^A)^2 + (\partial_\mu \alpha + MA_\mu)^2 \right.$$ 

$$+ \alpha (d_3 F^A \wedge F^A + d_2 F^A \wedge F_Y + d_1 F^Y \wedge F^Y + d_0 trG \wedge G) \}$$

(2.27)
where $d_0$, $d_1$, $d_2$, $d_3$, $M$ are constants. The axion $\alpha$ transforms as
\begin{equation}
\alpha \rightarrow \alpha - M \epsilon .
\end{equation}
and we are assuming that $\alpha$ does not shift under non-anomalous gauge transformations parameterized by $\zeta$. Although $Q_A$ and $Q_Y$ mix in a certain sense ($\text{Tr}[Q_A Q_Y] \neq 0$), we will confirm that $\alpha$ does not couple à la Stückelberg with $Y$.

It is obvious that the axionic transformation does not cancel all the anomalies. Therefore, it is necessary to add not invariant $\zeta$ terms, generalized Chern–Simons terms:
\begin{equation}
S_{\text{GCS}} = \int Y \wedge A \wedge \left\{ d_4 F^A - d_5 F^Y \right\}
\end{equation}
where all $d_i$ are constants. The gauge variation of the classical action (now $S_{\text{class}} = S_{\text{axion}} + S_{\text{GCS}}$):
\begin{align}
\delta S_{\text{class}} &= -\int \epsilon \left\{ d_3 F^A \wedge F^A + (d_2 - d_4) F^A \wedge F^Y + (d_1 + d_3) F^Y \wedge F^Y + d_0 \text{tr}(G \wedge G) \right\} \\
&\quad - \int \zeta \left\{ d_4 F^A \wedge F^A - d_5 F^Y \wedge F^A \right\} .
\end{align}

Anomaly cancellation implies:
\begin{equation}
d_0 = \xi , \quad d_1 = 2c_1 , \quad d_2 = 2c_2 , \quad d_3 = \frac{c_3}{3} , \quad d_4 = c_2 , \quad d_5 = -c_1 .
\end{equation}
The presence of the generalized CS terms seems due to the non-vanishing of $c_1$ and $c_2$.

This is a generic situation. Anomalous and non-anomalous factors usually appear in orientifolds. The presence of GCS terms is necessary to cancel all the anomalies.

**General formuli of generalized Chern–Simons terms in orientifolds models**

In the general case, generalized Chern–Simons terms appear in the effective action as
\begin{equation}
\frac{1}{24\pi^2} E_{ij,l} \int A^i \wedge A^j \wedge F^l
\end{equation}
where $ijl$ span the $U(1)$s. $E_{ij,k}$ is antisymmetric in $ij$. For gauge groups coming from D-branes in type II orientifold models, $E_{ij,k}$ can arise only from a non-planar cylinder diagram that contains the (antisymmetrized) Chan-Paton traces:
\begin{equation}
E_{ijl} = \frac{1}{3} \sum_k \eta_k |\sqrt{N_k}| \text{tr}[\gamma_k \lambda_i \lambda_j] \text{tr}[\gamma_k \lambda_l] .
\end{equation}
Here $k = 1 \cdots N - 1$ denotes the different type of twisted sectors propagating in the tree-level channel cylinder diagram, whereas
\begin{equation}
N_k = \begin{cases} 
\prod_{\Lambda=1}^3 (2 \sin[\pi k \nu_{\Lambda}])^2 & \text{for D9 – D9 and D5 – D5 sectors}, \\
(2 \sin[\pi k \nu_3])^2 & \text{for D9 – D5 sectors}
\end{cases}
\end{equation}
denote the number of fixed points in the internal space and in the third internal torus, respectively (we consider for simplicity D5 branes whose world-volume span the third internal torus $T^3_3$). Also, $\eta_k$ takes the values of: $\text{sign}(\prod_{\Lambda=1}^3 \sin[\pi k \nu_{\Lambda}])$ for all sectors of D9-D9, D5-D5, D9-D5 where the orbifold action twists all tori, $(-1)^{k\nu_3}$ for all sectors of D9-D5 where the orbifold
action leaves untwisted a perpendicular torus $T^2$ to the D5 brane (all the above are $\mathcal{N} = 1$ sectors), and zero for sectors of D9-D9, D5-D5, D9-D5 where the orbifold action leaves untwists the longitudinal torus $T^2$ to the D5-brane (which are $\mathcal{N} = 2$ sectors). Notice that particles and antiparticles contribute to the anomaly with different signs as it should be. Let us stress that the interpretation of the factors $N_k$ is different for D9 and D5 branes. Whereas D9 branes fill the whole space-time and therefore couple to twisted axions localized at all fixed points, the D5 branes can only probe some fixed points and their associated axions. Correspondingly, their couplings to such axions are different with respect to the D9 brane couplings.

All open string models which could eventually describe the Standard Model contain extra abelian factors, one for each stack of branes. GCS terms provide very particular anomaly-related couplings to such axions are different with respect to the D9 brane couplings. D5 branes can only probe some fixed points and their associated axions. Correspondingly, their couplings to such axions are different with respect to the D9 brane couplings.

In this section we study D-brane configuration that can describe the Standard Model. The ten-dimensions of string theory are split into four flat non-compact and six compact dimensions. D-branes are inserted and they are longitudinal to the four non-compact dimensions. In addition, there are Orientifold planes, non-dynamical hyperplanes which change the orientation of the strings and are required for the consistency and the stability of the theory [30]. Since open strings are proportional to their lengths, the branes that give rise to the SM must be very close together in the internal space. Therefore, we can focus in this particular area since all other branes further away may affect the global (the stability and consistency of the configuration) rather than the local properties of the model. These are called the bottom-up approaches.

If the end of a string is attached on a stack of $m$ branes it can take $m$ different values. It is easy to see that an oriented string with both ends on the same stack of branes can take $m \times m$ and transforms in the Adjoint of an $U(m)$ group. However, we can split $U(m) \equiv SU(m) \times U(1)\pi_m$ where the index $m$ in the $U(1)$ denotes that the abelian factor is coming from the $m$-stack of branes. An oriented string which starts from an $m$ and ends on an $n$ stack transforms in the bifundamental $(m, n) \equiv (m, 1_m; n, -1_n)$ in the two representations. As we mentioned before, unoriented strings are those which pass through an O-plane. Therefore, we can have unoriented strings which transform as $(m, n) \equiv (m, 1_m; n, 1_n)$ or $(\bar{m}, \bar{n}) \equiv (\bar{m}, -1_m; \bar{n}, -1_n)$. Finally, unoriented strings with both ends on the same stack of branes transform in the antisymmetric or symmetric representation of $U(m)$, depending on the O-plane: $\begin{pmatrix} m \end{pmatrix} \equiv (\bar{m}, 2_m)$ and $\begin{pmatrix} n \end{pmatrix} \equiv (\bar{m}, 2_m)$. All the above are illustrated in figure 1.

Keeping all these in mind we can study D-brane configurations that describe the SM spectrum which contains:

$$3 \times \left[ Q \left(3, 2, \frac{1}{6} \right) + u^c \left(3^*, 1, -\frac{2}{3} \right) + d^c \left(3^*, 1, \frac{1}{3} \right) + L \left(1, 2, -\frac{1}{2} \right) + \ell^c \left(1, 1, 1 \right) \right]$$

and in addition, three left-handed anti-neutrinos $\nu^c$ in the representation $(1, 0, 0)$, and three MSSM Higgs pairs, $H \left(1, 2, \frac{1}{2} \right) + H' \left(1, 2, -\frac{1}{2} \right)$.

3.1. Three stacks: $U(3) \times U(2) \times U(1)$ Models

The minimal D-brane configuration that can describe the SM consists from three stacks of branes, one with 3 (the “QCD” branes), one with 2 (the “weak” branes) and a single brane. We denote these stacks with $a, b, c$ respectively. Obviously, the $SU(3)_c \times SU(2)_L$ part of the SM is described by the non-abelian factors of the QCD-a and weak-b branes respectively and the
Figure 1. Possible representations that can appear in orientifold models. Notice the change of orientation of the strings that pass throw O-planes, which are denoted by the small gray segment.

The hypercharge $U(1)_Y$ is a linear combination of the three abelian factors which are coming from each stack of branes.

The quark doublet $Q$ $(u, d)$ is stretched between the QCD and the weak branes. The antiquarks $u^c$, $d^c$ can be stretched between the QCD and a single brane. However, a new possibility emerge due to the change of orientation (O-plane). There can be also strings which pass through an O-plane and have both ends attached onto the QCD branes. The last possibility provides an antisymmetric representation of the SU(3) and breaks baryon number symmetry, which is related to the abelian charge coming from the QCD stack $B = Q_a/3$. The $L$ $(\nu, l)$ is stretched between the weak and the single brane. Similarly for the Higgses. However, we should mention that we can allow $H$ and $H'$ to be conjugate to each other. This possibility will be explicitly noted in the next models. The lepton singlet $l^c$ could have both ends onto the weak (providing an antisymmetric representation of SU(2)) or onto the single branes (bifundamentals). Finally, the right-handed netrino $\nu^c$ can come from the hidden sector or from the configuration. However, if it is described by a string attached onto the three stacks of branes, it can be stretched between the weak or single branes. In our study we separate these two possibilities.

Requiring that the particles have the proper hypercharge we find two possible ways to embed the SM in this D-brane system of three stacks\(^3\). We also provide the corresponding charge assignments for the SM particles:

$$Y = -\frac{1}{3}Q_a + \frac{1}{6}Q_b$$

$$M = -\frac{1}{6}Q_a + \frac{1}{2}Q_c$$

| $Q$ | $(V, V, 0)$ | $Q$ | $(V, \bar{V}, 0)$ |
| $u^c$ | $(A, 0, 0)$ | $u^c$ | $(V^*, 0, V^*)$ |
| $d^c$ | $(V^*, 0, \bar{V})$ | $d^c$ | $(A, 0, 0)$ or $(V^*, 0, V)$ |
| $L^c$ | $(0, V^*, \bar{V})$ | $L$ | $(0, \bar{V}, V^*)$ |
| $H$ | $(0, V, \bar{V})$ | $H$ | $(0, \bar{V}, V)$ |
| $H'$ | $(0, V^*, \bar{V})$ | $H'$ | $(0, \bar{V}, V^*)$ |
| $[\nu^c]$ | $(0, 0, \bar{S})$ | $[\nu^c]$ | $(0, \bar{A}, 0)$ |

The right-hand neutrino is placed in brackets to emphasize the fact that it can also come from

\(^3\) By tilde representation $R$ we denote the possibility of fundamental or conjugate representation, $\tilde{R} \to (R$ or $R^*)$
the hidden sector. Notice the ambiguity of the representations (tilted) when a brane does not contribute to the hypercharge, and also the two different possibilities for the charges of $d^c$: $(V^*, 0, V)$ or $(A, 0, 0)$ in the second case.

From the above charge assignments we can construct families and search for triplets of these families where all cubic anomalies cancel. The conditions for these cancellations are derived from tadpole cancellation and they are the usual ones for the non-abelian subgroups of $U(N)$, $N > 2$: Vectors contribute 1, symmetric tensors $N + 4$ and anti-symmetric tensors $N - 4$, and conjugates contribute with opposite signs. However, the same condition emerges even if $N = 1$ and $N = 2$. This means that for example a combination of three vectors and an anti-symmetric tensor is allowed in a $U(1)$ factor.

For that first embedding there are 10 different anomaly-free spectra which describe the SM. For the second, there are 24 different anomaly-free models.

The baryon number $B = Q_a/3$ is a gauge symmetry only in models where $d^c$ is string with the two ends onto different branes. The lepton number is also a linear combination of the abelian factors. However, solving the resulting systems we find that in none of the above models lepton number is a symmetry.

3.2. Four stacks: $U(3) \times U(2) \times U(1) \times U(1)'$ Models

In this section, we provide the statistics of four-stack realizations of the SM, without getting in details (the fourth single brane is denoted by $d$). The possible embeddings of the hypercharge to the four abelian factors are eight [15]:

- Hypercharge $Y = (x - \frac{1}{3})Q_a + (x - \frac{1}{2})Q_b + xQ_c + (x - 1)Q_d$

  Here the right-handed neutrino must necessarily arise in the hidden sector. Following the same spirit as in the tree-stack models, we can form families from the corresponding charge assignments and require that triplets of them are free of irreducible anomalies. For the present hypercharge embedding there is only one anomaly-free model which can describe the SM.

- Hypercharge $Y = -\frac{1}{3}Q_a - \frac{1}{2}Q_b - Q_d$

  If $\nu^c$ is coming from the hidden sector, there are 302 anomaly-free models which can describe the SM particles. Among them, there are 62, 72, 96 and 72 models with tree, two, one and none chiral Higgs pairs. On the other hand, if $\nu^c$ is attached onto branes of the above stacks, it can only transform as a symmetric rep of $U(1)c$ which does not contribute to the hypercharge. In that case, there are 1208 different anomaly-free models which can describe the SM particles (including $\nu^c$). Among them, there are 240, 384, 288 and 248 models with tree, two, one and none chiral Higgs pairs.

- Hypercharge $Y = \frac{1}{2}Q_a + \frac{1}{2}Q_b + Q_c$

  In total, there are 6 different anomaly-free models which can describe the SM particles with chiral Higgs-pairs and the left-hand neutrino coming from the hidden sector. A $\nu^c$, as a string attached on this stack of branes, can only transform as a symmetric of $U(1)'$. In that case, there are 24 different anomaly-free models with chiral Higgs-pairs (including $\nu^c$) and they all have baryon number $B = Q_a/3$.

- Hypercharge $Y = \frac{1}{2}Q_a + \frac{1}{2}Q_c - \frac{1}{2}Q_d$

  There are 8552 different anomaly-free models with chiral Higgs pairs which can describe the SM particles. If the right-handed neutrino $\nu^c$ is attached onto the SM branes, it can be described by an antisymmetric of the $U(2)$ or a string stretched between the two single branes. Including $\nu^c$, there are 150672 different anomaly-free models. Among them, there are 29360, 61344, 48800 and 11168 models with tree, two, one and none chiral Higgs pairs.

- Hypercharge $Y = \frac{1}{6}Q_a + \frac{1}{2}Q_c - \frac{3}{2}Q_d$
In that case, there are 4 different anomaly-free models with chiral Higgs Pairs which can describe the SM. A $\nu^c$ which is stretched between the four stacks can only transform in the antisymmetric of $U(2)$. Including $\nu^c$, the number of different charge assignments is 24 (8 of them have two chiral Higgs pairs and the other 16 have non chiral Higgs pairs). Half of these states have baryon number $Q_B = Q_a/3$ and in none lepton number is a symmetry. All models have one non-anomalous $U(1)$.

- **Hypercharge** $Y = -\frac{1}{3} Q_a - \frac{1}{2} Q_b$
  
  There are 936 anomaly-free models with $\nu^c$'s coming from the hidden sector. Among them, there are 256, 120, 120 and 440 models with tree, two, one and none chiral Higgs pairs. A $\nu^c$ which will be stretched between the four branes will transform in a symmetric rep of the single branes or it might be stretched between them. Including $\nu^c$, there are 106792 different anomaly-free models. Among them, there are 15072, 32332, 36228 and 23160 models with tree, two, one and none chiral Higgs pairs.

- **Hypercharge** $Y = -\frac{5}{6} Q_a - Q_b - \frac{1}{2} Q_c + \frac{3}{2} Q_d$
  
  In that case, there are 2 different anomaly-free models which can describe the SM and they have baryon number $Q_B = Q_a/3$. Lepton number is not a symmetry.

- **Hypercharge** $Y = \frac{5}{6} Q_a + Q_b + \frac{1}{2} Q_c + \frac{1}{2} Q_d$
  
  The above hypercharge embedding is allowed only in cases where the right-handed neutrino is coming from the hidden sector. In that case, there are 2 different anomaly-free models which can describe the SM particles and they have baryon number $Q_B = Q_a/3$. Lepton number is not a symmetry.

3.3. $Z'$s in D-brane models

In this section we will investigate the effective action of D-brane models focusing in two representative models (similar properties hold for all above models). Therefore, consider the following models:

\[
Y = \frac{1}{3} Q_a - \frac{1}{2} Q_b - Q_d \\
B = \frac{1}{3} Q_a \\
L = \frac{1}{2} (Q_a + Q_b - Q_c - Q_d) \\
PQ = -\frac{1}{2} (Q_a - Q_b - 3Q_c + 3Q_d) \\
\]

\[
Y = \frac{2}{3} Q_a + \frac{1}{2} Q_b + Q_c \\
B = \frac{1}{3} Q_a \\
L = -\frac{1}{2} (Q_a + Q_b + Q_c + Q_d) \\
PQ = -\frac{1}{2} (Q_a + 3Q_b - Q_c - Q_d) \\
\]

where $Y$, $B$, $L$, $PQ$ are the hypercharge, the baryon number, the lepton number and the Peccei-Quinn global symmetries respectively. We can easily realize that only the hypercharge $Y$ is free
of anomalies. Therefore, apart from the SM particles, we have three more U(1)’s, three axions that mix with the U(1)’s and three pairs of Higgs doublets.

Next, we would like to diagonalize the mass matrix of the U(1)’s and go to the so called photon basis. There are two mass-origins for the U(1)’s. One is the electroweak symmetry breaking and the other the Stückelberg anomaly-related terms [24]:

$$|D_\mu H|^2 + |D_\mu H'|^2 + \frac{1}{2} \sum I (\partial \alpha_I + M_I A_I)^2$$

(3.37)

where $D_\mu H = (\partial_\mu + \frac{i}{2} g_2 \tau^a W_\mu^a + \frac{i}{2} g_Y A_\mu^Y + \frac{i}{2} \sum_I q_I^I g_1 A_\mu^I) H$ (similarly for $H'$). Index $I$ spans on the three anomalous U(1)’s. We should mention that the Higgses are charged only under $Y$ and $PQ$ and after electroweak symmetry breaking only $Y$ and $PQ$ are spontaneously broken. Therefore:

- The photon $A$, the $Z^0$ and the PQ-related $Z'$-boson are the three eigenstates of the mass matrix of $Y$, $PQ$ and $W_3$. This change of basis

$$\begin{pmatrix} W^3 \\ Y \\ PQ \end{pmatrix} = \begin{pmatrix} c_{11} & c_{12} & c_{13} \\ c_{21} & c_{22} & c_{23} \\ c_{31} & c_{32} & c_{33} \end{pmatrix} \begin{pmatrix} A \\ Z^0 \\ Z' \end{pmatrix}$$

(3.38)

where

$$c_{11}, c_{12}, c_{21}, c_{22}, c_{33} \sim O(1) \quad c_{13}, c_{23}, c_{31}, c_{32} \sim O \left( \frac{M_s^2}{M_s^2} \right) < 10^{-4}$$

(3.39)

- On the other hand the $B$ and $L$ gauge bosons are not affected by the Higgs mechanism. They give two extra massive $Z'$ gauge bosons with masses proportional to the string scale $M_s$.

As we have argued in section 2.4, generalized Chern-Simons terms are necessary to cancel all the anomalies, since the above models contain anomalous and non-anomalous U(1)’s. These couplings are new and they distinguish these D-brane models from other $Z'$ models which have been studied in the past. Consider for example $PQ \wedge Y \wedge dY$, which after a rotation to the photon basis it generates various anomaly cancelling Chern-Simons-like terms:

$$\begin{array}{rcl}
PQ \wedge Y \wedge dY & \longrightarrow & \begin{cases} 
Z^0 \wedge A \wedge dA & \Rightarrow & Z^0 \rightarrow \gamma \gamma & \sim & O \left( \frac{M_s^2}{M_s^2} \right) \\
A \wedge Z^0 \wedge dZ^0 & \Rightarrow & Z^0 \rightarrow Z^0 \gamma & \sim & O \left( \frac{M_s^2}{M_s^2} \right) \\
Z' \wedge A \wedge dA & \Rightarrow & Z' \rightarrow \gamma \gamma & \sim & O (1) \\
Z' \wedge Z^0 \wedge dZ^0 & \Rightarrow & Z' \rightarrow Z^0 Z^0 & \sim & O (1) \\
Z' \wedge Z^0 \wedge dA & \Rightarrow & Z' \rightarrow Z^0 \gamma & \sim & O (1) 
\end{cases}
\end{array}$$

(3.40)

We should mention that some decays vanish on-shell. However, $Z' \rightarrow Z \gamma$ and $Z' \rightarrow ZZ$ produce new signals that distinguish such models from other $Z'$ models. If the string scale is of order of a few TeV, such signals may be visible in LHC.

4. Conclusions

As we have mentioned, all open string models that approach the Standard Model contain anomalous U(1) gauge fields. The anomaly is cancelled via the Green-Schwarz mechanism that
generates a mass for the corresponding anomalous gauge boson. We provide the bare masses of the anomalous U(1)s in four-dimensional supersymmetric orientifolds. However, there are cases where even non-anomalous U(1)s acquire a mass due to six-dimensional anomalies that upon decompactifications affect the four-dimensional theory.

However, Green-Schwarz mechanism is not enough to cancel all the anomalies. Mixed abelian anomalies between anomalous and non-anomalous factors need generalized Chern-Simons terms to be cancelled.

Next, we provide results of a systematic study of the Standard Model embeddings in type I string theory at TeV scale. We focus in configurations of three and four stacks of D-branes and we discuss the statistics of these models. We focus in two representative models which apart from the Standard Model particles, they also contain three anomalous U(1), three axions and three pairs of Higgs doublets. The presence of anomalous and non-anomalous U(1)’s requires generalized Chern-Simons terms. These anomaly related couplings produce new signals that distinguish such models from other $Z'$ models. If the string scale is of order of a few TeV, such signals may be visible in LHC.

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