A D-brane alternative to unification

I. Antoniadis $^a$, E. Kiritsis $^b$ and T.N. Tomaras $^b$

$^a$Centre de Physique Théorique, Ecole Polytechnique, 91128 Palaiseau, France

$^b$Department of Physics, University of Crete, and FO.R.T.H.
P.O. Box 2208, 710 03 Heraklion, Crete, Greece

ABSTRACT

We propose a minimal embedding of the Standard Model spectrum in a D-brane configuration of type I string theory. The $SU(3)$ color and $SU(2)$ weak interactions arise from two different collections of branes. The model is neither grand unified nor supersymmetric but it naturally leads to the right prediction of the weak angle for a string scale of the order of a few TeV. It requires two Higgs doublets and guarantees proton stability.
String theory is the only known framework for quantizing gravity. If its fundamental scale is of the order of the Planck mass, stability of the hierarchy of the weak scale requires low energy supersymmetry. This framework fits nicely with the apparent unification of the gauge couplings in the minimal supersymmetric standard model. However, breaking supersymmetry at low energies is a hard problem, which in string perturbation theory implies a large extra dimension. Recently, an alternative approach has been put forward in which stabilization of the hierarchy is achieved without supersymmetry, by lowering the string scale down to a few TeV. A natural realization of this possibility is offered by weakly coupled type I string theory, where gauge interactions are described by open strings whose ends are confined on D-branes, while gravity is mediated by closed strings in the bulk. The observed hierarchy between the Planck and the weak scales is then accounted for by two or more large dimensions, transverse to our brane-world, with corresponding size varied from a millimeter to a fermi.

One of the main questions with such a low string scale is to understand the observed values of the low energy gauge couplings. One possibility is to have the three gauge group factors of the Standard Model arising from different collections of coinciding branes. This is unattractive since the three gauge couplings correspond in this case to different arbitrary parameters of the model. A second possibility is to maintain unification by imposing all the Standard Model gauge bosons to arise from the same collection of D-branes. The large difference in the actual values of gauge couplings could then be explained either by introducing power-law running from a few TeV to the weak scale, or by an effective logarithmic evolution in the transverse space in the special case of two large dimensions. However, no satisfactory model built along these lines has so far been presented.

In this work, we propose a third possibility which is alternative to unification but nevertheless maintains the prediction of the weak angle at low energies. Specifically, we consider the strong and electroweak interactions to arise from two different collections of coinciding branes, leading to two different gauge couplings. Assuming that the low energy spectrum of the (non-supersymmetric) Standard Model can be derived by a type I/I' string vacuum, the normalization of the hypercharge is determined in terms of the two gauge couplings and leads naturally to the right value of $\sin^2 \theta_W$ for a string scale of the order of a few TeV. The electroweak gauge symmetry is broken by the vacuum expectation values of two Higgs doublets, which are both necessary in the present context to give masses to all quarks and leptons.
Another issue of this class of models with TeV string scale is to understand proton stability. In the model presented here, this is achieved by the conservation of the baryon number which turns out to be a perturbatively exact global symmetry, remnant of an anomalous $U(1)$ gauge symmetry broken by the Green-Schwarz mechanism. Specifically, the anomaly is canceled by shifting a corresponding axion field that gives mass to the $U(1)$ gauge boson.

The model and the weak angle

The gauge group closest to the $SU(3) \times SU(2) \times U(1)$ of the Standard Model one can hope to derive from type I/I' string theory in the above context is $U(3) \times U(2) \times U(1)$. The first factor arises from three coincident D-branes ("color" branes). An open string with one end on them is a triplet under $SU(3)$ and carries the same $U(1)$ charge for all three components. Thus, the $U(1)$ factor of $U(3)$ has to be identified with gauged baryon number. Similarly, $U(2)$ arises from two coincident "weak" D-branes and the corresponding abelian factor is identified with gauged weak-doublet number. A priori, one might expect that $U(3) \times U(2)$ would be the minimal choice. However, this is not good enough because the hypercharge cannot be expressed as a linear combination of baryon and weak-doublet numbers. Therefore, at least one additional $U(1)$ factor corresponding to an extra D-brane ("$U(1)$" brane) is necessary in order to accommodate the Standard Model. In principle this $U(1)$ brane can be chosen to be independent of the other two collections with its own gauge coupling. To improve the predictability of the model, here we choose to put it on top of either the color or the weak D-branes. In either case, the model has two independent gauge couplings $g_3$ and $g_2$ corresponding, respectively, to the gauge groups $U(3)$ and $U(2)$. The $U(1)$ gauge coupling $g_1$ is equal to either $g_3$ or $g_2$.

Let us denote by $Q_3$, $Q_2$ and $Q_1$ the three $U(1)$ charges of $U(3) \times U(2) \times U(1)$, in a self explanatory notation. Under $SU(3) \times SU(2) \times U(1)_3 \times U(1)_2 \times U(1)_1$, the members of a family of quarks and leptons have the following quantum numbers:

\[
\begin{align*}
Q & \quad (3, 2; 1, w, 0)_{1/6} \\
\bar{u}^c & \quad (\bar{3}, 1; -1, 0, x)_{2/3} \\
\bar{d}^c & \quad (\bar{3}, 1; -1, 0, y)_{1/3}
\end{align*}
\]

See nevertheless the comments at the end of this section for a string embedding of the Standard Model based on $U(3) \times U(2)$, where the two $U(1)$’s are not the baryon and weak-doublet numbers. The model is though unsatisfactory for phenomenological reasons.
Here, we normalize all $U(N)$ generators according to $\text{Tr}T^aT^b = \delta^{ab}/2$, and measure the corresponding $U(1)_N$ charges with respect to the coupling $g_N/\sqrt{2N}$, with $g_N$ the $SU(N)$ coupling constant. Thus, the fundamental representation of $SU(N)$ has $U(1)_N$ charge unity. The values of the $U(1)$ charges $x, y, z, w$ will be fixed below so that they lead to the right hypercharges, shown for completeness as subscripts.

The quark doublet $Q$ corresponds necessarily to a massless excitation of an open string with its two ends on the two different collections of branes. The $Q_2$ charge $w$ can be either $+1$ or $-1$ depending on whether $Q$ transforms as a $2$ or a $\bar{2}$ under $U(2)$. The antiquark $\bar{u}^c$ corresponds to fluctuations of an open string with one end on the color branes and the other on the $U(1)$ brane for $x = \pm 1$, or on other branes in the bulk for $x = 0$. Ditto for $d^c$. Similarly, the lepton doublet $L$ arises from an open string with one end on the weak branes and the other on the $U(1)$ brane for $z = \pm 1$, or in the bulk for $z = 0$. Finally, $l^c$ corresponds necessarily to an open string with one end on the $U(1)$ brane and the other in the bulk. We defined its $Q_1 = 1$.

The weak hypercharge $Y$ is a linear combination of the three $U(1)$’s:

$$Y = c_1 Q_1 + c_2 Q_2 + c_3 Q_3.$$  

$c_1 = 1$ is fixed by the charges of $l^c$ in eq. (1), while for the remaining two coefficients and the unknown charges $x, y, z, w$, we obtain four possibilities:

\begin{align*}
  c_2 &= \frac{1}{2}, \quad c_3 = \frac{1}{3}; \quad x = -1, \quad y = 0, \quad z = 0, \quad w = -1 \\
  c_2 &= \frac{1}{2}, \quad c_3 = \frac{1}{3}; \quad x = -1, \quad y = 0, \quad z = -1, \quad w = 1 \\
  c_2 &= -\frac{1}{2}, \quad c_3 = \frac{2}{3}; \quad x = 0, \quad y = 1, \quad z = 0, \quad w = 1 \\
  c_2 &= \frac{1}{2}, \quad c_3 = \frac{2}{3}; \quad x = 0, \quad y = 1, \quad z = -1, \quad w = -1
\end{align*}

Orientifold models realizing the $c_3 = -1/3$ embedding in the supersymmetric case with intermediate string scale $M_s \sim 10^{11}$ GeV have been described in [12].

$^3$A study of hypercharge embeddings in gauge groups obtained from M-branes was considered in Ref. [11]. In the context of Type I groundstates such embeddings were considered in [12].
To compute the weak angle $\sin^2 \theta_W$, we use from eq. (2) that the hypercharge coupling $g_Y$ is given by

$$\frac{1}{g_Y^2} = \frac{2}{g_1^2} + \frac{4c_2^2}{g_2^2} + \frac{6c_3^2}{g_3^2}, \quad (4)$$

with $g_1 = g_2$ or $g_1 = g_3$ at the string scale. On the other hand, with the generator normalizations employed above, the weak $SU(2)$ gauge coupling is $g_2$. Thus,

$$\sin^2 \theta_W \equiv \frac{g_Y^2}{g_2^2 + g_Y^2} = \frac{1}{1 + 4c_2^2 + 2g_2^2/g_1^2 + 6c_3^2g_2^2/g_3^2}, \quad (5)$$

which for $g_1 = g_2$ reduces to:

$$\sin^2 \theta_W(M_s) = \frac{1}{4 + 6c_3^2g_2^2(M_s)/g_3^2(M_s)}, \quad (6)$$

while for $g_1 = g_3$ it becomes:

$$\sin^2 \theta_W(M_s) = \frac{1}{2 + 2(1 + 3c_3^2)g_2^2(M_s)/g_3^2(M_s)}. \quad (7)$$

We now show that the above predictions agree with the experimental value for $\sin^2 \theta_W$ for a string scale in the region of a few TeV. For this comparison, we use the evolution of gauge couplings from the weak scale $M_Z$ as determined by the one-loop beta-functions of the Standard Model with three families of quarks and leptons and one Higgs doublet,

$$\frac{1}{\alpha_i(M_s)} = \frac{1}{\alpha_i(M_Z)} - \frac{b_i}{2\pi} \ln \frac{M_s}{M_Z}; \quad i = 3, 2, Y \quad (8)$$

where $\alpha_i = g_i^2/4\pi$ and $b_3 = -7$, $b_2 = -19/6$, $b_Y = 41/6$. We also use the measured values of the couplings at the $Z$ pole $\alpha_3(M_Z) = 0.118 \pm 0.003$, $\alpha_2(M_Z) = 0.0338$, $\alpha_Y(M_Z) = 0.01014$ (with the errors in $\alpha_{2,Y}$ less than 1%) [13].

In order to compare the theoretical relations for the two cases (6) and (7) with the experimental value of $\sin^2 \theta_W = g_Y^2/(g_2^2 + g_Y^2)$ at $M_s$, we plot in Fig. 1 the corresponding curves as functions of $M_s$. The solid line is the experimental curve. The dashed line is the plot of the function (6) for $c_3 = -1/3$ while the dotted-dashed line corresponds to the function (7) for $c_3 = 2/3$. Thus, the second case, where the $U(1)$ brane is on top of the color branes, is compatible with low energy data for $M_s \sim 6 - 8$ TeV. This selects the last two possibilities of charge assignments in Eq. (3). Indeed, the curve corresponding to $g_1 = g_3$

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4The gauge couplings $g_{2,3}$ are determined at the tree-level by the string coupling and other moduli, like radii of longitudinal dimensions. In higher orders, they also receive string threshold corrections.
Figure 1: The experimental value of $\sin^2 \theta_W$ (thick curve), together with the theoretical predictions (6) with $c_3 = -1/3$ (dashed line) and (7) with $c_3 = 2/3$ (dotted-dashed), are plotted as functions of the string scale $M_s$.

and $c_3 = -1/3$ intersects the experimental curve for $\sin^2 \theta_W$ at a scale $M_s$ of the order of a few thousand TeV. Since this value appears to be too high to protect the hierarchy, it is less interesting and is not shown in Fig. 1. The other case, where the $U(1)$ brane is on top of the weak branes, is not interesting either. For $c_3 = 2/3$, the corresponding curve does not intersect the experimental one at all and is not shown in the Fig. 1, while the case of $c_3 = -1/3$ leads to $M_s$ of a few hundred GeV and is most likely excluded experimentally. In the sequel we shall restrict ourselves to the last two possibilities of Eq. (3).

From the general solution (3) and the requirement that the Higgs doublet has hypercharge 1/2, one finds the following possible assignments for it, in the notation of Eq. (1):

\[
c_2 = -\frac{1}{2}: \quad H \quad (1, 2; 0, 1, 1)_{1/2} \quad H' \quad (1, 2; 0, -1, 0)_{1/2} \quad (9)
\]

\[
c_2 = \frac{1}{2}: \quad \tilde{H} \quad (1, 2; 0, -1, 1)_{1/2} \quad \tilde{H}' \quad (1, 2; 0, 1, 0)_{1/2} \quad (10)
\]
It is straightforward to check that the allowed (trilinear) Yukawa terms are:

\[ c_2^2 = -\frac{1}{2} : \quad H' Q u^c , \quad H L l^c , \quad H Q d^c \]  
\[ c_2^2 = \frac{1}{2} : \quad \tilde{H}' Q u^c , \quad \tilde{H} \tilde{L} \tilde{l}^c , \quad \tilde{H} \tilde{Q} \tilde{d}^c \]

Thus, two Higgs doublets are in each case necessary and sufficient to give masses to all quarks and leptons. Let us point out that the presence of the second Higgs doublet changes very little the curves of Fig. 1 and consequently our previous conclusions about \( M_s \) and \( \sin^2 \theta_W \).

A few important comments are now in order:

(i) The spectrum we assumed in Eq. (1) does not contain right-handed neutrinos on the branes. They could in principle arise from open strings in the bulk. Their interactions with the particles on the branes would then be suppressed by the large volume of the transverse space [14]. More specifically, conservation of the three U(1) charges allow for the following Yukawa couplings involving the right-handed neutrino \( \nu_R^c \):

\[ c_2 = -\frac{1}{2} : \quad H' L \nu_L \quad ; \quad c_2 = \frac{1}{2} : \quad \tilde{H} L \nu_R \]

These couplings lead to Dirac type neutrino masses between \( \nu_L \) from \( L \) and the zero mode of \( \nu_R^c \), which is naturally suppressed by the volume of the bulk.

(ii) Implicit in the above was our assumption of three generations (1) of quarks and lepton in the light spectrum. They can arise, for example, from an orbifold action along the lines of the model described in Ref. [12].

(iii) From Eq. (7) and Fig. 1, we find the ratio of the \( SU(2) \) and \( SU(3) \) gauge couplings at the string scale to be \( \alpha_2/\alpha_3 \approx 0.4 \). This ratio can be arranged by an appropriate choice of the relevant moduli. For instance, one may choose the weak branes to extend along one extra dimension transverse to the color branes, with size around twice the string length. Another possibility would be to move slightly off the orientifold point which may be necessary also for other reasons (see discussion towards the end of the paper).

(iv) Finally, it should be stressed that the charge assignments (1) were based on the assumption that the anti-quarks \( u^c \) and \( d^c \) arise as excitations of open strings with only one end on the color D-branes. This is not however the only possibility. The fact that the \( \bar{3} \) of \( SU(3) \) can also be obtained as the antisymmetric product of two \( 3 \)'s implies that \( u^c \) and \( d^c \) may also arise as fluctuations of open strings with both ends on the color branes. Similarly, the anti-lepton \( l^c \) which is \( SU(2) \) singlet can be obtained as the antisymmetric product of two doublets and consequently it may arise as a fluctuation of an open string with both ends on
the weak branes. In these cases, the quantum numbers of the corresponding particles will be:

\[ u^c : (3, 1; 2, 0, 0)_{-2/3} \quad d^c : (3, 1; 2, 0, 0)_{1/3} \quad l^c : (1, 1; 0, \mp 2, 0)_1 \]  

(14)

One should then repeat the previous analysis from the beginning, with any combination of the particles \( u^c, d^c \) and \( l^c \) in Eq. (11) replaced by the corresponding \( u^c, d^c \) and \( l^c \). However, as we argue next, the only physically viable alternative scenario to the one discussed above is just to replace \( l^c \) by \( l^c \).

First observe that \( d^c \) cannot be replaced by \( d^c \). Indeed, this would fix \( c_3 = 1/6 \) and the \( Q \) hypercharge would set \( c_2 = 0 \). It is then easy to see that one cannot satisfy the hypercharge assignments of leptons for either choice of \( l^c \) or \( l^c \). Next, let us replace \( u^c \) by \( u^c \). This fixes \( c_3 = -1/3 \). If we keep \( l^c \), then \( c_1 = 1 \) and one is left with the first two lines of Eq. (3) as the two possible solutions for \( y, z, w \) (\( x \) is absent in this case). From our previous analysis of \( \sin^2 \theta_W \), these solutions are not very interesting since the string scale comes out to be either too low or too high. On the other hand, if we substitute also \( l^c \) by \( l^c \), the solutions for the remaining parameters are: (a) the second line of Eq. (3) with \( c_1 = 1 \) as before, which is uninteresting, and (b) the first line of Eq. (3) with \( c_1 \) undetermined. In this case, the \( Q_1 \) charges of all particles vanish, the corresponding gauge field decouples, and the gauge group becomes effectively \( U(3) \times U(2) \) with \( Y = -Q_2/2 - Q_3/3 \). At first sight, this seems to be a more economical embedding of the Standard Model than the one based on \( U(3) \times U(2) \times U(1) \). In this case, the \( g_1 \) term drops from Eq. (11) and the weak angle is given by \( 1/\sin^2 \theta_W(M_s) = 2 + 2g_2^2(M_s)/3g_3^2(M_s) \). Unfortunately, comparison with the experimental value of \( \sin^2 \theta_W \) at \( M_Z \) requires a string scale of the order of \( 10^{14} \) GeV. An additional unsatisfactory feature of the models obtained by replacing \( u^c \) with \( u^c \) is the absence of appropriate Yukawa couplings to give masses to the up-quarks.

The last case to be examined is the substitution of \( l^c \) alone by \( l^c \). This leads to the same four solutions (3) as with \( l^c \), and thus, to the same conclusions for \( \sin^2 \theta_W \) and \( M_s \). However, the case with \( c_2 = 1/2 \) is problematic because the charge assignments do not allow tree-level Yukawa interactions to give masses to the leptons. In the case with \( c_2 = -1/2 \), the Yukawa couplings of the leptons (12) are slightly modified to

\[ c_2 = -\frac{1}{2} : \quad H^\dagger L l^c, \]  

(15)

implying that they acquire masses from the Higgs which gives masses also to the up-quarks.
Extra $U(1)$’s, anomalies and proton stability

The model under discussion has three $U(1)$ gauge interactions corresponding to the generators $Q_1, Q_2, Q_3$. From the previous analysis, the hypercharge was shown to be either one of the two linear combinations:

$$Y = Q_1 \pm \frac{1}{2} Q_2 + \frac{2}{3} Q_3.$$ 

(16)

It is easy to see that the remaining two $U(1)$ combinations orthogonal to $Y$ are anomalous. Indeed, the generic two-parameter generator orthogonal to $Y$ is

$$\tilde{Q} = \left( \pm \frac{\beta}{2} - \frac{2\gamma}{3} \right) Q_1 + \beta Q_2 + \gamma Q_3,$$

(17)

which satisfies $\text{Tr} \tilde{Q} T_{SU(2)}^2 = \pm 5\beta/4 - \gamma/3$ and $\text{Tr} \tilde{Q} T_{SU(3)}^2 = 2\beta + 3\gamma/2$ (for $c_2 = -1/2$), or $-3\beta/4 + 11\gamma/6$ (for $c_2 = 1/2$); they can both vanish only for $\beta = \gamma = 0$.

We have assumed throughout that this model can be derived as a consistent type I/I$'$ string vacuum without additional light states charged under $U(3) \times U(2) \times U(1)$. In such a vacuum, the anomalies should be canceled by appropriate shifts of Ramond-Ramond axions in the bulk \[13\]. Since we have two independent anomalous $U(1)$ currents, we need two axions $a, a'$ that couple to the non-abelian Pontryagin densities with coefficients fixed by the anomalies. The relevant part of the low-energy effective lagrangian can be written as:

$$L_{eff}^{(1)} = \frac{1}{2} (\partial a + \lambda M_s A)^2 + \frac{1}{32\pi^2} \frac{a}{\lambda M_s} (k_2 \text{Tr} F_2 \wedge F_2 + k_3 \text{Tr} F_3 \wedge F_3)$$

$$+ \frac{1}{2} (\partial a' + \lambda M_s A')^2 + \frac{1}{32\pi^2} \frac{a'}{\lambda M_s} (k'_2 \text{Tr} F_2 \wedge F_2 + k'_3 \text{Tr} F_3 \wedge F_3),$$

(18)

where $F_2$ and $F_3$ are the $SU(2)$ and $SU(3)$ gauge field strengths, and $A, A'$ the gauge fields corresponding to two independent and mutually orthogonal anomalous abelian charges $Q_A, Q_{A'}$ of the form \[17\]. $k_2, k'_2, k_3, k'_3$ are their respective mixed anomalies with $SU(2)$ ($SU(3)$) given by

$$k_2 = \text{Tr} Q_A T_{SU(2)}^2, \quad k'_2 = \text{Tr} Q_{A'} T_{SU(2)}^2, \quad k_3 = \text{Tr} Q_A T_{SU(3)}^2, \quad k'_3 = \text{Tr} Q_{A'} T_{SU(3)}^2,$$

(19)

while $\lambda$ is a calculable parameter in every particular string vacuum. The theory is invariant under the gauge transformation $A \rightarrow A + \partial \Lambda / g_A, a \rightarrow a - \lambda M_s \Lambda$, together with appropriate transformations of the fermion fields. Indeed, under this transformation the action \[18\] changes by exactly the amount necessary to cancel the phase of the fermionic determinant. Ditto for $A'$. According to Eq. \[18\], the gauge fields $A$ and $A'$ become massive with masses
\[ \lambda g_A M_s \] and \[ \lambda g_{A'} M_s \], respectively, with \( g_A \) and \( g_{A'} \) the corresponding gauge couplings. The axions \( a \) and \( a' \) become their longitudinal components. Note that we have chosen \( A \) and \( A' \) so that their mass matrix is diagonal. Gravitational anomalies proportional to the trace of a single charge are also canceled by similar axionic couplings to \( R \wedge R \).

This mechanism can be generalized to show the cancellation of the mixed \( U(1) \) anomalies. These are of two types. First, the ones associated to the non-vanishing traces \( \text{Tr} Q_A Y^2 \equiv k_Y \) and \( \text{Tr} Q_{A'} Y^2 \equiv k'_Y \) can be canceled by introducing in \( \mathcal{L}_{\text{eff}}^{(1)} \) the additional terms

\[
\mathcal{L}_{\text{eff}}^{(1)} \rightarrow \mathcal{L}_{\text{eff}}^{(1)} + \frac{1}{32\pi^2 \lambda M_s} (k_Y a + k'_Y a') F_Y \wedge F_Y ,
\]

which are needed to cancel the \( F_Y \wedge F_Y \) contribution to the divergence of the two anomalous currents. The coefficients \( k_Y \) and \( k'_Y \) can be deduced from the anomaly of the generic current \( \mathcal{L}_{\text{eff}}^{(1)} \). In the case of \( l^c \) we obtain \( \text{Tr} \tilde{Q} Y^2 = 4 \beta/3 - 43 \gamma/18 \) (for \( c_2 = -1/2 \)), or \(- \beta/12 - 37 \gamma/18 \) (for \( c_2 = 1/2 \)), while for \( l^e \) (and \( c_2 \) necessarily \(-1/2 \)) \( \text{Tr} \tilde{Q} Y^2 = -7 \beta/6 - 31 \gamma/18 \).

Second, there are mixed anomalies related to the non-vanishing trace \( \text{Tr} \tilde{Q} \tilde{Y}^2 = \beta^2/2 + 20 \gamma^2/9 - \beta \gamma/3 \) (for \( c_2 = -1/2 \)), or \(-3 \beta^2/4 + 16 \gamma^2/9 - 5 \beta \gamma/3 \) (for \( c_2 = 1/2 \)), or \( 17 \beta^2/4 + 16 \gamma^2/9 + \beta \gamma/3 \) (in the case of \( l^e \) for \( c_2 = -1/2 \)). Using this general formula, we can uniquely determine the two orthogonal combinations \( Q_A \) and \( Q_{A'} \) in such a way that the triple mixed trace vanishes. We thus have:

\[
\text{Tr} Y Q_A^2 \equiv \xi_A , \quad \text{Tr} Y Q_{A'}^2 \equiv \xi_{A'} , \quad \text{Tr} Y Q_A Q_{A'} = 0 .
\]

These mixed anomalies seem to violate the hypercharge gauge invariance of the Standard Model. However, in the context of a consistent string theory, they should also be eliminated. This can be achieved without giving mass to the hypercharge gauge field \( A_Y \) if the low-energy effective lagrangian contains Chern-Simons terms of the form \( A_Y \wedge \omega_A \) and \( A_Y \wedge \omega_{A'} \). Finally, the violation of the \( U(1)_A \) and \( U(1)_{A'} \) gauge invariances introduced by these new terms can be remedied by adding non-diagonal axionic couplings proportional to \( a F_Y \wedge F_A \) and \( a' F_Y \wedge F_{A'} \). To summarize, one may cancel all anomalies of our model by modifying the relevant to the anomaly part of the effective lagrangian \( \mathcal{L}_{\text{eff}}^{(1)} \) in Eq. (18) to:

\[
\mathcal{L}_{\text{eff}}^{\text{anom}} = \mathcal{L}_{\text{eff}}^{(1)} + \frac{1}{32\pi^2 \lambda M_s} (k_Y a + k'_Y a') F_Y \wedge F_Y - \frac{1}{32\pi^2} A_Y \wedge (\xi_A \omega_A + \xi_{A'} \omega_{A'}) + \frac{1}{32\pi^2 \lambda M_s} F_Y \wedge (\xi_A a F_A + \xi_{A'} a' F_{A'}) .
\]

For completeness, we give the linear combinations \( Q_A \) and \( Q_{A'} \) that satisfy Eq. (21):

\[
Q_A \sim \frac{3}{2} Q_1 \pm \frac{13}{3} Q_2 + Q_3 + t \left( -\frac{2}{3} Q_1 + Q_3 \right)
\]
\[ Q_{A'} \sim -t \left( \frac{3}{2}Q_1 \pm \frac{13}{3}Q_2 + Q_3 \right) + \frac{61}{4} \left( -\frac{2}{3}Q_1 + Q_3 \right) \]  

(23)

where the \( \pm \) signs correspond to \( c_2 = \mp 1/2 \), respectively. For \( c_2 = 1/2 \) the value of \( t \) is

\[ t = \frac{1159 \pm 13\sqrt{21533}}{388} \]  

For \( c_2 = -1/2, t = \frac{427 \pm 13\sqrt{1342}}{54} \) for \( l^c \), and \( t = \frac{-671 \pm 91\sqrt{61}}{60} \) for \( l'^c \).

An important property of the above Green-Schwarz anomaly cancellation mechanism is that the two \( U(1) \) gauge bosons \( A \) and \( A' \) acquire masses leaving behind the corresponding global symmetries (23) [15]. This is in contrast to what would had happened in the case of an ordinary Higgs mechanism. These global symmetries remain exact to all orders in type I string perturbation theory around the orientifold vacuum. This follows from the topological nature of Chan-Paton charges in all string amplitudes. On the other hand, one expects non-perturbative violation of global symmetries and consequently exponentially small in the string coupling, as long as the vacuum stays at the orientifold point. Once we move sufficiently far away from it, we expect the violation to become of order unity. This can be justified in a supersymmetric theory as follows. Every Ramond-Ramond axion \( a \) is part of a chiral superfield \( a + im/g_s \) with \( g_s \) the string coupling. Its scalar component \( m \) is a NS-NS (Neveu-Schwarz) closed string modulus, whose vacuum expectation value (VEV) blows up the orbifold singularities moving away from the orientifold point. Using the shift of the axion under gauge transformations, one can form the complex field \( e^{(ia−m/g_s)/M_s} \) that transforms covariantly with charge \(-\lambda\). A matter interaction term with charge \( n\lambda \) (with integer \( n \)), multiplied by the \( n\)-th power of this field forms a neutral operator which can appear in the effective action. For \( < m > \neq 0 \), one thus obtains charge violating non-perturbative interaction terms with strength \( \mathcal{O}(e^{−<m>/g_sM_s}) \). A small \( < m > \) of order \( g_sM_s \) leads therefore to charge violations of order unity.

So, as long as we stay at the orientifold point, all three charges \( Q_1, Q_2, Q_3 \) are conserved and since \( Q_3 \) is the baryon number, proton stability is guaranteed.

To break the electroweak symmetry, the Higgs doublets in Eq. (9) or (10) should acquire non-zero VEV’s. Since the model is non-supersymmetric, this may be achieved radiatively [14]. From Eqs. (11) and (12), to generate masses for all quarks and leptons, it is necessary for both Higgses to get non-zero VEV’s. The baryon number conservation remains intact because both Higgses have vanishing \( Q_3 \). However, the linear combination \((t − 61/4)Q_A + (t + 1)Q_{A'}\) which does not contain \( Q_3 \), will be broken spontaneously, as follows from their quantum numbers in Eqs. (9) and (10). This leads to an unwanted massless Goldstone boson of the Peccei-Quinn type. A possible way out is to break this global symmetry explicitly, by moving
away from the orientifold point along the direction of non-vanishing $(t - 61/4)m + (t + 1)m'$, so that baryon number remains conserved.

In conclusion, we presented a particular embedding of the Standard Model in a non-supersymmetric D-brane configuration of type I/I' string theory. The strong and electroweak couplings are not unified because strong and weak interactions live on different branes. Nevertheless, $\sin^2 \theta_W$ is naturally predicted to have the right value for a string scale of the order of a few TeV. The model contains two Higgs doublets needed to give masses to all quarks and leptons, and preserves baryon number as a (perturbatively) exact global symmetry. The model satisfies the main phenomenological requirements for a viable low energy theory and its explicit derivation from string theory deserves further study.

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