# ON THE D-BRANES STANDARD-LIKE MODELS*

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Based on the low-energy effective field theory of D-branes, the mass spectrum of an extended Standard Model with two-Higgs doublets used to generate all the mass terms is investigated. Besides the gauge bosons, the fermion mass spectrum is weighted by the Higgs VEVs with a partial hierarchy and the smallness of neutrino masses is exhibited. With reference to the known data, the involved scales of the model are approached.

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## 1. Introduction

Up to date, all low-energy phenomena have been successfully predicted by the Standard Model of particle physics (SM) according to a vast amount of experimental data [1–4]. The Higgs interactions with the gauge bosons and fermions are completely determined by theory and are yet to be experimentally established [3, 4]. Despite its experimental success, the SM is not satisfactory as a fundamental theory, and there are several attempts with supplementary symmetries, matters or dimensions which have been developed to address its open issues with new phenomenological features [5–7]. The scalar-extended SMs constitute one of these most well-motivated extensions bringing many interesting phenomenological features such as fermion masses hierarchy and dark matter candidates [6–9].

All these ingredients naturally occur in some string theory models involving more U(1)s in the gauge group and leading to effective low-energy theories containing more scalar fields beyond the SM. A feature of such scalar-extended models is the existence of tree-level flavour changing neutral currents (FCNC) whose the potentially dangerous interactions can be avoided here by the presented extra symmetries [10, 11].

Intensive works with orientifold constructions has been recently developed based on intersecting D-branes embedded in superstrings theory [12–15], where $N$ coincident D-branes typically generate a unitary group $\mathrm{U}(N) \simeq \mathrm{SU}(N) \times \mathrm{U}(1)$ and, hence, every stack of branes supplies the model with an extra Abelian factor in the gauge group. Such U(1) fields have generically four-dimensional anomalies which are cancelled via the Green–Schwarz mechanism [16, 17]. This mechanism gives a mass to the anomalous U(1) fields and breaks the associated gauge symmetry [18, 19]. In this string vacua, SM particles are considered as open string states attached to different stacks of D-branes and their interactions are subject to additional restriction from these remnant global Abelian U(1) symmetries [20–24]. This approach provides an acceptable effective low-energy description reproducing SM-like or some extensions and offering a simple low $M_{\mathrm{s}}$ string scale energy framework.

In this paper, motivated by the progressive efforts towards obtaining the SM in string theory constructions, we present a quiver gauge theory of a Higgs-extended SM below a low-mass scale $M_{\mathrm{s}}$ based on the low-energy effective field theory of intersecting D-branes giving rise to the gauge symmetry $\mathrm{U}(3) \times \mathrm{SU}(2) \times \mathrm{U}(1)^n$. The four-stack configuration, $n = 2$, contains two Higgs doublets $H_u$ and $H_d$ to generate all fermion Yukawa couplings. Besides the gauge bosons getting their masses from the two Higgs VEVs $v_u$, $v_d$, the resulting fermion masses are such that the heavy fermions are weighted by the Higgs VEVs $v_u$, while the light fermions including neutrinos are weighted by the second Higgs vev $v_d$ giving a partial solution to fermion masses hierarchy as well as the smallness of neutrino masses. According to the known experimental data, we bound the Higgs VEVs $v_u$, $v_d$ and the mass scale $M_{\mathrm{s}}$.

## 2. D-brane inspiration

### 2.1. D-branes and fields content

String theory realizations of important particle physics models have recently grown much interest. Growing studies in gauge theories have been developed based on intersecting D-branes embedded in Type II superstrings where it is possible to investigate deeper details seen in particle physics [13–15]. In these compactifications, the gauge groups

$$\mathrm{G}_{2+n} = \mathrm{U}(3) \times \mathrm{U}(2) \times \mathrm{U}(1)^n \,, \qquad n \geqslant 0 \tag{1}$$

arise from stacks of D-branes that fill out four-dimensional spacetime and wrap three-cycles in the internal Calabi–Yau threefold. Chiral matter arises at the intersection of two different D-brane stacks in the internal space, and their multiplicity is given by the topological intersection number of the respective three-cycles. Their interactions are subject to additional restrictions of the global U(1)s exhibited by orientifold compactifications. The vacua we consider here are obtained from four stacks of intersecting D-branes giving rise to the symmetry group[1]

$$\mathrm{G}_4 = \mathrm{U}(3)_a \times \mathrm{Sp}(1)_b \times \mathrm{U}(1)_c \times \mathrm{U}(1)_d \,. \tag{2}$$

Given this brane content and conditions that we require of the SM realizations such as tadpole cancellation and the presence of a massless hypercharge [16–19], we can build an intersecting brane model with the chiral content of the SM. Since the Sp(1) does not exhibit a $\mathrm{U}(1)_b$ which could contribute to the hypercharge, the mixed anomalies and anomalous Abelian $\mathrm{U}(1)_{a,c,d}$ parts are cancelled by the Green–Schwarz mechanism and promoted to global $\mathrm{U}(1)_{a,c,d}$ symmetries which are respected by all perturbative couplings and a linear combination $\mathrm{U}(1)_Y = q_a\mathrm{U}(1)_a \times q_c\mathrm{U}(1)_c \times q_d\mathrm{U}(1)_d$ of them does not acquire a Stuckelberg mass and remains massless to be identified as the hypercharge. The tadpole cancellations, which are conditions on the cycles the D-branes wrap, imply restrictions on the transformation properties of the chiral spectrum and guarantee the cancellation of gauge anomalies $\mathrm{U}(N_\alpha)$. Here, it is seen that the SU(2) is realized as Sp(1), all representations are real and the tadpole equations do not impose any condition on the transformation behavior under the Sp(1). Moreover, the fact that the stack $b = b^*$ limits the potential origins of fields charged under the $b$-brane. These conditions are used to fit the $\mathrm{U}(1)_{a,c,d}$ charged SM particles

$$\begin{aligned}
Q^i &= (u,d)^i_{\mathrm{L}} = (u,d)_{\mathrm{L}}\,,\ (c,s)_{\mathrm{L}}\,,\ (t,b)_{\mathrm{L}}\ , \\
\overline{u}^{\,i} &= u^{c^i}_{\mathrm{R}} = \overline{u}, \overline{c}, \overline{t}\,, \\
\overline{d}^{\,i} &= d^{c^i}_{\mathrm{R}} = \overline{d}, \overline{s}, \overline{b}\,, \\
L^i &= (v_e,e)^i_{\mathrm{L}} = (v_e,e)\,,\ (v_\mu,\mu)\,,\ (v_\tau,\tau)\ , \\
\overline{e}^{\,i} &= e^{c^i}_{\mathrm{R}} = \overline{e}, \overline{\mu}, \overline{\tau}\,.
\end{aligned} \tag{3}$$

In our description, the three left-handed quarks $Q^i$ are localized at the intersections of D6-branes $a$ and $b$, while right-handed quarks $\overline{u}^{\,i}$ and $\overline{d}^{\,i}$ split into two up quarks $\overline{u}^{\,2,3}$ and one down quark $\overline{d}^{\,3}$ at the intersection of the D6-branes $a$ and $c/c^*$, and two down quarks $\overline{d}^{\,1,2}$ and one up quark $\overline{u}^{\,1}$ at the intersection of the D6-branes $a$ and $d/d^*$. The three left-handed leptons $L^i$

[1] Where the $\mathrm{Sp}(1) \simeq \mathrm{SU}(2)$ weak symmetry arise from D6-wrapped on an orientifold invariant three-cycle ($b = b^*$) and $\mathrm{U}(1)_d$ is a gauged flavor symmetry distinguishing various quarks from each others.

arise at the intersection of branes $b$ and $d$ respectively, and the three right-handed electrons $\overline{e}^{\,i}$ arise at the intersection of D6-branes $d$ and $c^*$. Finally, the two Higgs doublets $H_c$ and $H_d$ arise at intersection of D6-branes $b$ and $c/d^*$ respectively. These correspond to the following fermion intersection numbers[2]

$$\begin{aligned} I_{ab} &= 3\,, \qquad I_{ac} = -2\,, \qquad I_{ac^*} = -1\,,\\ I_{ad} &= -1\,, \qquad I_{ad^*} = -2\,, \qquad I_{db} = 3\,,\\ I_{dc^*} &= -3 \end{aligned} \tag{4}$$

which is indeed obeyed by the above spectrum. From these intersection numbers, we summarize in the following table the fields content and the corresponding charges which depend on the anomaly-free hypercharge linear combination for which all the matter particles have the proper electroweak hypercharge.

Since any realistic string vacua have to exhibit the phenomenologically desired terms in its low-energy physics, we need to fit the Higgs sector. An examination of the associated quantum numbers in Table I shows that the communication of the electroweak symmetry breaking to all fermion requires, in addition to the SM Higgs doublet, a second Higgs doublet. The latter carries the following U(1) charges, given in Table II, leading to one-hypercharges $Y = 1/2$, two Higgs doublets $H_c = (H_c^+, H_c^0)$ and $H_d = (H_d^+, H_d^0)$, with the subscripts $c$ and $d$ referring to the $c$-brane and $d$-brane under which the Higgses are charged.

TABLE I

The fields content corresponding to the anomaly-free hypercharge linear combination $Y = \frac{1}{6}Q_a - \frac{1}{2}Q_c - \frac{1}{2}Q_d$. The index $i(=1,2,3)$ denotes the family index.

| Fields | $Q^i$ | $\overline{u}^{\,1}$ | $\overline{d}^{\,1,2}$ | $\overline{u}^{\,2,3}$ | $\overline{d}^{\,3}$ | $L^i$ | $\overline{e}^{\,i}$ |
|---|---|---|---|---|---|---|---|
| $\mathrm{U}(1)_c$ | 0 | 0 | 0 | 1 | $-1$ | 0 | $-1$ |
| $\mathrm{U}(1)_d$ | 0 | 1 | $-1$ | 0 | 0 | 1 | $-1$ |
| $Y$ | 1/6 | $-2/3$ | 1/3 | $-2/3$ | 1/3 | $-1/2$ | 1 |

[2] We have not included those involving $b^* = b$. The other intersection numbers are set to zero and as we discussed, the cancellation of $\mathrm{U}(N_{a,c,d})$ anomalies is

$$\sum_{\beta=a,b,c,d} N_\beta \left(I_{\alpha\beta} + I_{\alpha\beta^*}\right) = 0\,.$$

TABLE II

The Higgs sector with their U(1) charges.

| Fields | $H_c$ | $H_d$ |
|---|---|---|
| $\mathrm{U}(1)_c$ | $-1$ | $0$ |
| $\mathrm{U}(1)_d$ | $0$ | $-1$ |
| $Y$ | $1/2$ | $1/2$ |

Based on these data and string theory description, this model can be obtained from an orientifold compactification with D-branes wrapping non-trivial cycles as represented in Fig. 1 where bold lines represent the relevant D-branes, which are not distinguished from their orientifold images, giving rise to the gauge symmetry of Eq. (2); solid and dashed thin lines indicate the chiral intersections of those branes and refer to matter fields and Higgs fields respectively; arrows directions indicate fundamental (antifundamental) representations of the U($N$) gauge groups.

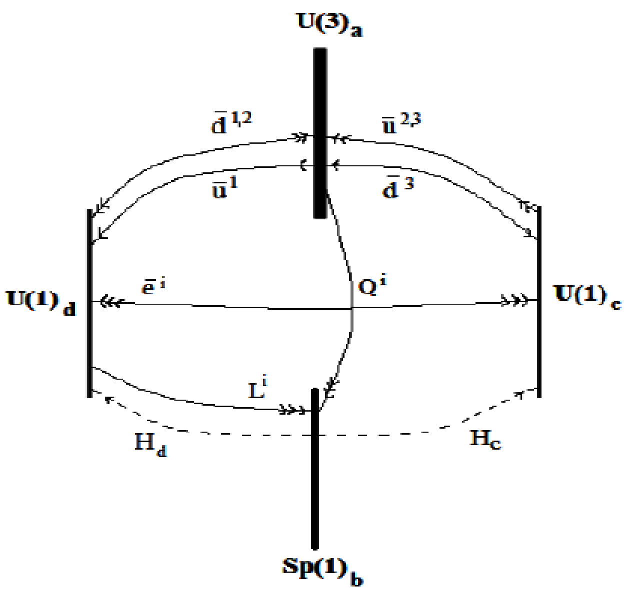


Fig. 1. Four-stack quiver SM.

Although this is a simple scalar extension of the SM with two one-hypercharge Higgses and thus with more parameters than in the minimal SM, it provides a rich mass spectrum.

## 3. Mass spectrum

### *3.1. Higgs spectrum and electroweak boson masses*

The Higgs scalar fields of the model consist of two complex isodoublets, or eight real, scalar degrees of freedom. We have then in terms of physical fields

$$H_{i=c,d} = \begin{pmatrix} H_i^+ \\ H_i^0 \end{pmatrix} = \begin{pmatrix} h_{1i}^+ + ih_{2i}^+ \\ (v_i + h_{1i}^0 + ih_{2i}^0) \end{pmatrix} . \tag{5}$$

When the electroweak symmetry is broken by the VEVs of the two Higgses $v_c$ and $v_d$ minimizing the corresponding renormalizable and gauge invariant Higgs potential

$$\begin{aligned} V(H_c, H_d) &= \mu_c^2 \left(H_c^\dagger H_c\right) + \mu_d^2 \left(H_d^\dagger H_d\right) + \tfrac{1}{2}\lambda_c \left(H_c^\dagger H_c\right)^2 + \tfrac{1}{2}\lambda_d \left(H_d^\dagger H_d\right)^2 \\ &\quad +\lambda_{cd} \left(H_c^\dagger H_c\right)\left(H_d^\dagger H_d\right) + \lambda_{cd}' \left(H_c^\dagger H_d\right)\left(H_d^\dagger H_c\right) , \end{aligned} \tag{6}$$

three of the eight real scalar degrees are the would-be charged and neutral Nambu–Goldstone bosons $G^\pm$ (mixture of $h_{1i}^+$ and $h_{2i}^+$) and $G^0$ (mixture of $h_{2i}^0$) which become the longitudinal modes of the $Z_\mu = g_2 W_\mu^3 - gB_\mu/\sqrt{g_2^2 + g_1^1}$ and $W_\mu^\pm = W_\mu^1 \mp iW_\mu^2/\sqrt{2}$ electroweak bosons getting masses through the kinetic part of the Higgs potential

$$\begin{aligned} |D_\mu H_c|^2 + |D_\mu H_d|^2 &= \left|\left(\partial_\mu + ig_2\frac{\tau_a}{2}W_\mu^a - ig_1\frac{1}{2}B_\mu\right)\right|^2 \left(|H_c|^2 + |H_d|^2\right) \\ &= \frac{1}{4}g_2^2\left(v_c^2 + v_d^2\right)^2 W_\mu^+ W^{-\mu} + \frac{1}{8}\left(g_2^2 + g_1^2\right)\left(v_c^2 + v_d^2\right)^2 Z_\mu Z^\mu + \dots \end{aligned} \tag{7}$$

With the SM constraint $\sqrt{v_c^2 + v_d^2} = v = \left(G_{\mathrm{F}}\sqrt{2}\right)^{-1/2} = 246$ GeV limiting the parameters of the Higgs potential, the electroweak boson masses are

$$m_W = \tfrac{1}{2}gv , \qquad m_Z = \tfrac{1}{2}\sqrt{g_2^2 + g_1^2}\,v . \tag{8}$$

The remaining five degrees of freedom are the physical Higgs bosons; they consist of three neutral scalars $h_1^0$, $h_1^{0\prime}$, $h_2^0$ (mixture of $h_{1i}^0$, $h_{2i}^0$ with $m_{h_2^0}$, $m_{h_1^{0\prime}} > m_{h_1^0}$) and two charged scalars $h^\pm$ (mixture of $h_{1i}^+$, $h_{2i}^+$ with $m_{h_1^{0\prime}} > m_{h_1^0}$) which are important from an experimental point of view.

### *3.2. Fermion masses and hierarchy*

In all 2HDMs, the most serious potential problem faced is the possibility of FCNC which can cause severe phenomenological difficulties [10, 11]. Nonetheless, under reasonable assumptions, models with these problems may still be viable. In the presented model, the tree-level FCNC are completely absent due to the extra continuous symmetries $\mathrm{U}(1)_{c,d}$ distinguishing the SM fermions such that all fermions with the same quantum numbers couple to the same Higgs multiplet [10, 11]. Since the presence of the two Higgs doublets was required to generate the Yukawa couplings for the all fermions, according to the D-brane model building selection rules, the allowed quark and lepton Yukawas are

$$
\begin{aligned}
&Q_{c,d} = \left(H_c Q^i \overline{u}^{2,3}\right) = Q_{c,d}\left(H_c^\dagger Q^i \overline{d}^{3}\right) = Q_{c,d}\left(H_d Q^i \overline{u}^{1}\right) = Q_{c,d}\left(H_d^\dagger Q^i \overline{d}^{1,2}\right) = 0\,,\\
&Q_{c,d} = M_\mathrm{s}^{-1}\left(H_d L^i\right)^2 = Q_{c,d}\left(H_c^\dagger L^i \overline{e}^{i}\right) = 0\,,
\end{aligned}
\tag{9}
$$

and thus the corresponding Yukawa Lagrangian for the three families of quarks and leptons including left-handed neutrinos reads

$$
\begin{aligned}
L_\mathrm{Y} \;=\; & y_{u^{2,3}} H_c Q^i \overline{u}^{2,3} + y_{d^3} H_c^\dagger Q^i \overline{d}^{3} + \; y_{e^i} H_c^\dagger L^i \overline{e}^{i} + y_{u^1} H_d Q^i \overline{u}^{1}\\
& + y_{d^{1,2}} H_d^\dagger Q^i \overline{d}^{1,2} + y_{\nu^i} M_\mathrm{s}^{-1}\left(H_d L^i\right)^2\,,
\end{aligned}
\tag{10}
$$

where $y$s are coupling constants and $M_\mathrm{s}$ being the string mass scale. After the electroweak symmetry breaking and explicating the resulting coupling terms by using the attributed fermion family indices, we get

$$
\begin{aligned}
L_{\mathrm{Y}_\mathrm{mass}} \;=\; & \left(y_c Q\overline{c} + y_t Q\overline{t} + y_b Q\overline{b} + y_{e^i} L^i \overline{e}^{i}\right) v_c + \left(y_u Q\overline{u} + y_d Q\overline{d} + y_s Q\overline{s}\right) v_d\\
& + \left(\frac{y_{\nu^i}}{M_\mathrm{s}} L^i L^i\right) v_d^2\,.
\end{aligned}
\tag{11}
$$

This Lagrangian exhibits clearly 3 fermion mass scales. The first related to $v_c$, the second to $v_d$ and the third to the highly suppressed term $v_d^2/M_\mathrm{s}$. The resulting fermion masses read then

$$
\begin{aligned}
m_c &= y_c v_c\,, \qquad m_b = y_b v_c\,, \qquad m_\tau = y_t v_c\,,\\
m_u &= y_u v_d\,, \qquad m_d = y_d v_d\,, \qquad m_s = y_s v_d\,,\\
m_{e^i} &= y_{e^i} v_c\,, \qquad m_{v^i} = \frac{1}{M_\mathrm{s}} y_{v^i} v_d^2\,.
\end{aligned}
\tag{12}
$$

In this view, with the Yukawa constants hierarchy: $y_t \sim 1 > y_c > y_b > y_s \gg y_d \sim y_u \gg y_{e^i}$ and the extremely small effective neutrino Yukawa constant $y_{v^i} v_d/M_\mathrm{s} \ll y_{e^i}$, the heavy known quarks $t$, $b$ and $c$ and leptons $e^i = e$, $\mu, \tau$ get their masses from the first Higgs VEV $v_c$, the light known quarks $u$, $d$ and $s$ get their masses from the second Higgs VEV $v_d$ and the left-handed neutrinos $v^i = v^e, v^\mu$, $v^\tau$ get their masses from the second Higgs

VEV $v_d$ through a high-dimension operator suppressed by the mass scale $M_\mathrm{s}$ explaining their tiny masses. This leads to

$$v_c \sim \mathrm{GeV} > v_d \sim \mathrm{MeV} \gg \frac{v_d^2}{M_\mathrm{s}} \sim \mathrm{eV} \tag{13}$$

which shows that we have indeed three scales responsible for the generation of fermion masses and giving a partial explanation for the observed mass hierarchies. Since the mass scale $M_\mathrm{s}$ is taken as the low-string scale at which neutrino masses have origin, it is interesting to approximate its value. For neutrino masses upper bound $m_{\nu_\tau} \lesssim 2$ eV from (12), we can get

$$M_\mathrm{s} \sim 10^6 \ \mathrm{GeV}\,. \tag{14}$$

This is a stringy prediction for new physics beyond SM that we have been able to derive from the known data; it would be useful to ask what is the physical implication of these deductions for low-energy physics and whether the Large Hadron Collider experiment is able to confirm the stringy physics directly.

## 4. Conclusions

In this work, we have discussed a string-inspired extended SM where the corresponding effective low-energy theory emerges from intersecting D-branes. The interaction terms are subject to additional restrictions from U(1) symmetries and anomaly cancellation conditions. In particular, we have discussed a four stacks of intersecting D-brane configuration in orientifolded geometries realizing the SM spectrum with two Higgs doublets necessary for generating the Yukawa couplings for all fermion families. The extended Higgs sector generates rich mass spectrum and presents interesting features. In addition to the charged physical Higgses that are experimentally important, the fermion masses have been split into three scales GeV, MeV and eV, that we have bounded from the known data. This hierarchy offers a partial solution to the fermion masses problem as well as the smallness of neutrino masses.

This is among others one simple string-inspired model for flavours that is able to address some open issues in the SM within the allowed window by assuming that the scale of new physics, taken as the string scale, is closely related to the scale at which the neutrino masses are generated.

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