Quark mass hierarchies in D-brane realizations of the Standard Model

P. Anastasopoulos[1,2]*and Andrea Lionetto[1]†

[1] *Dip. di Fisica & Sez. I.N.F.N. Università di Roma "Tor Vergata",
Via della Ricerca Scientifica, 1 - 00133 Roma, ITALY*

[2] *Department of Physics, CERN-Theory Division,
1211 Geneva 23, SWITZERLAND*

ABSTRACT: This proceeding is based on arXiv:0905.3044. We analyze the problem of the hierarchy of masses and mixings in orientifold realizations of the Standard Model. We present a bottom-up brane configuration that can generate such hierarchies.

*Pascal.Anastasopoulos@roma2.infn.it
†Andrea.Lionetto@roma2.infn.it

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1. Introduction, motivation and scope

One of the biggest puzzles in the Standard Model (SM) is the origin and hierarchy of masses and mixings. When it comes to masses the scale of the problem is enormous: one needs to explain a range of masses that spans fifteen orders of magnitude between the mass of the lightest neutrino to the top quark. Moreover the pattern of mixings is interesting. In the quark sector the first and second family mix strongly while all other mixings are small. In the lepton sector, all mixings measured so far are maximal. It seems to suggest that as we move up in mass mixings tend to become smaller.

There are several ideas on the origin of mass, which can be roughly lumped into four classes: radiative mechanisms, texture zeros, family symmetries and seesaw mechanisms. However, the classes are not completely disjoint. In particular texture zeros can be considered as a class of family symmetries as they are usually implemented via a discrete symmetry.

In this work we analyze SM mass patterns and hierarchies in the context of open string theory vacua, alias orientifolds [1]. These vacua allow for a small string scale M_s and provided a fresh new perspective in the search for the SM [2]. Moreover they allowed a bottom-up approach [3, 4] to building the SM, by utilizing the geometrized language offered by D-branes supporting the SM interactions and particles.

In this framework, the SM particles are strings which are stretched between some stacks of D-branes (aka the "SM branes"). Typically there is a need of some extra stacks (aka the "hidden sector") which do not include light observable-hidden strings [5].

Anomalous U(1) symmetries are ubiquitous in orientifolds [6]. It has been argued early on [2, 3], that any SM orientifold realization must have at least one and generically three anomalous U(1) symmetries, that make the most characteristic signature of orientifold vacua. Their phenomenological implications are diverse [7].

Their most important property, that impacts importantly on the dynamics of the D-brane stack is that they provide numerous selection rules on the effective couplings. In particular, they may be responsible for the absence of the μ-term, Yukawa couplings,
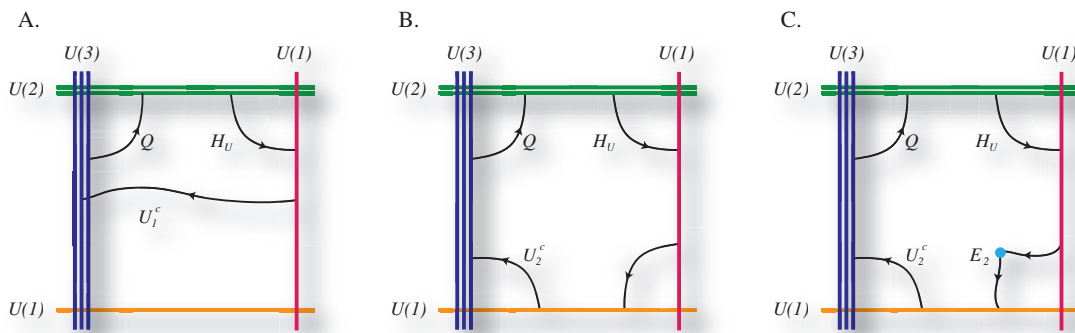
Figure 1: The three types of mass generating terms: The configuration A allows for a Yukawa term. However, in the B and C cases no Yukawa terms can be generated. In the B case there is a higher order term due to the presence of a field Φ, while in the C case there is a contribution from an instanton term E_2.

baryon and lepton violating couplings etc. However, as anomalous U(1)'s are effectively broken as gauge symmetries, the selection rules they provide need qualification. As the breaking of the gauge symmetry happens via the mixing with RR forms, the global U(1) symmetry remains at this stage intact. There are two types of realizations of anomalous U(1) symmetries as global symmetries. If D-terms force charged fields to obtain vev's then the global U(1) symmetry is broken. If on the other hand no vev's are generated the anomalous U(1) global symmetry remain intact in perturbation theory.

However, the story must change beyond perturbation theory for two reasons. The first is that we do not expect exact (compact) global symmetries to survive in a gravitational theory. The second (in agreement with the first) is that there are always non-perturbative effects that violate the associated global symmetry. The argument is simple. A U(1) transformation involves a shift of RR field. The associated D-instanton effect which is charged under the same RR field (the Stuckelberg axion) will violate by definition the associated global U(1) symmetry. The effect is a D-instanton effect, whose field theory limit sometimes may admit a gauge instanton interpretation [8].

Therefore, there are several effects that can produce hierarchically different Yukawa-like couplings (figure 1)[1] [12]-[15].

- Tree-level cubic Yukawa couplings: This is the generic case when such couplings are allowed. Their coefficient depends in general on several ingredients. It is always proportional to the ten-dimensional dilaton but also internal volumes, and other backgrounds fields (internal magnetic fields, fluxes) enter. They may be correlated

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[1]We would like to mention that higher order or instantonic terms might bring in the potential unwanted couplings too. By "unwanted couplings" we mention all terms that could make the model phenomenologically uninteresting, like baryon or lepton number violating terms or couplings that lead to fast proton decay. One can easily assume that such couplings are not present due to the absence of the related instanton. However, if an unwanted coupling in the superpotential has the same violation of U(1) charges as a mass term that has a non-zero instanton contribution, then there is a non-trivial instanton contribution for this term and thus it will generated with a similar strength. Such models must be excluded.

with the associated gauge couplings if the fields participating come from overlapping D-branes. They may also be free of volumes if the branes intersect at points. Such variations are enough some times to explain the mass hierarchy inside a family [3].

- Higher order couplings: These are couplings that appear beyond the cubic level. They necessarily involve more fields than the SM fields. These extra fields must obtain an expectation value in order for an effective Yukawa coupling to be generated. Then such couplings compared to the previous case carry an extra factor of $(\langle\phi\rangle/M_s)^n$ with n a positive integer. Depending on the compactification the string scale may be replaced by a compactification scale. If $\langle\phi\rangle \ll M_s$ this generates a hierarchy in the associated Yukawa coupling.

- D-Instanton-generated couplings: Such couplings violate the anomalous U(1) symmetries. They are suppressed by exponential instanton factors of the form $e^{-1/g}$ where g is linearly related to the ten-dimensional coupling constant and depends also on the volume of the cycle the D-instanton is wrapped-on, as well as on magnetic fields, fluxes, etc. In the particular case of gauge instantons g is the square of the associated gauge coupling. In the well-controlled regime, $g \ll 1$ and multi-instantons are suppressed. Beyond the instanton-action factor, instanton-generated couplings carry a characteristic scale. This is determined by the string scale, or other volume factors affecting the world-volume factor of the D-instanton. Finally there is a one-loop determinant that is generically of order $\mathcal{O}(1)$.

- Not generated couplings: Remain zero as no vev or instanton can generate them.

In [14], we explore different effects that are prone to generate interesting hierarchies between fermion masses. In particular, for four brane stacks we found vacua in which for each fermionic mass matrix, the highest mass scale is related to Yukawa terms, the intermediate mass scale to instantons while the lowest scale to higher order terms. The CKM matrix computed for this model agrees with the experimental result. We would like to mention that our study was based on the bottom-up approach. The goal was to identify D-brane configurations that are promising when it comes to generating the fermion hierarchy. The next step will be to construct such interesting D-bane configurations.

2. An example of the new hierarchy idea

In order to show our results we focus on a vacuum with four stacks of branes: a stack of three (color stack), a stack of two (weak stack), and two single branes. The hypercharge

embedding is $Y = \frac{1}{6}Q_3 + \frac{1}{2}Q_1 + \frac{1}{2}Q'_1$. The SM particles are:

$$
\begin{array}{llll}
Q_1 & : (\ V, \ V, \ 0, \ 0) & , \ Q_2, Q_3 & : (\ V, \ \bar{V}, \ 0, \ 0) \\
U_1^c & : (\ \bar{V}, \ 0, \ \bar{V}, \ 0) & , \ U_2^c \, U_3^c & : (\ \bar{V}, \ 0, \ 0, \ \bar{V}) \\
D_1^c & : (\ \bar{V}, \ 0, \ V, \ 0) & , \ D_2^c \, D_3^c & : (\ \bar{V}, \ 0, \ 0, \ V) \\
L_1^c & : (\ 0, \ V, \ \bar{V}, \ 0) & , \ L_2^c \, L_3^c & : (\ 0, \ V, \ 0, \ \bar{V}) \\
E_1^c & : (\ 0, \ 0, \ 0, \ S) & , \ E_2^c & : (\ 0, \ 0, \ 0, \ S) \quad , \quad E_2^c \ : (\ 0, \ 0, \ S, \ 0) \\
N_{1,2,3}^c & : (\ 0, \ A, \ 0, \ 0) & & \\
H_u & : (\ 0, \ \bar{V}, \ V, \ 0) & , \ H_d & : (\ 0, \ \bar{V}, \ \bar{V}, \ 0)
\end{array}
\tag{2.1}
$$

where V, A, S are the fundamental, antisymmetric and symmetric representation of the corresponding group with U(1) charges: $V \to 1$, $A \to 2$, $S \to 2$. All bars denote the conjugates rep with opposite U(1) charge. We also consider two additional scalars with zero hypercharge ϕ_1 and ϕ_2, coming from the non chiral part of the spectrum:

$$
\phi_1 \ : \ (0, 0, \ \bar{V}, V) \quad , \quad \phi_2 \ : \ (0, 0, V, \ \bar{V}) \ .
\tag{2.2}
$$

The above spectrum satisfies all irreducible anomaly cancellation conditions and additionally some extra conditions that are coming from tadpole cancellation.

For the above spectrum (2.1): Q_1, U_1^c, and H_u form a perturbative Yukawa term $Q_1 U_1^c H_u \to m_{1,1} \sim \langle H_u \rangle$. On the other hand Q_1, $U_{2,3}^c$, and H_u do not allow for such term. The contribution of the extra field ϕ_1 provides a higher term $Q_1 U_1^c H_u \phi_1 \to m_{1,2} \sim \langle H_u \rangle \langle \phi_1 \rangle / M_s$. For the rest of the entries in the mass matrix, neither $\phi_{1,2}$ can form gauge invariant couplings, thus we assume the presence of instantons, like: $Q_2 U_1^c H_u \to m_{2,1} \sim \langle H_u \rangle e^{-Vol}$. Following this spirt we evaluate the corresponding mass matrices for the quarks:

$$
M_U = \langle H_u \rangle \begin{pmatrix} g_1 & g_2 v_{\phi_1} & g_3 v_{\phi_1} \\ g_4 E_1 & g_5 E_2 & g_6 E_2 \\ g_7 E_1 & g_8 E_2 & g_9 E_2 \end{pmatrix} \quad , \quad M_D = \langle H_d \rangle \begin{pmatrix} q_1 & q_2 v_{\phi_2} & q_3 v_{\phi_2} \\ q_4 E_1 & q_5 E_3 & q_6 E_3 \\ q_7 E_1 & q_8 E_3 & q_9 E_3 \end{pmatrix}
\tag{2.3}
$$

$$
M_L = \langle H_d \rangle \begin{pmatrix} l_1 E_4 & l_2 v_{\phi_1} & l_3 \\ l_4 E_4 & l_5 v_{\phi_1} & l_6 \\ l_7 E_4 & l_8 v_{\phi_1} & l_9 \end{pmatrix}
$$

$$
M_N = \begin{pmatrix}
0 & 0 & 0 & g_{11}\langle H_u \rangle E_1 & g_{12}\langle H_u \rangle E_1 & g_{13}\langle H_u \rangle E_1 \\
0 & 0 & 0 & g_{21}\langle H_u \rangle E_1 & g_{22}\langle H_u \rangle E_1 & g_{23}\langle H_u \rangle E_1 \\
0 & 0 & 0 & g_{31}\langle H_u \rangle E_1 & g_{32}\langle H_u \rangle E_1 & g_{33}\langle H_u \rangle E_1 \\
g_{11}\langle H_u \rangle E_1 & g_{21}\langle H_u \rangle E_1 & g_{31}\langle H_u \rangle E_1 & q_{11}M_s E_5 & q_{12}M_s E_5 & q_{13}M_s E_5 \\
g_{12}\langle H_u \rangle E_1 & g_{22}\langle H_u \rangle E_1 & g_{32}\langle H_u \rangle E_1 & q_{21}M_s E_5 & q_{22}M_s E_5 & q_{23}M_s E_5 \\
g_{13}\langle H_u \rangle E_1 & g_{23}\langle H_u \rangle E_1 & g_{33}\langle H_u \rangle E_1 & q_{31}M_s E_5 & q_{32}M_s E_5 & q_{33}M_s E_5
\end{pmatrix}
$$

where g_i, q_i, l_i, g_{ij} and q_{ij} are dimensionless couplings assumed to be of the same order $[0.1 - 0.6]$. Also $v_{\phi_1} = \langle \phi_1 \rangle / M_s$, $v_{\phi_2} = \langle \phi_2 \rangle / M_s$ and $E_i = e^{-Vol_{I_i} I_i}$ are the dimensionless instantons.

For the present vacuum we were able to find solutions where there is 1-1 correspondence between the fermion masses in each family and the Yukawa, higher order and instantonic

terms:

$$\begin{aligned}
\langle H_u \rangle &\sim m_t &, \qquad \langle H_d \rangle &\sim m_b \\
E_1 &\sim E_2 \sim m_c/m_t , &\qquad E_3 &\sim E_4 \sim m_s/m_b \quad , E_5 \sim 0.654 \\
v_{\phi_1} &\sim m_u/m_t &, \qquad v_{\phi_2} &\sim m_d/m_b
\end{aligned} \tag{2.4}$$

where m_i are the masses of the corresponding quarks [11], and all couplings $|g_i|$, $|q_i|$, $|l_i|$, $|g_{ij}|$, $|q_{ij}|$ are within the range $[0.1, 0.6]$. Notice that E_5 appears always multiplied by M_s and it's value triggers the seesaw mechanism. It is the only value that changes if we repeat the same mechanism at other scales apart from 1 TeV.

Using the above values for the couplings and vev's, we can proceed and evaluate the Cabibbo - Kobayashi - Maskawa Matrix (CKM). For the above vacuum, the matrix is:

$$\mathrm{CKM}(1\,\mathrm{TeV}) = \begin{pmatrix} 0.970 & 0.240 & 0.007 \\ 0.240 & 0.970 & 0.013 \\ 0.010 & 0.011 & 0.999 \end{pmatrix} \tag{2.5}$$

that has to be compared with the experimental data [16]:

$$\mathrm{CKM(Data)} = \begin{pmatrix} 0.97419 \pm 0.00022 & 0.2257 \pm 0.0010 & 0.00359 \pm 0.00016 \\ 0.2256 \pm 0.0010 & 0.97334 \pm 0.00023 & 0.0415 \pm 0.001 \\ 0.00874^{+0.00026}_{-0.00037} & 0.0407 \pm 0.0010 & 0.999133^{+0.000044}_{-0.000043} \end{pmatrix} \tag{2.6}$$

In this vacuum, only the μ-term and same specific Yukawa couplings share the same instanton E_1 and therefore $\mu \sim E_1 \sim m_c/m_t$. This is not problematic for low string scale models.

3. Conclusions

In the field theory context several ideas has been proposed to explain the fermionic mass hierarchy in the SM such as radiative mechanisms, texture zeros, family symmetries and seesaw mechanisms. None of them can be separately implemented in D-brane constructions whose aim is to reproduce the SM. New types of textures arise due to extended symmetries, i.e. several (anomalous) $U(1)$'s.

In this proceeding we argued that SM mass hierarchies can come from different kind of terms that are generically present in D-brane vacua. Such terms are the perturbative Yukawas, higher order terms (with the contribution of an extra scalar field) and instantonic terms. Hierarchies are due to the 1-1 correspondence between the masses of the fermions in each mass matrix with the three terms mentioned above.

In this framework we have shown a model that materialize this idea. However the model is constructed on the bottom-up fashion and a consistent open string vacuum has to be still constructed.

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