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Standard model on D-branes

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Abstract. I briefly outline previous work on getting the (supersymmetric) standard model from string theory, and then describe two ecent attempts using D-branes. The first uses D3- and D7-branes and gives a supersymmetric standard model with extra vector-like matter and an intermediate unification scale. The second uses intersecting D4-branes and yields a non-supersymmetric spectrum with TeV-scale unification.

Keywords. Standard; D-branes.

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1. Introduction

The earliest attempts to get realistic phenomenology from string theory used a geometric approach. The pioneering work of Candelas *et al* [1] showed that the requirement of a four-dimensional space-time supersymmetric spectrum in (weakly-coupled) heterotic string theory necessitates compactifying the unobserved six spatial dimensions on a Calabi–Yau manifold $\mathscr{M}$. The number of fermion generations is given by $\frac{1}{2}|\chi(\mathscr{M})|$, where $\chi(\mathscr{M})$ is the Euler character of $\mathscr{M}$. Since $\chi(\mathscr{M})$ is generically a multiple of 12, the construction of a realistic model (with $\chi(\mathscr{M}) = \pm 6$) requires the deletion of unwanted chiral matter using a (freely-acting) point group and/or Wilson lines; see ref. [2] for a review. Similar observations apply to orbifold constructions [3], obtained by quotienting a six-torus with a non-freely-acting point group, for example $\mathscr{M} = T^6/Z_N$ with $N = 3,4,6,7,8,12$; see [4] for a review. It was later realised that such a geometrical interpretation could be dispensed with. The extra dimensions permit the cancellation of the ghost central charge, and this can alternatively be done using world-sheet degrees of freedom that do not admit of a geometrical interpretation. Specifically, the Gepner models [5] use $\mathscr{N} = 2$ superconformal coset theories to supply the nine units of central charge, previously contributed by the six compact bosonic dimensions and their fermionic counterparts. In the real fermion models [6,7], all of the erstwhile bosonic degrees of freedom are fermionised, and different models arise from imposing different boundary conditions on these fermioms. The most realistic models to date have probably been constructed using the real fermion approach [8].

All of these attempts use a top-down approach, in which the initial set-up generally has too much symmetry and too much matter. The task of the model 'builder' is to reduce the symmetry and excise or hide the unwanted matter. Nevertheless, there is often residual unwanted matter, sometimes with exotic quantum numbers; the string scale is required to be about 4×10^{17} GeV, far larger than the 'observed' unification scale of 2×10^{16} GeV; and breaking supersymmetry requires non-perturbative physics, not obviously included in the model. The development of (strongly-coupled) heterotic M-theory [9] reverted to the geometrical approach and solved the problem of the disparity between the string and unification scales. There is an additional eleventh (spatial) dimension, compactified on an orbifold interval S^1/Z_2, that is 'seen' by gravity, but not by the standard-model fields, and the disparity can be removed by adjusting the length of the additional dimension. However, although there are prescriptions for ensuring that there are just three chiral generations of matter, getting the standard model gauge group seems to require the spontaneous breaking of a grand unification theory (GUT) using Wilson lines symmetry [10], reminiscent of the technique used in the weakly-coupled case. Another new feature of heterotic M-theory is the possibility of objects with five spatial dimensions, 5-branes, situated in the orbifold interval. These too can accommodate gauge symmetries and chiral matter that is hidden from matter on the pair of 9-branes at the ends of the interval. The D-brane models that we are principally concerned with share this feature.

The discovery [11] that, in general, compactified type II superstring theories also contain D-branes has led to the re-evaluation of the phenomenological possibilities of type I/II theories. Initially the approach followed [12–15] was to start with an orientifold compactification of the theory $\mathscr{M}=T^6/G$, where $G=G_1+\Omega G_2$ with $G_{1,2}$ discrete internal symmetry groups, and Ω the world-sheet parity operator that exchanges left- and right-movers. The appearance of Ω means that some D9-branes are required in order to cancel the (untwisted) Ramond–Ramond (RR) tadpole divergences that arise from vacuum amplitudes with Klein-bottle world sheets. Depending on G, models with both D9- and D5-branes arise, the latter deriving from the necessity to cancel twisted RR tadpoles. (There is a dual picture containing D7- and D3-branes.)

The most realistic orientifold models constructed to date [16] have contained anti-D-branes as well as D-branes. In that case, supersymmetry is broken explicitly in an anti-D-brane hidden sector, with supersymmetry breaking being transmitted gravitationally to the observable sector. Since we expect the scale of supersymmetry breaking in the anti-D-brane sector to be the string scale m_{s}, the mass of the sparticles in the observable sector may be expected to be of order $m_{\mathrm{s}}^2/m_{\mathrm{P}}$, where m_{P} is the four-dimensional Planck mass. For sparticle masses of order 1 TeV we require $m_{\mathrm{s}}\sim10^{11}$ GeV. Then, unification of gauge coupling constants at a scale of order 10^{11} GeV is to be expected (modulo some subtleties to do with Kaluza–Klein modes and winding modes [14]). Interestingly, orientifold models can contain the extra matter required to allow the renormalization group equations to run to unification at this lower scale, though with an unconventional value for the coupling constant ratio g_1^2/g_3^2 which arises from the identification of the weak hypercharge with a specific non-anomalous $U(1)$. The string scale is necessarily of order the Planck mass in weakly coupled heterotic theories. However, in type II orientifold theories with isotropic compactification on a scale $R=m_c^{-1}$, and gauge fields for the observable sector arising from D9-brane open strings, we have the relationship

$$\frac{m_c^3}{m_{\mathrm{s}}^2}=\frac{\alpha_X m_{\mathrm{P}}}{2\sqrt{2}} \tag{1}$$

where $\alpha_X\simeq\frac{1}{24}$. Then, the string scale can be adjusted to an intermediate scale by adjusting m_c.

There is a potential problem of too rapid proton decay with such a low string scale. However, type IIB orientifold theories differ from weakly coupled heterotic theories in a vital respect. Whereas in the latter case anomalous $U(1)$ factors in the gauge group disappear completely after the anomaly has been cancelled by a Green–Schwarz mechanism, in the former case anomalous $U(1)$ gauge symmetries are able to survive as global symmetries in the low energy theory. This different behaviour is a consequence of the fact that in the heterotic case the Fayet–Iliopouls term for the anomalous $U(1)$ involves the dilaton field whose expectation value is controlled by the strength of the gauge coupling constant, whereas in the type IIB orientifold case it involves orbifold blowing-up modes whose expectation values can approach zero as the orbifold limit is approached. These global $U(1)$ symmetries may forbid dangerous terms in the superpotential. They are also a promising way of obtaining a quark–lepton mass matrix with appropriate hierarchies.

2. Three-generation type IIB orbifold model

We noted in the previous section that by considering toroidal compactifications of type I superstrings we are led naturally to type II strings with D-branes. It is well known that type IIB theory quotiented with the Z_2 generated by the world sheet parity operator Ω is just type I theory. Thus orbifold compactifications of type I are orientifold compactifications of type IIB. However, if we consider only orbifold compactifications of type IIB and their T-duals we will never encounter open strings or D-branes, but there is no reason why we should not *start* from type IIB with D-branes. This avoids an unwelcome feature of the orientifold models, which is that they start with a large ($SO(32)$) gauge group with no immediate resemblence to the standard-model group. Such models necessarily use a top-down approach. In contrast, the orbifold models permit a bottom-up approach [17].

It is convenient to work in the T-dual picture in which we start with a stack of six D3-branes at a Z_3 orbifold fixed point that without loss of generality may be taken to be the origin O. Then the gauge vector bosons (and gauginos) of $U(6)$ arise from open strings that begin and end on the D3-branes. The gauge group is broken to $U(3)\times U(2)\times U(1)$ when the action of the generator θ of the Z_3 point group is embedded in the Chan–Paton indices by the matrix

$$\gamma_{\theta,3} = \mathrm{diag}(I_3, \alpha I_2, \alpha^2) \tag{2}$$

where $\alpha \equiv \mathrm{e}^{2\pi i/3}$. Each of the $U(n)$ groups has its own $U(1)$ charge Q_n, and the only non-anomalous combination is proportional to (the weak hypercharge)

$$-Y \equiv \frac{1}{3}Q_3 + \frac{1}{2}Q_2 + Q_1. \tag{3}$$

It follows that at the unification scale m_X the three standard model gauge couplings satisfy

$$\alpha_3(m_X) = \alpha_2(m_X) = \frac{11}{3}\alpha_1(m_X) \tag{4}$$

corresponding to

$$\sin^2\theta_W(m_X) = \frac{3}{14} = 0.214. \tag{5}$$

The 33 sector also produces $n_G = 3$ generations of (supersymmetric) chiral matter in the (bi-findamental) representations $(\mathbf{3},\mathbf{2})_{1/6} + (\mathbf{2},\mathbf{1})_{1/2} + (\mathbf{1},\bar{\mathbf{3}})_{-2/3}$ of the gauge group, corresponding to $Q_L + H_u + u_L^c$. It is easy to see that this has non-Abelian anomaly, which derives from the existence of uncancelled Ramond–Ramond (RR) twisted tadpoles at O. To cancel them, D7-branes passing through O must be introduced. We choose to introduce only $D7_3$-branes with $z_3 = 0$, wrapping the $z_{1,2}$ complex planes, with point-group embedding [18]

$$\gamma_{\theta,7_3} = \mathrm{diag}(I_u, \alpha I_{u+3}, \alpha^2 I_{u+6}) \tag{6}$$

so that the twisted-tadpole cancellation condition

$$3\mathrm{Tr}\gamma_{\theta,3} - \gamma_{\theta,7_1} - \gamma_{\theta,7_2} + \gamma_{\theta,7_3} = 0 \tag{7}$$

is satisfied (for any integer $u \geq 0$). The D7-branes generate their own gauge group, and the $37_3 + 7_3 3$ states complete the $n_G = 3$ chiral generations together with extra vector-like matter. In all we get

$$\begin{aligned} &3(Q_L + u_L^c + d_L^c + L + H_u + H_d + e_L^c) \\ &+(u+3)(d_L^c + \bar{d}_L^c) + u(H_u + H_d) + u(e_L^c + \bar{e}_L^c). \end{aligned} \tag{8}$$

The introduction of the D7-branes leads to uncancelled twisted tadpoles at the eight other orbifold fixed points with $z_3 = 0$, namely at $(p, \pm 1, 0)$ where $p = \pm 1, 0$, and at $(\pm 1, 0, 0)$. Introducing a suitably chosen Wilson line in the z_2-direction cancels the RR charge at the six points with $z_2 \neq 0$, and at the same time reduces the size of the D7-brane gauge group; we take $u = 3\tilde{u}$ and the D7-brane gauge group is reduced to $[U(\tilde{u}) \times U(1+\tilde{u}) \times U(2+\tilde{u})]^3$. The cancellation at the remaining two fixed points is easily arranged, for example by introducing a stack of six D3-branes at each of them, with θ embedded as at O. The important point is that so far there are no branes, and so no twisted RR charge, at any of the fixed points in the planes $z_3 = \pm 1$. Thus we are free finally to cancel the untwisted tadpoles by putting an appropriate array of anti-branes in (say) the $z_3 = 1$ plane. Precisely what array is used does not matter. The essential point is that it *can* be done, and in a way that ensures cancellation of all twisted RR charges in the $z_3 = 1$ plane. The anti-branes are non-supersymmetric, but this lack of supersymmetry is transmitted to the observable sector (in the plane $z_3 = 0$) only gravitationally, by closed string states. Thus, as noted earlier, requiring that the sparticles have TeV-scale masses requires a string scale (m_s) of order 10^{11} GeV. The question, therefore, is whether the matter content we have derived leads to unification at a scale (m_X) that is consistent with such a string scale.

The (one-loop) renormalization group equations give

$$\sin^2\theta_W(m_Z) = \frac{3}{14} + \frac{11b_2 - 3b_1}{3b_1 + 3b_2 - 14b_3}\left[\frac{3}{14} - \frac{\alpha_{em}(m_Z)}{\alpha_3(m_Z)}\right] \tag{9}$$

where b_i $(i = 1,2,3)$ are the beta function coefficients. The matter content (8) gives $(11b_2 - 3b_1)/(3b_1 + 3b_2 - 14b_3) < 0$, so that $\sin^2\theta_W(m_Z) < 0.214$, whereas

$$\sin^2\theta_W(m_Z)|_{\exp} = 0.231. \tag{10}$$

However, if we assume that we are able to turn on string-scale VEVs for the D7-brane matter fields, then, since the $37_3 + 7_33$ matter fields are coupled to the 7_37_3 matter fields, we are free to give string-scale masses to some or all of the vector-like matter. We therefore consider the possibility of giving masses to α copies of $e^c_L + \overline{e^c_L}$, β copies of $H_u + H_d$, and γ copies of $d^c_L + \overline{d^c_L}$ on a scale larger than m_X. Then using the particle data group values $\alpha^{-1}_{em}(m_Z) = 128.9$ and $\alpha_3(m_Z) = 0.119$, the above renormalization group equation gives

$$\sin^2\theta_W(m_Z) = 0.2275 \quad \text{and} \quad m_X = 1.3\times 10^{10}\ \text{GeV} \tag{11}$$

when $\alpha - \beta = 2$ and $\gamma - \beta = 4$. We can obtain a larger unification scale, and exactly the measured value (10) of $\sin^2\theta_W(m_Z)$, if we assume in addition that some other 77 sector scalars give masses to all of the extra matter, over and above that in the minimum supersymmetric model (MSSM), on a scale m_Y not very much larger than m_Z. Then with $\alpha - \beta = 3$ and $\gamma - \beta = 3$ we get

$$\frac{m_Y}{m_Z} = 13.9 \quad \text{and} \quad m_X = 1.1\times 10^{12}\ \text{GeV}. \tag{12}$$

To avoid a Landau pole in the renormalisation group equation we require $\tilde{u} \leq 1$. Since we also require $\tilde{u} > 0$ so that $\alpha > 0$, which is needed for (11) and (12), the value $\tilde{u} = 1$ is uniquely selected. The values calculated in (11) and (12) are (just about) consistent with the observable supersymmetry breaking arising from gravitational interactions with an anti-brane sector.

The situation with respect to quark and lepton mass hierarchies is the same as for the model discussed by Aldazabal *et al* [17]. In particular, superpotential terms coupling the chiral superfields L, e^c_L and H_d are forbidden at tree level, and to all orders string perturbation theory. As discussed in the Introduction, such $U(1)$ symmetries are expected to survive as *global* symmetries in type I/IIB theories. However, the global and other symmetries do allow the coupling of L and e^c_L to a composite effective Higgs field constructed from $\overline{e^c_L}, H_u$ and some D7-brane chiral fields. Since baryon number $B = \frac{1}{3}Q_3$, it is perturbatively conserved if conservation of Q_3 survives as a global symmetry. Thus the proton is stable. On the other hand, lepton number is not conserved, so another solution of this problem is required if such models are to survive. In any case, the survival of global $U(1)$ symmetries brings other problems. Generically, we expect electroweak spontaneous symmetry breaking to break *both* the local $SU(2)\times U(1)$ gauge symmetry and any global $U(1)$ symmetry. (This certainly applies to Q_2, since H_u and H_d have $Q_2 = \pm 1$.) Spontaneous breaking of a global symmetry gives a massless Goldstone boson at the Lagrangian level, the 'axion', which acquires a non-zero mass because the global $U(1)$ is anomalous. Unfortunately, electroweak spontaneous symmetry breaking generates a 'visible' axion with a mass predicted to be of order 10 keV, and there is by now overwhelming evidence that such an axion does not exist [19].

3. Intersecting brane model

A different route to the standard model was motivated by Berkooz *et al*'s [20] observation that intersecting D-branes can give rise to chiral fermions propagating in the intersection of their world volumes. This led Aldazabal *et al* [21,22] to develop (non-supersymmetric) four-dimensional chiral models from intersecting D4-branes wrapped on 1-cycles of a two-dimensional torus T^2 sitting at a singular point in the transverse four-dimensional space B. The local geometry of B near the singularity is modelled as C^2/Z_N. This gives an $\mathcal{N}=2$ supersymmetric gauge sector. Since the models are generically non-supersymmetric, by virtue of their matter content, solution of the hierarchy problem requires a string scale of not more than a few TeV, although the usual four-dimensional Planck mass can be obtained by making the volume of B large enough. The generic appearance at tree level of scalar tachyons at some intersections allows the interpretation of doublets as electroweak Higgs scalars, provided $|m_{\mathrm{H}}^2| \simeq m_{\mathrm{ewk}}^2$. However, colour-triplet and/or charged singlet tachyons are potentially lethal for theories purporting to be realistic, and we require that such modes are absent [23]. Although these models are non-supersymmetric, they typically contain extra matter, besides that in the standard model, some of which can be light. In particular there are towers of 'gonions', massive (fermionic and bosonic), vector-like matter arising from excitations associated with the angle at which the D-branes intersect, and which are charged with respect to the standard model gauge group. There are also Kaluza–Klein and winding-mode towers of gauge bosons which might also exist below the string scale.

All of these features make intersecting brane models susceptible to experimental investigation in the near future. So it is clearly desirable to construct models that are consistent with existing data. A first step in this programme was made in [22] where some three-generation models with the standard model gauge group were constructed using five stacks of D4-branes, each stack wrapping a different 1-cycle of T^2. Our objective is to construct models which besides satisfying the technical constraints of Ramond–Ramond tadpole and anomaly cancellation [23a] also satisfy more of the phenomenological constraints. Specifically, we require that the models allow the (renormalisable) Yukawa couplings needed to generate mass terms for all standard model matter. The gauge couplings in these models are *not* unified at the string scale; they depend upon the wrapping numbers of the particular stack of D4-branes with which they are associated. So we also require that these values are achieved at a string scale of not more than a few TeV using renormalization group equations starting from the measured values of the gauge coupling strengths at the weak scale.

We take the space B transverse to the D4-branes to be C^2/Z_3, where the point group generator θ of Z_3 acts on the two complex coordinates of B with twist vector $v=\frac{1}{3}(1,-1)$. We too choose the first stack to have $N_1=3$ D4-branes and to be sitting at a singular point of B with wrapping numbers $(n_1,m_1)=(1,0)$. In general the integers n_a and m_a count the number of times that the 1-cycle wrapped by the stack a wraps the two basis 1-cycles defining the torus T^2; there is no loss of generality in this choice of wrapping numbers for the first stack. θ is embedded in the stack as $\gamma_\theta=\alpha^p\mathbf{I}_3$ which generates a $U(3)\supset SU(3)_c$ gauge group. Here $\alpha=\mathrm{e}^{2\pi i/3}$ and $p=0,1$ or 2. As in [22], we also choose the second stack to generate a $U(2)\supset SU(2)_L$ gauge group by taking $N_2=2$ D4-branes and embedding θ as $\gamma_\theta=\alpha^q\mathbf{I}_2$, where $q\neq p$ mod 3. The wrapping numbers are $(n_2,3)$, so that the (12) intersection produces $I_{12}\equiv n_1m_2-n_2m_1=3$ quark doublets Q_L. Note that, unlike in the bottom-up approach [17], the number of generations is not determined by the choice of the point group Z_3. Since further non-Abelian gauge groups are not required, all remaining stacks have just one D4-brane.

We find that at least six stacks are required to satisfy all of our constraints, and the six-stack models that do are parametrised by a single integer n_2. They are given in table 1. The (massless) matter content of these models is easily determined using the results of

Table 1. Multiplicities, wrapping numbers and Chan–Paton phases for the six-stack models ($p \neq q \neq r \neq p$).

Stack a	N_a	(n_a, m_a)	γ_θ
1	3	$(1,0)$	$\alpha^p \mathbf{I}_3$
2	2	$(n_2,3)$	$\alpha^q \mathbf{I}_2$
3	1	$(1-n_2,-3)$	α^q
4	1	$(1-n_2,-3)$	α^q
5	1	$(2,0)$	α^r
6	1	$(-1,0)$	α^p

[21], and in our case it turns out to be independent of n_2. Besides the $n_G = 3$ generations of chiral matter, we get the following extra vector-like fermions and (scalar) Higgs doublets

$$3(e_L^c + \overline{e_L^c}) + 9(L + \overline{L}) + 6H.$$

The gauge group is that of the standard model apart from some extra $U(1)$ factors provided $n_2 \neq 0$ mod 3. The weak hypercharge Y is given by the superposition of the charges Q_a that generate the $U(1)_a$ factors associated with the ath stack:

$$-Y = \frac{1}{3}Q_1 + \frac{1}{2}Q_2 + Q_3 + Q_5 + Q_6. \tag{13}$$

The mass of the tachyonic Higgs doublets [22] *does* depend on n_2 and is given by

$$m_{\mathrm{H}}^2 = -\frac{m_{\mathrm{s}}^2}{2\pi}\frac{\varepsilon(n_2-\delta)}{|\delta||1-\delta|} \tag{14}$$

where m_{s} is the string scale; ε and δ are related to the parameters defining the torus wrapped by the D4-branes. If R_1 and R_2 are the radii of the two fundamental 1-cycles, and θ is the angle between the two vectors defining the lattice, then

$$\varepsilon \equiv 2|\cos(\theta/2)| \tag{15}$$

$$\delta \equiv n_2 - 3R_2/R_1. \tag{16}$$

The above formula for m_{H}^2 is valid so long as $\varepsilon \ll 1$, but in any case $m_{\mathrm{H}}^2 \ll m_{\mathrm{s}}^2$ is required for consistency of the standard model without major contamination by string effects. m_{H}^2 also sets the scale for the various vector-like gonions that arise below the string scale.

The matter content detailed above contributes to the running of the coupling strengths $\alpha_i(\mu)$ $(i = 3, 2, Y)$ when the renormalization scale μ is greater than the mass of the relevant matter, and we can hence evaluate the coupling strengths at the (unknown) string scale $\mu = m_{\mathrm{s}}$ at which their values are given in terms of the type II string coupling λ_{II} and the wrapping numbers relevant to the stack producing the gauge group. We assume that the $\mathcal{N} = 2$ gauge vector supermultiplet is massless over the entire range starting from $\mu = m_Z$. There are also Kaluza–Klein and winding mode partners of gauge bosons, but we can arrange that these have masses above m_{s}, so do not contribute to the renormalisation group running between m_Z and m_{s}. The ratios of the gauge coupling strengths, which depend on the length of the cycles (n_a, m_a), are independent of the unknown λ_{II}:

Table 2. Predicted values of the physical Higgs mass $m_h \equiv |2m_{\rm H}^2|^{1/2}$ and the ratio $a \equiv R_1/R_2$ of the compactification radii in the case of an $\mathcal{N}=2$ supersymmetric gauge sector.

n_2	$m_{\rm s}$ (TeV)	ε	$m_{\rm h}$ (GeV)	a
5	1.0	0.185	137	0.294
5	1.2	0.182	165	0.297
5	1.9	0.176	265	0.304
5	3.0	0.170	425	0.312
10	1.0	0.123	137	0.198
10	1.2	0.121	165	0.199
10	1.9	0.116	265	0.202
10	3.0	0.111	425	0.206
20	1.0	0.074	137	0.119
20	1.2	0.073	165	0.120
20	1.3	0.072	179	0.120
20	1.9	0.069	265	0.121
20	3.0	0.066	425	0.122

$$\frac{\alpha_2^{-1}(m_{\rm s})}{\alpha_3^{-1}(m_{\rm s})} = \left[\delta^2 + n_2(n_2-\delta)\varepsilon^2\right]^{1/2} \tag{17}$$

$$\frac{\alpha_Y^{-1}(m_{\rm s})}{\alpha_3^{-1}(m_{\rm s})} = \frac{10}{3} + \frac{\alpha_2^{-1}(m_{\rm s})}{2\alpha_3^{-1}(m_{\rm s})} + \left[(1-\delta)^2 + (n_2-1)(n_2-\delta)\varepsilon^2\right]^{1/2} \tag{18}$$

where ε and δ are defined in eqs (15) and (16). Thus, by substituting the solutions of the renormalisation group equations into (17) and (18), we obtain two constraint equations on ε and δ. We solve these constraints numerically for a range of values of the parameters n_2 and $m_{\rm s}$, and using the latest values for the coupling strengths at the weak scale. For given values of the parameters n_2 and $m_{\rm s}$, the calculated values of ε and δ determine the mass of the tachyonic Higgs doublets using eq. (14), and we assume that this determines the mass $m_{\rm h}$ of the physical Higgs particle from

$$m_{\rm h}^2 = |2m_{\rm H}^2|. \tag{19}$$

The ratio a of the compactification radii is then determined from eq. (16). Our results are summarised in table 2 for the case of an $\mathcal{N}=2$ supersymmetric gauge sector. We find that it is easy to obtain values of $\varepsilon \ll 1$ and $m_{\rm H}^2 \ll m_{\rm s}^2$ consistently with a string scale $m_{\rm s}$ not more than a few TeV. Physical Higgs masses not much greater than the current LEP bound ($m_{\rm h} > 114$ GeV) [24] are found for some values of the parameters, but these are probably excluded because they are necessarily accompanied by charged, vector-like gonions with similar masses. Similar results are obtained [23] if we assume that all of the $\mathcal{N}=2$ superpartners have masses greater than $m_{\rm s}$, and so do not contribute to the renormalisation group running below $m_{\rm s}$. The general behaviour is that with $m_{\rm s}$ fixed, $m_{\rm h}$ falls slowly as n_2 increases, whereas with n_2 fixed, $m_{\rm h}$ increases as $m_{\rm s}$ increases.

Like the bottom-up models [17,18], with D-branes located at R/Z_3 orbifold fixed points, the mixed anomalies in the present intersecting brane model are also cancelled by a generalised Green–Schwarz mechanism mediated by twisted Ramond–Ramond fields. However, because of the lack of supersymmetry in the present set-up, it is unclear whether the anomalous $U(1)$s survive as global symmetries. Since the Yukawa terms needed to generate the required masses certainly conserve all of the $U(1)$ symmetries, it seems likely that they *are* global symmetries. If so, then keV scale axions also arise in these models. Unlike the models of [17,18], the present model possesses lepton mass terms as well as quark mass terms, at renormalisable level provided we assign the lepton doublets to (25) interesections rather than the (24) intersection. Also unlike these models our model is free from lepton number violating terms at renormalisable level. On the other hand, it shares with those models the virtue of proton decay being perturbatively forbidden [22], although baryon number B non-conservation with ΔB even *is* allowed. Masses for the vector-like matter are more problematic because of the lack of gauge-singlet scalars amongst the tachyonic states to develop vacuum expectation values. Three of the $L+\bar{L}$ states and three of the $e^c_L + \bar{e}^c_L$ states acquire masses at tree-level via the usual $Le^c_L H$ and $\bar{L}\bar{e}^c_L H^\dagger$ couplings. The remaining six $L+\bar{L}$ pairs are able to acquire masses via non-renormalisable terms of the form $L\bar{L}HH^\dagger$; the latter are suppressed by a factor $\langle H\rangle/m_s$.

In conclusion, we have found a unique class of six-stack standard-like models that give masses to all of the matter, and are free of charged singlet (and colour triplet) tachyons. We find it relatively easy to ensure that the Higgs mass is small compared with the string scale in the range 1–3 TeV, while consistently reproducing the observed values of the gauge coupling constants at the electroweak scale, and achieving the values required by the string theory at the string scale. However, our models *do* have extra light, vector-like lepton states (but not quark states), and Higgs doublets, as well as the towers of gonions characteristic of all theories of this type. There are also three extra non-anomalous $U(1)$ gauge groups under which the matter is charged. One problem that is common to all such intersecting brane models is that the complex structure moduli of the wrapped torus is not stabilised. (The dilaton and Kähler moduli are not stabilised either, but that is not peculiar to these models.) The complex structure moduli *are* stabilised if the wrapped space is an orbifold or an orientifold [25,26], and more exotic possibilities that stabilise more moduli have also been proposed [27,28].

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