

More about the Standard Model at Intersecting Branes

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Abstract

Intersecting D-brane models seem to be one of the most promising avenues to embed the standard model physics within the string context. We review here different aspects of these models. Topics include the question of SUSY and quasi-SUSY in intersecting brane models, model-building, the brane recombination interpretation of the SM Higgs mechanism, Yukawa couplings, the lowering of the string scale and possible new Z 's accessible to accelerators.

In the last couple of years there have been renovated efforts in looking for D-brane configurations with a low-energy effective theory resembling the standard model (SM). One approach which looks particularly successful is that of intersecting D-brane models [1, 2, 3, 4, 5, 6] (see also [7, 8, 9, 10, 11, 12, 13, 14, 15, 16]). We will not attempt to give here an introduction to the subject (see e.g. [6] and references therein). Rather we will concentrate on giving a brief report on some of the work in the subject that we have been involved with in the last year.

In intersecting brane models (fig.1) the different gauge interactions live on different stacks of D-branes, the simplest configurations having four stacks: *baryonic*, *left*, *right* and *leptonic*. In particular one considers stacks of branes with multiplicities $N_a = 3$, $N_b = 2$, $N_c = 1$, $N_d = 1$, yielding initially a gauge group $U(3) \times U(2) \times U(1) \times U(1)$. Up to three of the $U(1)$'s become massive by combining with some closed string (Ramond-Ramond) fields so that in the simplest situation one is just left with standard hypercharge and the full group is that of the SM. The D_p -branes (with $p=6,5,4$) worldvolumes contain Minkowski space and the remaining $(p-3)$ dimensions wrap cycles on compact dimensions. At the brane intersections in extra dimensions live quarks and leptons and the triplication of generations appear because in the compact space the different branes intersect three times. In particular, if we denote by I_{ab} the number of times that branes (a) and (b) intersect, the following intersection numbers [5]:

$$\begin{aligned} I_{ab} &= 1 ; I_{ab^*} = 2 \\ I_{ac} &= -3 ; I_{ac^*} = -3 \\ I_{bd} &= 0 ; I_{bd^*} = -3 \\ I_{cd} &= -3 ; I_{cd^*} = 3 \end{aligned} \tag{1}$$

give rise to the fermion spectrum of the SM. Specific D6-brane models in which the compact space is just a 6-torus T^6 and yielding the above SM spectrum were provided in ref.[5]. One can also find D5-brane models in which the compact space is $T^4 \times (T^2/Z_N)$ and

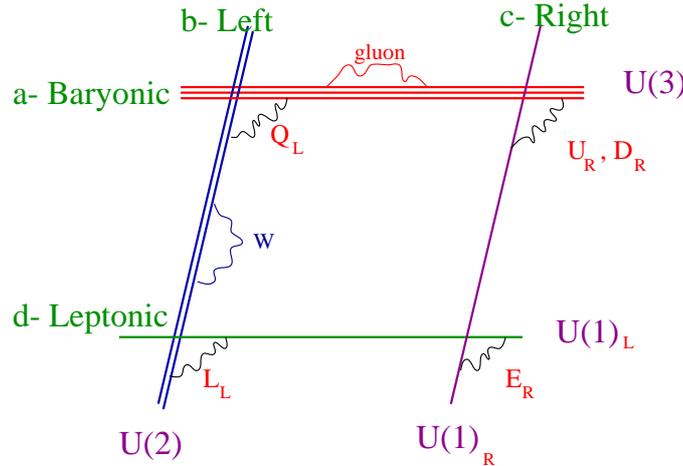


Figure 1: The SM spectrum at intersecting branes.

one obtains the same SM chiral fermion spectrum [15]. In both classes of constructions there is an interesting connection between the number of generations and the number of colours. Indeed, in order to cancel anomalies the net number of $U(2)_b$ doublets has to equal that of anti-doublets [5], which in these models happens only because the number of generations equals the number of colours. In addition, one of the nicest features of these constructions is that the proton is automatically stable since baryon number ($U(1)_a$) is a gauged symmetry [5].

We must emphasize that the intersecting brane setting just described is quite general. As fig.(2) illustrates, one may consider for example four stacks a,b,c,d of D6-branes wrapping 3-cycles on a complicated Calabi–Yau. As long as the intersection numbers are as above, the chiral fermion spectrum is going to be the one of the SM independently of the details of the compactification. As an example the above SM spectrum has also been recently obtained from D6-branes wrapping 3-cycles on the quintic CY [17] or in certain non-compact manifolds [18].

The mentioned specific toroidal models are generically non-SUSY due to the presence of the intersections. Thus, for example, associated to each of the intersections there are massive scalar fields which in some sense may be considered "SUSY-partners", squarks and sleptons, of the massless chiral fermions, have the same multiplicity $|I_{ij}|$ and carry the same gauge quantum numbers. The lightest of those states have masses [2]

$$\begin{aligned}
 (\text{Mass})_1^2 &= \frac{1}{2}(-\vartheta_1 + \vartheta_2 + \vartheta_3)M_s^2 \\
 (\text{Mass})_2^2 &= \frac{1}{2}(\vartheta_1 - \vartheta_2 + \vartheta_3)M_s^2 \\
 (\text{Mass})_3^2 &= \frac{1}{2}(\vartheta_1 + \vartheta_2 - \vartheta_3)M_s^2
 \end{aligned}
 \tag{2}$$

where ϑ_i are the intersection angles (in units of π) at each of the three subtori. As is obvious from these formulae the masses depend on the angles at each intersection and hence on the relative size of the radii. In principle some of the scalars could be tachyonic, but in general it is possible to vary the compact radii in order to get rid of all tachyons of a given model (see [2, 5]). On the other hand, one can also adjust the radii so that there is one massless scalar at the intersection. Then one gets $N = 1$ SUSY *at that*

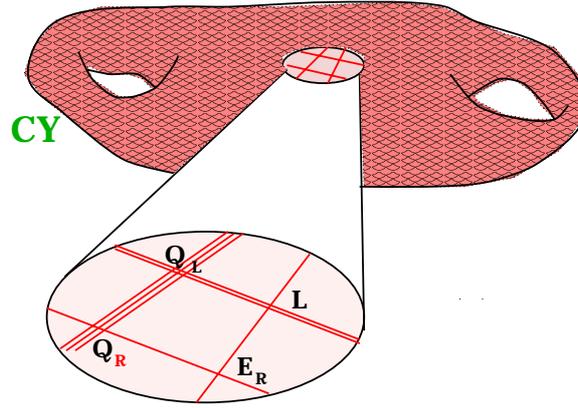


Figure 2: The general intersecting brane system may be embedded i.e., into a general Calabi-Yau compact space.

specific intersection. Is it possible to get a fully $N = 1$ SUSY model, i.e., a model in which all intersections respect the *same* $N = 1$ supersymmetry? The answer is no, at least in the purely toroidal examples as in ref.[5]. The reason for this is that D-brane configurations wrapping compact spaces as in here have to respect the conditions of RR-tadpole cancellation. The overall charge of the configuration with respect to certain tensorial RR-fields has to vanish. In the purely toroidal case those conditions turn out to be incompatible with the geometrical configurations required to get $N = 1$ SUSY. Fully $N = 1$ SUSY intersecting brane models may be built in the case with an added $Z_2 \times Z_2$ orbifold twist [8]. However in that case additional chiral exotics beyond the SM content seem unavoidable.

On the other hand there is an interesting possibility termed quasi-SUSY in ref.[10] and pseudo-SUSY in ref.[19]. The possibility exists that all intersections respect some $N = 1$ SUSY but *different ones*. In the case of toroidal models this possibility is still compatible with cancellation of RR tadpoles. For example, certain subset of the models in ref.[5] can be made quasi-SUSY by choosing appropriate radii. Let us denote by (n_a^i, m_a^i) , $i=1,2,3$ the wrapping numbers of each brane $D6_a$, $n_a^i(m_a^i)$ being the number of times the brane is wrapping around the $x(y)$ -coordinate of the i -th torus. Consider in particular the wrapping numbers for the different branes given in table 1. The number of times the two branes $D6_a$ and $D6_b$ intersect in T^6 is given by the intersection number [1] $I_{ab} = (n_a^1 m_b^1 - m_a^1 n_b^1)(n_a^2 m_b^2 - m_a^2 n_b^2)(n_a^3 m_b^3 - m_a^3 n_b^3)$. Then one can easily check that this brane setting yields the chiral spectrum of the SM. Now, the masses of the scalars depend on the ratios $U^i = R_2^i / R_1^i$, $i = 1, 2, 3$. One can easily check that if we set $U^1 = (n_b^1/2)U^3$ and $U^2 = (n_a^2/6\beta^2)U^3$ massless scalars appear at each intersection and some $N = 1$ SUSY is preserved at each of them. This property may be depicted in terms of a *square quiver diagram* which is shown in fig.(3). Thus now each quark and lepton has a massless SUSY-partner, very much as in the SUSY-SM. The model is however not fully $N = 1$ supersymmetric because each intersection respects a different SUSY. This kind of *quasi-SUSY* theories have some interesting properties. In particular, loop corrections to scalar masses appear only at two loops, since only at that order the global non-SUSY

N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
$N_a = 3$	$(1, 0)$	(n_a^2, β^2)	$(3, -1/2)$
$N_b = 2$	$(n_b^1, 1)$	$(1/\beta^2, 0)$	$(1, -1/2)$
$N_c = 1$	$(0, 1)$	$(1/\beta^2, 0)$	$(0, 1)$
$N_d = 1$	$(1, 0)$	$(n_a^2, 3\beta^2)$	$(1, 1/2)$
N_h	$(1, 0)$	$(1/\beta^2, 0)$	(n_h^3, m_h^3)

Table 1: D6-brane wrapping numbers giving rise to a Q-SUSY SM spectrum for a square quiver. For the sake of generality we have also considered the possible presence of an extra brane with no intersection with the SM branes.

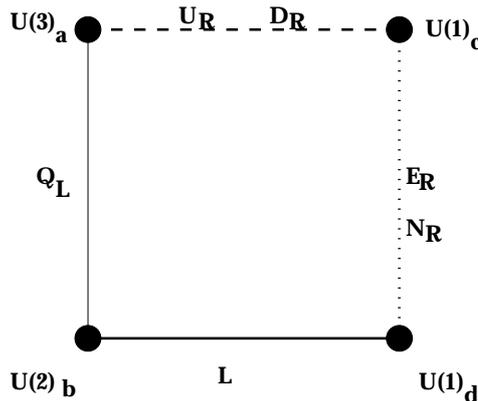


Figure 3: A square SUSY-quiver.

structure of the configuration may be noticed [10, 13, 19]. This loop suppression of scalar masses may be interesting phenomenologically in order to address the “modest hierarchy problem”, i.e., in order to maintain a hierarchy between a string scale of order 10-100 TeV and the weak scale.

There is a variety of SUSY-quivers that one may consider leading to different low-energy models [13]. One can also find D6-brane configurations wrapping T^6 leading to the massless chiral spectrum of the MSSM. Some examples were presented in ref.[10, 13] but we will present here a new and remarkably simple model which will be discussed in more detail elsewhere [20]. Consider the stacks of D6-branes with the wrapping numbers of table 2. Generically the gauge group of this configuration is $U(3) \times U(1)^3$. However, one can check that the symmetry is enhanced to $U(3)_a \times SU(2)_b \times U(1)_c \times U(1)_d$ if the brane (b) is located on top of its orientifold mirror (b^*). Computing the intersection numbers as

N_i	(n_i^1, m_i^1)	(n_i^2, m_i^2)	(n_i^3, m_i^3)
$N_a = 3$	$(1, 0)$	$(1/\rho, 3\rho)$	$(1/\rho, -3\rho)$
$N_b = 1$	$(0, 1)$	$(1, 0)$	$(0, -1)$
$N_c = 1$	$(0, 1)$	$(0, -1)$	$(1, 0)$
$N_d = 1$	$(1, 0)$	$(1/\rho, 3\rho)$	$(1/\rho, -3\rho)$

Table 2: D6-brane wrapping numbers giving rise to a the chiral spectrum of the MSSM. Here $\rho = 1, 1/3$.

above one gets the result

$$\begin{aligned}
I_{ab} &= 3 ; I_{ab^*} = 3 \\
I_{ac} &= -3 ; I_{ac^*} = -3 \\
I_{db} &= 3 ; I_{db^*} = 3 \\
I_{dc} &= -3 ; I_{dc^*} = 3 \\
I_{bc} &= -1 ; I_{bc^*} = 1
\end{aligned} \tag{3}$$

which corresponds to the chiral fermion spectrum of the SM (plus right-handed neutrinos). In addition there is a minimal set of Higgs multiplets if one locates the brane (b) on top of the brane (c) along the first torus. In other words, there is a minimal Higgs sector with a μ -parameter given by the distance between branes (b) and (c) along the first torus. If the ratios of radii in the second and third torus are equal ($U^2 = U^3 = \chi$) one can check that *the same* $N=1$ SUSY is preserved at all intersections. So this configuration is (locally) $N = 1$ supersymmetric, and *the massless chiral spectrum is that of the MSSM with a minimal Higgs set*. The spectrum is anomaly-free in the sense that there are as many fundamentals as antifundamentals of any of the groups. On the other hand the configuration cannot be made fully $N = 1$ supersymmetric because it turns out that in order to cancel RR-tadpoles an additional massive $N = 0$ sector has to be added (see ref.[13] for a discussion of this point). In this model there are three $U(1)$'s and only one of them ($3B + L$) is anomalous and gets massive by combining with one RR-field. There are two massless $U(1)$'s corresponding to $(B - L)$ and the 3-d component of right-handed weak isospin ($U(1)_c$). So the actual low-energy gauge group is $SU(3) \times SU(2) \times U(1)_{B-L} \times U(1)_c$.

One of the nice features of the intersecting brane approach is that the low-energy Lagrangian parameters admit a simple geometrical interpretation. We already saw an example: the μ -parameter in this model corresponds to the distance between branes (b) and (c) in the first torus. Another example is the generation of tree level Fayet-Iliopoulos terms for the anomalous $U(1)$'s [8, 10]. If one has a small departure from the SUSY geometry, i.e., if $U^2 = U^3 + \delta$ with δ small, one finds a FI-term for the anomalous $U(1)$ [10] :

$$\xi = M_s^2 \times \frac{3\rho^2 \delta}{1 + (3\rho^2 U^3)^2} . \tag{4}$$

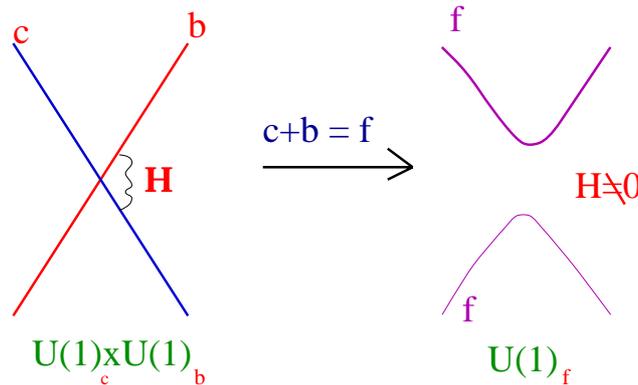


Figure 4: Branes (b) and (c) are recombined into a single brane f. The gauge symmetry is reduced.

If $\delta \neq 0$, the existence of this FI-term may trigger further gauge symmetry breaking. In particular the additional $U(1)$ may be broken down to standard hypercharge by inducing a vev to the right-handed sneutrino.

Before taking into account SUSY-breaking effects, the above local SUSY configuration has (for vanishing μ -term) a flat direction in which electroweak symmetry is broken by the vevs of the Higgs scalars lying at the (bc) and (bc^*) intersections H_{bc} . The Higgs mechanism has also a geometric interpretation in terms of branes [13]. As illustrated in fig.(4), a vev for the Higgs fields H_{bc} corresponds to a process $b + c \rightarrow f$ in which branes (b) and (c) recombine into a single brane f. Since we have started from two branes (plus orientifold mirrors) and end up with one brane and its mirror, the rank of the gauge group has been reduced. Altogether we are left at the end of the day with only three brane stacks, a,d,f and a gauge group $SU(3) \times U(1)_{em}$ at the massless level.

The Yukawa couplings also have an interesting geometrical interpretation in intersecting brane models [3]. A Yukawa coupling involves correlators of a Higgs field, a right-handed fermion and a left-handed fermion. The worldsheet of strings connecting those three vertices has a triangular shape, as in fig.(5), with open strings stretching between the three intersecting branes participating in the coupling. The Yukawa couplings are then proportional to $\exp(-S_{cl})$, S_{cl} being the classical string action, which is proportional to the area of the worldsheet. This provides a nice physical origin for the observed hierarchy of fermion masses since, as exemplified in fig.(5), the size of the relevant triangles for the different generations is in general different [3]. Thus, e.g., the triangle associated to the top-quark coupling would be smaller than the one associated to the c-quark which would be smaller than the one of the u-quark. One can also see that generation mixing as well as complex phases do in general appear. A systematic analysis of Yukawa couplings in intersecting brane models will appear in ref.[20].

One interesting question is whether in this class of intersecting brane models one can realize the low string scale scenario [21] with $M_s \sim 1-10$ TeV. This is particularly relevant in models which are not supersymmetric and in which lowering the string scale down to the TeV scale provides then a solution to the hierarchy problem. As is well known, this requires that at least 2 of the 6 compact dimensions become very large, so that we obtain

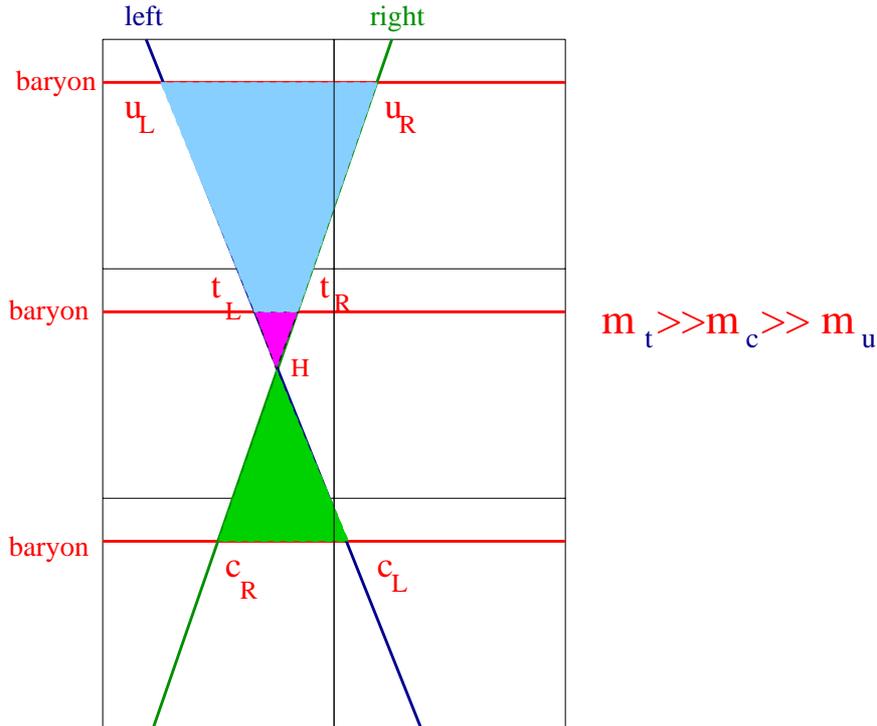


Figure 5: The Yukawa couplings are exponentially suppressed by the area of the triangular worldsheets. Thus large hierarchies originate from the different areas corresponding to different generations.

a large splitting between the string scale M_s and the effective 4-dimensional gravity scale M_{Planck} . In the simple case of D6-branes wrapping a 6-torus realizing the low string scale scenario is in principle complicated [1]. This is because there are no dimensions which can be made large and are orthogonal to the SM brane system. On the other hand, as pointed out in [2] and explained in more detail in [18], one can start from a toroidal model as above and construct a related model in which the compact volume can be made arbitrarily large without affecting the SM branes. This is done by gluing an infinite throat to the torus in a region far away from the branes and then connecting the throat to some large volume. Alternatively one can consider other intersecting brane constructions in which the presence of a couple of dimensions transverse to the SM branes which can be made large is more obvious. For example, one can consider Type IIB compactified on $T^2 \times T^2 \times (T^2/Z_N)$ with D5-branes wrapping 2-cycles on $T^2 \times T^2$ and located at a fixed point of the orbifold T^2/Z_N (see fig.6). One can now have a low string scale $M_s \propto 1 - 10$ TeV while maintaining the experimentally measured four-dimensional Planck mass $M_p = 1.18 \times 10^{19}$ GeV by some dimensions getting very large [21]. Let us denote by V_4 the volume of T^4 and by V_2 that of T^2/Z_N . Then the Planck scale is given by

$$M_p = \frac{2}{\lambda} M_s^4 \sqrt{V_4 V_2} \quad (5)$$

In order to avoid too light KK/Winding modes in the worldvolume of the D5-branes let

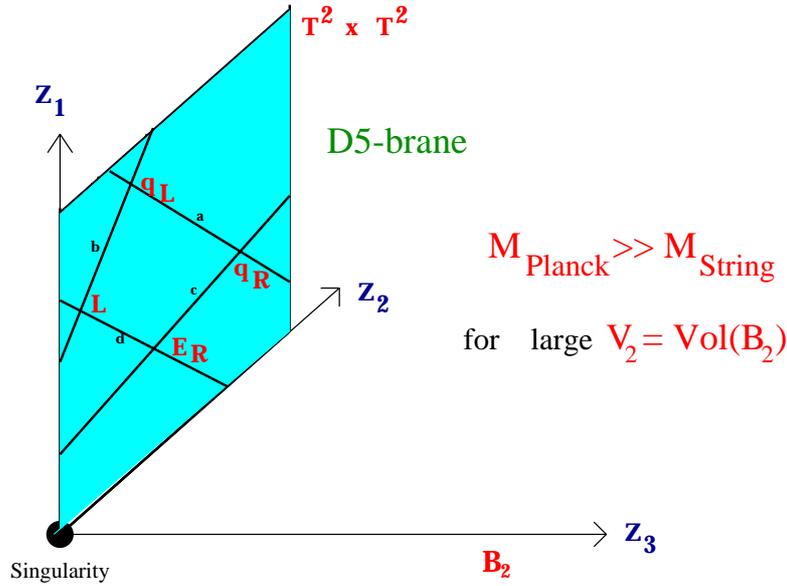


Figure 6: Intersecting D5-world set up. One can obtain a low string scale $M_s \ll M_p$ if the volume of the two transverse dimensions is large.

us assume $V_4 \propto 1/M_s^4$. Then one has

$$V_2 = \frac{M_p^2 \lambda^2}{4M_s^4} \quad (6)$$

and one can accommodate a low string scale $M_s \propto 1$ TeV by having the volume V_2 of the 2-dimensional manifold \mathbf{B}_2 large enough (i.e., of order $(mm.)^2$). Such D5-brane models with the chiral fermion content as in eq.(1) leading to the fermions of the SM have been recently constructed in ref.[15]. Note however that those D5-brane constructions are intrinsically non-SUSY since the Z_N orbifold projection projects out all gauginos.

In the case of a string scale M_s close to the 1-10 TeV range, one interesting feature of the intersecting brane models is the presence of a very well defined and model-independent class of extra TeV-scale Z' bosons. Indeed, in all models there are some extra $U(1)$ symmetries beyond hypercharge which seem rather model independent [22]. We have generators $U(1)_a$ and $U(1)_d$ which are gauged baryon and lepton numbers respectively. $U(1)_c$ correspond to the 3-d component of right-handed weak isospin and $U(1)_b$ is a PQ-like gauged symmetry. In the class of models in eq.(3) the latter $U(1)_b$ symmetry is absent. Hypercharge is a linear combination of $(B - L)$ and $U(1)_c$. The orthogonal $U(1)$'s may get Stueckelberg masses ¹ by combining with RR string fields $B_i^{\mu\nu}$, as shown in fig.(7). In this way one gets a mass matrix for the Abelian gauge bosons of the form:

$$M_{\alpha\beta}^2 = \sum_i g_\alpha g_\beta c_i^\alpha c_i^\beta M_s^2 \quad (7)$$

where α, β run over the $U(1)$ factors of each specific model and the c_i^α are model-

¹Note that all anomalous $U(1)$'s get masses through this mechanism but also some anomaly-free $U(1)$'s may get a mass, see ref.[5].

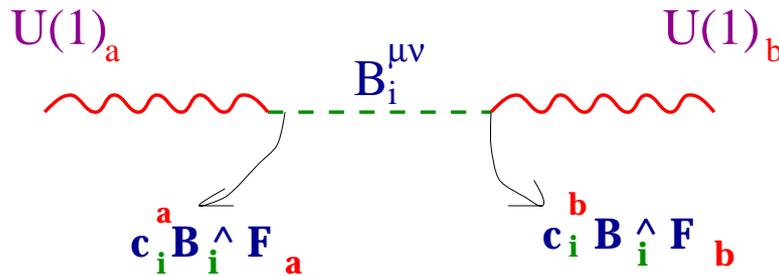


Figure 7: Mixing of $U(1)$'s and antisymmetric RR-fields $B_i^{\mu\nu}$ yielding masses to Abelian gauge fields.

dependent coefficients which may be computed in each particular brane setting [22, 23]. Thus for example, in the class of models discussed in ref.[5] or ref.[15] one sees that this matrix has four eigenvalues $M = (0, M_2, M_3, M_4)$, with the zero mode corresponding to standard hypercharge. It turns out that in those models typically one of the eigenvalues is well below the string scale M_s , so that one could detect the effects of such extra $U(1)$ before actually reaching the string threshold. One can also put constraints on those eigenvalues from ρ -parameter bounds [22]. It would be rather amusing if the first signature of string physics would come from the detection of any of these particular extra Z' bosons.

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