STANDARD-LIKE MODELS FROM INTERSECTING D5-BRANES

David BAILIN♣ [1], George V. KRANIOTIS♠ [2], and A. LOVE♢

♣ *Centre for Theoretical Physics, University of Sussex*
Brighton BN1 9QJ, U.K.
♠ *Fachbereich Physik, Martin-Luther-Universität Halle-Wittenberg*
Friedemann-Bach-Platz 6, D-06099, Halle, Germany
♢ *Centre for Particle Physics, Royal Holloway, University of London*
Egham, Surrey TW20-0EX, U.K.

ABSTRACT

We describe constructions of intersecting D5-brane orbifold models that yield the (non-supersymmetric) standard model up to vector-like matter and charged-singlet scalars. The models are constrained by the requirement that twisted tadpoles cancel, and that the gauge boson coupled to the weak hypercharge $U(1)_Y$ does not get a string-scale mass via a generalised Green-Schwarz mechanism. Gauge coupling constant ratios close to those measured are easily obtained for reasonable values of the parameters, consistently with having the string scale close to the electroweak scale, as required to avoid the hierarchy problem.

The D-brane world offers an attractive, bottom-up route to getting standard-like models from Type II string theory [1]. Open strings that begin and end on a stack of M D-branes generate the gauge bosons of the group $U(M)$ living in the world volume of the D-branes. So the standard approach is to start with one stack of 3 D-branes, another of 2, and n other stacks each having just 1 D-brane, thereby generating the gauge group $U(3) \times U(2) \times U(1)^n$. Fermions in bi-fundamental representations of the corresponding gauge groups can arise at the intersections of such stacks [2], but to get $D = 4$ *chiral* fermions the intersecting branes should sit at a singular point in the space transverse to the branes, an orbifold fixed point, for example. In general, such configurations yield a non-supersymmetric spectrum, so to avoid the hierarchy problem the string scale associated with such models must be no more than a few TeV.

Gravitational interactions occur in the bulk ten-dimensional space, and to ensure that the Planck energy has its observed large value, it is necessary that there are large dimensions transverse to the branes [3]. The D-branes with which we are concerned wrap the 3-space we inhabit and closed 1-, 2- or 3-cycles of a toroidally compactified T^2, $T^2 \times T^2$ or $T^2 \times T^2 \times T^2$ space. Thus getting the correct Planck scale effectively means that only

[1] D.Bailin@sussex.ac.uk
[2] kraniotis@physik.uni-halle.de.

D4- and D5-brane models are viable, since for D6-branes there is no dimension transverse to all of the intersecting branes. In a non-supersymmetric theory the cancellation of the closed-string (twisted) Ramond-Ramond (RR) tadpoles does *not* ensure the cancellation of the Neveu-Schwarz-Neveu-Schwarz (NSNS) tadpoles. There is a resulting instability in the complex structure moduli [4]. One way to stabilise some of the (complex structure) moduli is to use an orbifold, rather than a torus, for the space wrapped by the D-branes. The above instability has also recently inspired attempts for reconciling string theory with inflationary cosmology [5].

If the embedding is supersymmetric, then the instabilities are removed. This has been studied [7], using D6-branes, but it has so far proved difficult to get realistic phenomenology consistent with experimental data from such models [3].

During the past year orientifold models with intersecting D6- and D5-branes have been constructed that yield precisely the fermionic spectrum of the standard model (plus three generations of right-chiral neutrinos) [8]. The spectrum includes $SU(2)_L$ doublet scalar tachyons that may be regarded as the Higgs doublets that break the electroweak symmetry group, but also, unavoidably, colour-triplet and charged singlet tachyons either of which is potentially fatal for the phenomenology. The wrapping numbers of the various stacks are constrained by the requirement of RR tadpole cancellation, and this ensures the absence of non-Abelian anomalies in the emergent low-energy quantum field theory. A generalised Green-Schwarz (GS) mechanism ensures that that the gauge bosons associated with all anomalous $U(1)$s acquire string-scale masses [9], but the gauge bosons of some non-anomalous $U(1)$s can also acquire string-scale masses [8]; in all such cases the $U(1)$ group survives as a global symmetry. Thus we must also ensure the weak hypercharge group $U(1)_Y$ remains a gauge symmetry by requiring that its gauge boson does *not* get such a mass. In a recent paper [10] we constructed the first semi-realistic intersecting D4-brane orbifold models that satisfy these constraints.

We start with an array of D5-brane stacks, each wrapping closed 1-cycles in both 2-tori of $T^2 \times T^2$, and situated at a fixed point of the transverse $T^2/\mathbb{Z}_3$ orbifold. A stack a is specified by two pairs of wrapping numbers (n_a, m_a) and $(\tilde{n}_a, \tilde{m}_a)$ that specify the number of times a wraps the basis 1-cycles in each T^2. When n_a and m_a are coprime a single copy of the gauge group $U(N_a)$ occurs; if n_a and m_a have a highest common factor f_a there are f_a copies of $U(N_a)$, or if m_a (or n_a) is zero, then there are n_a (or m_a) copies of $U(N_a)$. The same thing happens if $\tilde{n}_a$ and $\tilde{m}_a$ are not coprime. The number of intersections of stack a with stack b is given by $\mathbf{I}_{ab} = I_{ab}\tilde{I}_{ab}$, where $I_{ab} \equiv n_a m_b - m_a n_b$ and $\tilde{I}_{ab} = \tilde{n}_a \tilde{m}_b - \tilde{m}_a \tilde{n}_b$ are the intersection numbers of the corresponding 1-cycles in each torus T^2. The generator θ of the Z_3 point group is embedded in the stack of N_a branes as $\gamma_{\theta,a} = \alpha^{p_a} I_{N_a}$, where $\alpha = e^{2\pi i/3}$, $p_a = 0, 1, 2$. The first two stacks $a = 1, 2$ defined above, that generate a $U(3) \times U(2)$ gauge group, are common to all models. We require that the intersection of these two stacks gives the three copies of the left-chiral quark doublet Q_L. Then, without loss of generality, we may take their wrapping numbers as shown in Table 1. We require that $n_2 \neq 0 \bmod 3$, so that the $U(2)$ group is not replicated. Besides the first two stacks we have, in general, three sets I, J, K of $U(1)$ stacks characterised by their Chan-Paton factors: $p_i = p_2\ \forall i \in I$, $p_1 \neq p_j \neq p_2\ \forall j \in J$, and $p_k = p_1\ \forall k \in K$;

[3] Another promising way for stabilising the moduli is to turn on fluxes for antisymmetric fields of type II-string theory. In this case a non-trivial superpotential is generated, and depending on the choice of fluxes various scenarios for expaining the hierarchy problem can in principle be obtained [6].

Stack a	N_a	(n_a, m_a)	$(\tilde{n}_a, \tilde{m}_a)$	$\gamma_{\theta,a}$
1	3	$(1,0)$	$(1,0)$	$\alpha^p \mathbf{I}_3$
2	2	$(n_2, 3)$	$(\tilde{n}_2, 1)$	$\alpha^q \mathbf{I}_2$
$i \in I$	1	(n_i, m_i)	$(\tilde{n}_i, \tilde{m}_i)$	α^q
$j \in J$	1	(n_j, m_j)	$(\tilde{n}_j, \tilde{m}_j)$	α^r
$k \in K$	1	(n_k, m_k)	$(\tilde{n}_k, \tilde{m}_k)$	α^p

Table 1: Multiplicities, wrapping numbers and Chan-Paton phases for the D5-brane models, $(p \neq q \neq r \neq p)$.

the sets I, J, K are each divided into two subsets $I_1 \cup I_2 = I$ etc., defined [11] so that the weak hypercharge Y is the linear combination

$$-Y = \frac{1}{3}Q_1 + \frac{1}{2}Q_2 + \sum_{i_1 \in I_1} Q_{i_1} + \sum_{j_1 \in J_1} Q_{j_1} + \sum_{k_1 \in K_1} Q_{k_1} \tag{1}$$

where Q_a is the $U(1)$ charge associated with the stack a; Q_a is normalised such that the $\mathbf{N}_a$ representation of $SU(N_a)$ has $Q_a = +1$.

The mixed $U(1)$ anomalies are cancelled by a generalised GS mechanism involving the exchange of fields that arise from the dimensional reduction of the twisted two-form RR fields $B_2^{(k)}$, $C_2^{(k)}$, $D_2^{(k)}$ and $E_2^{(k)}$ that live at the orbifold singularity [12]. These fields are coupled to the $U(1)_a$ field strength F_a of the stack by terms in the low energy action of the form

$$\begin{aligned} n_a\tilde{n}_a \int_{M_4} \mathrm{Tr}(\gamma_{k,a}\lambda_a) B_2^{(k)} \wedge \mathrm{Tr}F_a, \quad m_a\tilde{m}_a \int_{M_4} \mathrm{Tr}(\gamma_{k,a}\lambda_a) C_2^{(k)} \wedge \mathrm{Tr}F_a, \\ m_a\tilde{n}_a \int_{M_4} \mathrm{Tr}(\gamma_{k,a}\lambda_a) D_2^{(k)} \wedge \mathrm{Tr}F_a, \quad n_a\tilde{m}_a \int_{M_4} \mathrm{Tr}(\gamma_{k,a}\lambda_a) E_2^{(k)} \wedge \mathrm{Tr}F_a \end{aligned} \tag{2}$$

where $\gamma_{k,a} \equiv \gamma_{\theta,a}^k$ and λ_a is the CThe couplings determine the linear combination of $U(1)$ gauge bosons that acquire string-scale masses via the GS mechanism. We require that the $U(1)_Y$ gauge boson associated with the weak hypercharge given in eqn (1) remains massless. Consequently, the corresponding field strength must be orthogonal to those that acquire GS masses. It follows that the sets $J_1 \cup K_1$ and $J_2 \cup K_2$ are both non-empty.

We obtain the following constraints on the wrapping numbers deriving from twisted tadpole cancellation and requiring that the $U(1)_Y$ gauge boson associated with weak hypercharge does not acquire a GS mass:

$$n_2\tilde{n}_2 + \sum_{i_1} n_{i_1}\tilde{n}_{i_1} = \sum_{j_1} n_{j_1}\tilde{n}_{j_1} = 1 + \sum_{k_1} n_{k_1}\tilde{n}_{k_1} \equiv T_{11} \tag{3}$$

$$n_2\tilde{n}_2 + \sum_{i_2} n_{i_2}\tilde{n}_{i_2} = \sum_{j_2} n_{j_2}\tilde{n}_{j_2} = 2 + \sum_{k_2} n_{k_2}\tilde{n}_{k_2} \equiv T_{12} \tag{4}$$

$$3 + \sum_{i_1} m_{i_1}\tilde{m}_{i_1} = \sum_{j_1} m_{j_1}\tilde{m}_{j_1} = \sum_{k_1} m_{k_1}\tilde{m}_{k_1} \equiv T_{41} \tag{5}$$

$$3 + \sum_{i_2} m_{i_2}\tilde{m}_{i_2} = \sum_{j_2} m_{j_2}\tilde{m}_{j_2} = \sum_{k_2} m_{k_2}\tilde{m}_{k_2} \equiv T_{42} \tag{6}$$

$$n_2 + \sum_{i_1} n_{i_1} \tilde{m}_{i_1} = \sum_{j_1} n_{j_1} \tilde{m}_{j_1} = \sum_{k_1} n_{k_1} \tilde{m}_{k_1} \equiv T_{21} \tag{7}$$

$$n_2 + \sum_{i_2} n_{i_2} \tilde{m}_{i_2} = \sum_{j_2} n_{j_2} \tilde{m}_{j_2} = \sum_{k_2} n_{k_2} \tilde{m}_{k_2} \equiv T_{22} \tag{8}$$

$$3\tilde{n}_2 + \sum_{i_1} m_{i_1} \tilde{n}_{i_1} = \sum_{j_1} m_{j_1} \tilde{n}_{j_1} = \sum_{k_1} m_{k_1} \tilde{n}_{k_1} \equiv T_{31} \tag{9}$$

$$3\tilde{n}_2 + \sum_{i_2} m_{i_2} \tilde{n}_{i_2} = \sum_{j_2} m_{j_2} \tilde{n}_{j_2} = \sum_{k_2} m_{k_2} \tilde{n}_{k_2} \equiv T_{32} \tag{10}$$

Tachyonic scalars arise only at intersections between stacks a and b which have the same Chan-Paton factor $p_a = p_b$. Thus, Higgs doublets, which are needed to give mass to the fermionic matter, arise at $(2i_1)$ and $(2i_2)$ intersections. With D5-branes, chiral fermions arise only at intersections with $p_a \neq p_b$. Thus to obtain (doublet and singlet) leptons, we also require $U(1)$ stacks in the sets J and/or K.

The difference of the angles θ_{ab} and $\tilde{\theta}_{ab}$ between the intersecting 1-cycles on the two wrapped tori determine the squared mass of the tachyon at such an intersection, and when they are equal the tachyon is massless. Thus this mechanism in D5-brane models may remove potentially unwelcome (e.g. charged-singlet and/or colour triplet) tachyons. A further novelty compare to D4-brane models is that tachyons may also arise when the intersecting D5-branes have parallel 1-cycles in one, but not both, of the two wrapped tori, i.e. $\theta_{ab} = 0$ *or* $\tilde{\theta}_{ab} = 0$. The intersection number of the non-parallel 1-cycles determines the number of such states that arise. The distance between the parallel 1-cycles enters the mass formula, so the masses of these states are continuously adjustable. When it is zero (or small enough) the states are tachyonic, but when the cycles are far enough apart the states are massive and can be removed from the low-energy spectrum. Equations (3)-(10) are quadratic in the wrapping numbers, and the only six-stack models that can satisfy them have stacks $1, 2, i_1, i_2, k_1, k_2$. Unfortunately, the solution gives $\mathbf{I}_{2i_1} = \mathbf{I}_{2i_2} = 0$ and, since the 1-cycles in both tori are parallel, these models have no Higgs doublets (or charged-singlet or colour triplet tachyons). They are therefore not standard-like and we are forced to consider eight-stack models from the outset.

We concentrate in models with not more than one stack in each of the classes $I_1, I_2, J_1, J_2, K_1, K_2$. To ensure that Higgs doublets arise at the $(2i_1)$ intersection, either T_{21} or T_{31}, must be non-zero. Similarly, for the $(2i_2)$ intersection, either T_{22} or T_{32}, must be non-zero. Then the general solution for the wrapping numbers are presented in tables 2 and 3. To avoid triplication of gauge group factors in table 2, we require that $n_2 = -p \bmod 3$ and to avoid triplication of gauge group factors in table 3 we require $n_2 = q \bmod 3$ for $T_{22} = 0, n_2 = -q \bmod 3$ for $T_{32} = 0$.

There are four classes of models obtained by combining one of the options in Table 2 with one of the options in Table 3. For no choice of the parameters can any of the stacks in Table 2 or 3 be absent. Thus a minimum of eight stacks is required, and in fact this a general result, even if we allow more than one stack in some of the sets [13]. In the first instance all such eight-stack models have gauge group $U(3) \times U(2) \times U(1)^6$. By construction the standard model gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$ survives the GS mechanism that gives string-scale masses to some of the original (eight) $U(1)$ gauge bosons. In addition, it is easy to see that the symmetry $U(1)_X$, associated with the the sum of the charges $X = \sum_a Q_a$ also survives as a gauge symmetry. However, this is uncoupled to all of the matter and gauge fields, and so is physically unobservable.

Stack a	n_a	m_a	$\tilde{n}_a$	$\tilde{m}_a$
i_1	n_2	3	$p-\tilde{n}_2$	-1
j_1	n_2p	$3p$	1	0
k_1	n_2p-1	$3p$	1	0
i_1	n_2-p	3	$-\tilde{n}_2$	-1
j_1	1	0	$\tilde{n}_2p$	p
k_1	1	0	$\tilde{n}_2p-1$	p

Table 2: Wrapping numbers for the stacks i_1, j_1 and k_1. At the top $T_{21}=0,\ T_{31}=3p,\ T_{11}=n_2p$, and at the bottom $T_{21}=p,\ T_{31}=0,\ T_{11}=\tilde{n}_2p$. In both cases $p=\pm1$. A further overall, arbirary sign $\epsilon_a=\pm1$ is understood for each stack $a=i_1,j_1,k_1$.

Stack a	n_a	m_a	$\tilde{n}_a$	$\tilde{m}_a$
i_2	n_2	3	$q-\tilde{n}_2$	-1
j_2	n_2q	$3q$	1	0
k_2	n_2q-2	$3q$	1	0
i_2	$q-n_2$	-3	$\tilde{n}_2$	1
j_2	1	0	$\tilde{n}_2q$	q
k_2	1	0	$\tilde{n}_2q-2$	q

Table 3: Wrapping numbers for the stacks i_2, j_2 and k_2. At the top $T_{22}=0,\ T_{32}=3q,\ T_{12}=n_2q$, and at the bottom $T_{22}=q,\ T_{32}=0,\ T_{12}=\tilde{n}_2q$. In both cases $q=\pm1$. A further arbitrary, overall sign $\epsilon_a=\pm1$ is understood for each stack $a=i_2,j_2,k_2$.

It turns out that in the case $T_{21}=0=T_{32}$, 6 of the original 8 $U(1)$ gauge bosons get GS masses, leaving just the massless, gauged $U(1)_X$ and $U(1)_Y$. Potentially this can yield a standard-like model. The consistency conditions for avoiding triplication requires $n_2=-p \bmod 3=-q \bmod 3$ The spectrum, which is independent of p, has 3 generations of chiral matter that include right-chiral neutrino states, and there is additional vector-like leptonic, but not quark, matter. We find

$$3(L+\bar{L})+6(e^c_L+\bar{e}^c_L)+3(\nu^c_L+\bar{\nu}^c_L) \tag{11}$$

There is also one Higgs doublet at the $(2i_1)$ intersections, and 3 Higgs doublets at the $(2i_2)$ intersections. Less welcome are the 9 charged-singlet tachyons that arise, 3 at each of the (i_1i_2), (j_1j_2) and (k_1k_2) intersections. In principle, the contributions from the two tori can cancel leaving massless states rather than tachyons. For this to happen for the charged-singlet tachyons at the (i_1i_2), (j_1j_2) and (k_1k_2) intersections, the parameters must satisfy the conditions $\frac{\epsilon R_2/R_1}{\tilde{\epsilon}\tilde{R}_2/\tilde{R}_1}=\left|\frac{x(p-x)}{y(p-y)}\right|=\left|\frac{x}{y}\right|=\left|\frac{p-x}{2p-y}\right|$ where $x\equiv n_2-3R_2/R_1$, $y\equiv \tilde{n}_2-\tilde{R}_2/\tilde{R}_1$, $R_1,R_2,\tilde{R}_1,\tilde{R}_2$ are the radii of the fundamental 1-cycles on the two wrapping tori and $\epsilon\equiv 2|cos(\theta/2)|, \tilde{\epsilon}\equiv 2|cos(\tilde{\theta}/2)|$. These are all satisfied if

$$px=\frac{3}{2}=py \quad \text{or if} \quad px=\frac{2}{3}=\frac{1}{2}py \tag{12}$$

The 3 colour-triplet scalars that arise at the $(1k_1)$ intersection, as well as the one at the $(1k_2)$ intersection, can be expunged by taking the separation between the parallel 1-cycles to be sufficiently large.

For the case $T_{22} = 0 = T_{31}$, again only $U(1)_X$ and $U(1)_Y$ evade acquiring GS masses and so survive as gauged symmetries. To avoid triplication of gauge group factors requires that $n_2 = -p \bmod 3 = q \bmod 3$. In this case we find the vector-like matter to be: $3(L + \bar{L}) + 12(e^c_L + \bar{e}^c_L) + 3(\nu^c_L + \bar{\nu}^c_L)$. The total tachyonic Higgs, colour-triplet and charged-singlet scalar content is the same as for the above model. In this case though, there are 3 Higgs doublets at the $(2i_1)$ intersections, and one at $(2i_2)$; the colour-triplet scalars are similarly interchanged. They are massless if

$$px = -\frac{3}{2} = -py \quad \text{or if} \quad px = \frac{4}{3} = -2py \tag{13}$$

Ratios of the gauge coupling constants are independent of the Type II string coupling constant λ_{II}. Thus $\frac{\alpha_3(m_{\text{string}})}{\alpha_2(m_{\text{string}})} = |xy|$. Also, for the $T_{21} = 0 = T_{32}$ model we have that

$$\frac{\alpha_3(m_{\text{string}})}{\alpha_Y(m_{\text{string}})} = \frac{1}{3} + \frac{1}{2}|xy| + |x(p-y)| + |x| + |px-1| \tag{14}$$

Consistency with a low string scale requires that these ratios do not differ greatly from the values measured [14] at the electroweak scale m_Z, $\frac{\alpha_3(m_Z)}{\alpha_2(m_Z)} = 3.54$, $\frac{\alpha_3(m_Z)}{\alpha_Y(m_Z)} = 11.8$. It is easy to find all solutions of these:

$$\begin{aligned} px &= 4.78,\ -1.75,\ 2.41,\ -4.11 \\ py &= 0.74,\ -2.02,\ -1.46,\ 0.86 \end{aligned} \tag{15}$$

which can be satisfied with reasonable values of the parameters. For example, $n_2 = -2 = 1 \bmod 3,\ R_2/R_1 = 0.93,\ \tilde{n}_2 = 1,\ \tilde{R}_2/\tilde{R}_1 = 1.74$. None of these solutions is close to satisfying the conditions (12) for massless charged-singlet scalars, so the existence of these tachyons is unavoidable in this model unless there are very large radiative corrections.

For the $T_{31} = 0 = T_{22}$ model, eqn (14) is replaced by $\frac{\alpha_3(m_{\text{string}})}{\alpha_Y(m_{\text{string}})} = \frac{1}{3} + \frac{1}{2}|xy| + |(x-p)y| + |y| + |py-1|$ and the solutions are obtained by interchanging px and py in (15). Again, it is easy to satisfy these with reasonable values of the parameters. For example, $n_2 = 1,\ R_2/R_1 = 1.00,\ \tilde{n}_2 = -1,\ \tilde{R}_2/\tilde{R}_1 = 0.75$. In this case, the third solution is not too far removed from satisfying the first condition in (13) and it is possible that the effects of renormalisation group running and/or radiative corrections to the masses could remove the charged-singlet scalar tachyons from the low-energy spectrum.

In conclusion, we have obtained the standard model gauge group with *no* additional, unwanted gauged $U(1)$ symmetries, plus three (non-supersymmetric) generations of chiral matter that include right-chiral neutrino states using intersecting D5-branes. This required at least 8 stacks. This parallels the situation in orientifold models, in which a minimum of four stacks, plus their orientifold images, is required [8]. We found two distinct classes of model. Both classes of model have extra vector-like leptons, as well as charged-singlet scalar tachyons, and optional colour-triplet tachyons. There are radiative corrections to the masses of both that might be sufficient to render them non-tachyonic, without removing the Higgs doublets. For particular values of the wrapping numbers, both classes of model predict ratios of gauge coupling constants close to those measured at the electroweak scale, thereby allowing the possibility of a nearby string scale, such as is required of non-supersymmetric theories to avoid the hierarchy problem. For the model with $T_{31} = 0 = T_{22}$, one choice of the parameters comes close to removing the charged-singlet scalar tachyons from the low-energy spectrum. Both classes of model possess the Yukawa couplings of the tachyonic Higgs doublets needed to generate masses at renormalisable level for three generations of chiral matter including neutrinos. However, neither class of model has masses at renormalisable level for all of the vector-like matter. All models have anomalous $U(1)$s that survive as global symmetries. Their gauge bosons acquire string-scale masses via the generalised GS mechanism. It is expected that TeV-scale Z' vector bosons will be observable at future colliders, and precision electroweak data (on the ρ-parameter) already constrain [15] the string scale to be at least 1.5 TeV. In particular, baryon number $B = Q_1/3$ is anomalous and survives as a global symmetry. Consequently, the proton is stable despite the low string scale. As before in the D4-brane case [11], the Higgs boson fields are also charged under some of the anomalous $U(1)$s that survive as global symmetries. Thus a keV-scale axion is unavoidable.

Acknowledgements

This account is a modified version of the talk that G. Kraniotis gave in Hamburg in June, 2002. This research is supported in part by PPARC and the German-Israeli Foundation for Scientific Research (GIF).

References

[1] G. Aldazabal, L. E. Ibáñez, and F. Quevedo, JHEP 0001(2000)031, hep-th/9909172, JHEP 0002 (2000) 015, hep-th/0005067; N.D. Lambert, P.C. West, JHEP 9909 (1999)021, hep-th/9908129; I. Antoniadis, E. Kiritsis, and T. Tomaras, Phys. Lett. B486 (2000) 186, hep-th/0004214; D. Bailin, G.V. Kraniotis, and A. Love, Phys. Lett. B502 (2001) 209, hep-th/0011289

[2] M. Berkooz, M. R. Douglas, and R. G. Leigh, Nuclear Physics B480 (1996) 265, hep-th/9606139;C. Bachas, hep-th/9503030

[3] N. Arkani-Hamed, S. Dimopoulos, and G.R. Dvali, Phys. Lett. B429 (1998) 263, hep-ph/9803315; I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G.R.Dvali, Phys. Lett.B 436 (1998) 257, hep-ph/9804398.

[4] R. Blumenhagen, B. Körs, D. Lüst, and T. Ott, Nucl. Phys.B616 (2001) 3, hep-th/0107138

[5] R. Blumenhagen, B. Körs, D. Lüst, and T. Ott, hep-th/0202124; Burgess et al, hep-th/0111025

[6] S.B. Giddings, S. Kachru, and J. Polchinski, hep-th/0105097

[7] M. Cvetič, G.Shiu, and A. M. Uranga, Nucl. Phys. B615 (2001) 3, hep-th/0107166; M. Cvetič, P. Langacker, and G. Shiu, hep-ph/0205252, hep-th/0206115

[8] L. E. Ibáñez, F. Marchesano and R. Rabadan, JHEP 0111 (2001) 002, hep-th/0105155; D. Cremades, L. E. Ibáñez and F. Marchesano, hep-th/0205074;

[9] L. E. Ibáñez, R. Rabadán, and A. M. Uranga, Nucl. Phys. B542 (1999) 112, hep-th/9808139

[10] D. Bailin, G. V. Kraniotis, and A. Love, Phys. Lett. B547 (2002) 43, hep-th/0208103

[11] D. Bailin, G. V. Kraniotis, and A. Love, Physics Letters B530 (2002) 202, hep-th/0108131

[12] G. Aldazabal, S. Franco, L. E. Ibáñez, R. Rabadán, and A. M. Uranga, J. Math. Phys. 42 (2001) 3103, hep-th/0011073

[13] D. Bailin, G. V. Kraniotis, and A. Love, in preparation

[14] Reports of the working groups on precision calculations for LEP2 Physics, S. Jadach, G. Passarino, R. Pittau (Eds.), CERN Yellow Report, CERN, 2000-009

[15] D. M. Ghilencia, L. E. Ibáñez, N. Irges, and F. Quevedo, hep-ph/0205083