TeV-scale $Z'$ bosons in intersecting D-brane SM-like models

D.M. Ghilencea

DAMTP, CMS, University of Cambridge
Wilberforce Road, Cambridge, CB3 0WA, United Kingdom.

Abstract

Recent string constructions with intersecting D6 and D5 brane models succeeded in predicting in the low energy limit a symmetry group and a fermionic spectrum similar to that of the SM. In such constructions additional $U(1)$ fields are a generic presence and they (or linear combinations thereof) become massive through a string mechanism which couples them to 4D RR two-forms. $U(1)_Y$ of hypercharge emerges as a linear combination of the initial $U(1)$ symmetries while the massive $U(1)$ fields induce (through mixing with Z boson) corrections on Z mass. $\rho$ parameter constraints on the latter allow one to set lower bounds on the string scale in the TeV region. These can then be used to predict lower bounds on the masses of the additional $Z'$ bosons, without specific assumptions about the compactification volume.

Contribution to the Proceedings of the
1st International Conference on String Phenomenology, Oxford (U.K.), 6-11 July 2002.

---

1This contribution is based on work done in collaboration with L.E. Ibáñez, F. Quevedo and N. Irges.
1 Introduction.

The possibility of (large) extra space dimensions has attracted much research interest in the context of both string and effective field theories. In general the latter can be successfully used for many theoretical/phenomenological studies in this direction. However, it is generally thought that a more complete and consistent picture of the high energy physics is that brought by string theory. It is thus desirable to obtain viable field theory models as direct constructions of string theory. For this purpose, recent chiral D-brane string constructions with branes located at singularities \[1, 2, 3\] or models with D-branes intersecting at non-trivial angles \[4, 5, 6\] have brought in a new approach, alternative to that of (earlier) heterotic string compactifications. Finding out which of these approaches is more successful in obtaining in the low energy limit models close to the Standard Model (SM) or its minimal supersymmetric extension (MSSM) is a rather difficult problem. This is due to the large variety of different string vacua. For this reason, identifying some generic “predictions” of string compactifications may provide an useful approach to identifying viable string models.

In the following some phenomenological implications of the D-brane models are addressed. To obtain chiral D-brane models one uses either models with D-brane at singularities or with intersecting branes. These share some common properties: a low value for the string scale is allowed, the gauge symmetry contains direct products of groups \(U(N_\alpha) \times U(N_\beta)\) with \(U(N_\alpha)\) to arise from the stack “\(\alpha\)” of \(N\) individual \(U(1)\) branes; fermions can transform as bi-fundamental representations of these groups; \(U(1)\) group factors in addition to the non-Abelian SM gauge group are a generic presence and the \(U(1)\) of hypercharge emerges as a linear combination of these. Other characteristics are more model dependent, for example the models may be supersymmetric or non-supersymmetric \[4\]. The latter requires fixing a low value for the string scale, to avoid re-introducing the hierarchy problem.

D-branes models with a spectrum somewhat close to that of the SM include constructions with D4 branes and D5 or D6 branes intersecting at non-trivial angles. In the following we address the specific class of models with intersecting D6-branes of \[4\], although some of the results are more general. The models are constructions with D6 branes wrapping a 3-cycle on a six torus (orientifolded) type II-A string theory. It involves four stacks \(\alpha = a, b, c, d\) of D6-branes with gauge groups \(U(3)_a, U(2)_b, U(1)_c\) and \(U(1)_d\), containing inside the Minkowski space and wrapping each of the remaining three dimensions of the branes on a different torus \(T^2\). We denote by \(n_{\alpha i}, (m_{\alpha i}), i = 1, 2, 3, \alpha = a, b, c, d\) the number of times each brane \(D6_{\alpha}\) is wrapping around the \(x(y)\)-coordinate of the \(i-th\) torus. This construction with four stacks is the minimal one which can accommodate a fermionic spectrum similar to that of the SM case (plus three generations of right-chiral neutrinos).
However, the gauge symmetry contains in addition to the non-Abelian part of the SM gauge group four additional $U(1)\alpha$, $\alpha = a, b, c, d$ which emerge from the four stacks of branes considered. Such $U(1)$'s are a generic presence in models with intersecting branes and they play a central role in the discussion below. (Similar considerations apply to D5 brane models \[8\] which are type II-B compactifications on an orbifold $T^2 \times T^2 \times (T^2/Z_N)$).

An important feature of the models considered is the presence in the 4D action of the couplings $c_i^\alpha B_i \wedge F_\alpha$, where $F_\alpha$ is the field strength of a $U(1)\alpha$ field and $B_i$ is a 4D RR two-form. The couplings emerge, in the case of D6 branes from $C_5 \wedge F$ integrated over a three-cycle of the six-torus. Such couplings together with the kinetic term for $F_\alpha$ and $B_i$ can be re-written as a kinetic term for a $U(1)$ field together with an associated mass term. This is possible because the scalar dual to the $B_i$ two-form is “eaten” by the $U(1)\alpha$ field which thus becomes massive \[4\]. The procedure does not require a Higgs mechanism or a Higgs field be present in the spectrum. Note that the mechanism requires the coupling $B_i \wedge F_\alpha$, but not $\eta F_\alpha \wedge F_\alpha$ coupling. For a $U(1)$ field with a Green-Schwarz coupling $B \wedge F$, being anomalous (in four dimensions) is not a necessary condition to acquire a mass, and anomaly free $U(1)$’s can become massive. In previous (heterotic) string models anomalous $U(1)$ also had a Green-Schwarz term and only these became massive. Further, since there is no Higgs mechanism, the initial (or combinations thereof) $U(1)\alpha$ may survive as a perturbatively exact global symmetry of the models.

This observation can be useful for phenomenological purposes. For example the presence of $U(1)\alpha$ as a global symmetry ensures that baryon number is conserved, thus avoiding proton decay, a generic problem in models with a low string/UV cut-off scale. For this class of models (diagonal) lepton number is also an exact symmetry and Majorana neutrino masses are forbidden. Dirac neutrino masses are allowed and neutrino oscillations can take place (only the sum $L_e + L_\mu + L_\tau$ is an exact symmetry).

An important aspect of D6- and D5-brane models is that of the relation between anomaly cancellation and global and local tadpole cancellation. In these models non-Abelian gauge anomalies are cancelled by the requirement of global tadpole cancellation. $U(1)\alpha G_\beta G_\beta$ mixed anomalies are cancelled by a Green Schwarz mechanism together with tadpole cancellation requirement. Finally, cubic anomalies are cancelled by the Green Schwarz mechanism. One can draw the conclusion that tadpole cancellation is crucial for the overall consistency of these models. For a study of local anomalies in the context of type II B string constructions and their implications see reference \[13\].

Apart from ensuring anomaly cancellation, global tadpole cancellation imposes an additional condition such that for a $U(N)$ group the number of fundamental representations be equal to that of anti-fundamental representations. This requirement applies even to $U(2)$ and $U(1)$ gauge groups with the consequence of restricting the way we assign the quarks and leptons as $U(2)$ doubles and/or
anti-doublets. If all quarks were doublets and leptons anti-doublets, then the above condition would not be respected. The only possibility to equal the number of doublets and anti-doublets is to have two families of quarks as $U(2)$ doublets, with the third one and with the leptons as $U(2)$ anti-doublets. This only works for three generations and thus elegantly relates the number of colours to that of generations [4].

For the D6-brane models of [4] that we address in the following, the aforementioned couplings $c_i^\alpha B_i \wedge F_\alpha$ bring mass terms for three of the initial four $U(1)_\alpha$. However [4] one linear combination of these $U(1)$ remains massless, because its coupling to any RR two-form field $B_i$ vanishes, independent of the parameters of the model. An additional appropriate condition imposed on these allows one to identify the generator of this massless $U(1)$ with that of hypercharge, which remains a gauge symmetry of the model. The couplings $\sum_{i,\alpha} c_i^\alpha B_i \wedge F_\alpha$ provide a mass term in the action of the structure

\[
(M^2)_{\alpha \beta} = g_\alpha g_\beta M_S^2 \sum_{i=1}^{3} c_i^\alpha c_i^{\beta}, \quad c_i^\alpha = N_{\alpha n_i j n_{ak} m_{\alpha i}}, \quad i \neq j \neq k \neq i, \quad \alpha, \beta = a, b, c, d. \tag{1}
\]

where $g_\alpha$ is the coupling of associated $U(1)_\alpha$, $c_i^\alpha$ result from integrating for D6 branes over 3-cycle $C_i \wedge F_\alpha$ with $C_i$ RR 5-form, and $M_S$ is the string scale defined up to an overall (volume dependent) factor. The eigenvalues of this mass matrix (squared masses) are positive [7] irrespective of the parameters of the models. The masses (relative to the string scale) and associated eigenvectors of the three massive $U(1)$ fields can be computed explicitly [7]. Generic results for the masses are within a factor of 10 above and below the string scale respectively, with the third mass of the order of $M_S$. A fourth $U(1)$ field remains massless, and this is identifiable with the $U(1)_Y$ of hypercharge after an additional (model building) constraint on the wrapping numbers.

The next step is that of including the effects of the electroweak symmetry breaking. For this we considered the minimal Higgs content predicted in this class of models [4] which is at least two Higgs doublets[4]. Such effects bring small corrections to the massive $U(1)$ fields computed in the absence of EW symmetry breaking, since they are suppressed by $\eta \equiv <v>^2/M_S^2$. There is however a more important aspect to be addressed. The mass of the $Z$ boson itself receives corrections from mixing with the additional $U(1)$ fields, since the Higgs sector is also charged under $U(1)_{b,c}$. Since the masses of the $U(1)$ fields are essentially of string origin, induced by their initial couplings to $B_i$ fields, the mass of $Z$ boson will itself receive such a correction in addition to its SM value. The correction to the mass of the $Z$ boson can be computed as an expansion in $\eta$ (see details in [4])

\[
M_Z^2 = M_0^2 \left[ 1 + \eta \xi_{21} + \eta^2 \xi_{31} + \cdots \right] \tag{2}
\]

\[\text{This is a somewhat generic feature in string models [4].}\]
with \( M_0 \) the SM \( Z \) boson mass and \( \xi_{21} \) given by (for \( \xi_{31} \) see Appendix of [8])

\[
\xi_{21} = - \left\{ \beta_1^2 \left[ 2 \beta_1 g_y^2 \nu n_{c1} (1 + R^2) - (36 g_a^2 + g_y^2 R^2) \beta_2 n_{a2} \nu \right]^2 + 4 \beta_1^2 \beta_2^2 \epsilon^2 (36 g_a^2 + g_y^2 R^2)^2 + 9 \beta_2^4 \left[ g_y^2 n_{b1} R^2 + 12 g_a^2 \nu n_{b1} - n_{c1} (1 + R^2) \cos(2\theta) \right]^2 \right\} \left[ 5184 \beta_1^2 \beta_2^2 \epsilon^2 g_a^2 n_{c1} (1 + R^2)^2 \right]^{-1}
\]

where \( R = g_d/g_c \) and \( n_{a2} \) are model parameters, \( n_{c1} \) and \( n_{b1} \) are either 0 or \( \pm 1 \) function of the model, \( \nu = 1, 1/3, \beta_{1,2} = 1, 1/2, g_a^2 = g_{QCD}^2/6 \) and \( g_y \) is the hypercharge coupling. \( \theta \) is a mixing angle in the Higgs sector (which does not affect significantly the predictions). Correction (4) must be compatible with electroweak scale measurements, and this may be ensured using \( \rho \)-parameter constraints. From this a lower bound on the allowed values of the string scale \( M_S \) emerges. This bound should not be too large compared to the TeV scale because the models we address are non-supersymmetric and one would otherwise re-introduce a hierarchy problem. Detailed calculations show that generic lower bounds on \( M_S \) are within the range \( 1.5 - 5 \) TeV [8] and are to a great extent parameter-independent.

The analysis so far has ignored the fact that the coefficients \( c_i^a \) in eq.(1) depend on the normalisation of the kinetic terms of RR fields \( B_i \). In a canonical normalisation of the latter, extra volume factors \( \xi_i \) appear in the definition of \( c_i^a \). If equal for the three torii \( \xi_i = \xi_j, i, j = 1, 2, 3 \) (this can be easily respected if the ratio of the radii is the same for all torii) the volume factors can be “absorbed” in the re-definition of the string scale \( M_S \), as we already did in eq.(1). Therefore the lower bounds on \( M_S \) that we found in the region \( 1.5 - 5 \) TeV are subject to such additional corrections, difficult to evaluate numerically on string theory grounds.

One can take this analysis a step further [8] by using the bounds on (the re-scalled) \( M_S \) compatible with \( \rho \) parameter constraints to obtain lower bounds on \( U(1) \) masses. These will be independent of the volume factors mentioned and represent a prediction in this class of models. Of particular importance and experimental relevance is that of mixing of \( Z \) eigenvector with any of the massive \( U(1) \) fields, induced after electroweak symmetry breaking. This is usually of order \( 10^{-3} \) or possibly less [12]. Therefore, computing the eigenvectors corresponding to the mass eigenstates after electroweak symmetry breaking is necessary [8], which also helps identify specific signatures of the \( Z' \) bosons and compare them with other models.

Our results are presented in Table 1 and show that \( Z' \) bosons in the TeV-region may be present in these models. The bounds are somewhat larger than those of their counterparts in (heterotic) string models or extensions of SM (MSSM) which are in the range [12]: \( Z'_{SM} \) boson defined to have the same couplings to fermions as the SM \( Z \) boson: 809 GeV; \( Z_{LR} \) of \( SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_2 \): 564 GeV; \( Z_\chi \) of \( SO(10) \rightarrow SU(5) \times U(1)_\chi \): 545 GeV; \( Z_\psi \) of \( E_6 \rightarrow SO(10) \times U(1)_\psi \): 146 GeV.

So far we only outlined part of the phenomenological implications of the D6 brane models. For a full picture of the viability of the models there are additional issues to be addressed. These include...
the values of the gauge couplings at the TeV scale and the extent to which these can be compatible with the electroweak scale measurements. One potential drawback of the models discussed here is that unification of the gauge couplings is not possible (given the non-supersymmetric spectrum). However this is a more generic problem, manifest in all models with a low string scale, whether supersymmetric or not [11]. In the absence of unification, one should then ensure that the ratio of the couplings at the string scale is that obtained from the RG flow using low energy values as input. This is relevant since the value of gauge couplings at the string scale depends ultimately on the volume each D6 brane is wrapping, bringing additional constraints on the models. There is one final complication because threshold effects induced by additional momentum states can affect significantly the value of the couplings [11], making the whole analysis rather difficult.

The class of models considered here for the phenomenological analysis is free of tachyons and RR tadpoles, but the models are non-supersymmetric and there will in general be NSNS tadpoles.
which is generic in these constructions. For the D6-brane models addressed one may also ask how one could achieve a low value for the string scale. Since the models are non-supersymmetric, theory should fix this value in the few TeV region. (We have seen from the numerical analysis that low energy constraints as $\rho$ parameter can be accommodated with a TeV string scale, which is welcome for the viability and consistency of the models). However, for the D6 brane models yielding the SM fermionic spectrum there is no dimension transverse to all the intersecting branes, to allow one to predict the right value of 4D Planck scale. The problem is not present in the D5 brane models mentioned in this talk, which yield similar results for the $U(1)$ masses and bounds on $M_S$ from the $\rho$ parameter [4]. One solution for the D6-brane models to have a low string scale was suggested in [3]. The idea is to construct intersecting D6 branes wrapped on 3-cycles localised in a small region of a Calabi-Yau manifold. By enlarging the dimensions of the Calabi-Yau which are transverse to the set of D6 branes the 4D Planck mass is recovered even for low (“TeV-scale”) values of $M_S$. This would not require enlarging the size of (local) 3-cycles (and would not bring about unobserved light Kaluza-Klein modes that would appear in such case).

To conclude, the class of D6 and D5 brane models addressed provide an interesting alternative to previous heterotic constructions and are successful in deriving the symmetry and the fermionic spectrum of the SM. They may thus provide us an understanding of the string embedding of the Standard Model. While such models are intensively studied at the string level, their phenomenological viability was not investigated systematically. Our attempt to address it in [7], [8] (see also recent [17]) shows that for a suitable choice of model parameters, the models may be able to comply with the low energy constraints considered there.

Acknowledgements: The author thanks L. E. Ibáñez, F. Quevedo and N. Irges for their collaboration. Helpful discussions on related topics with S. Förste, G. Honecker, R. Rabadán and A. Uranga are acknowledged. This work was supported by PPARC (U.K.).

References

[1] G. Aldazabal, L. E. Ibáñez and F. Quevedo, JHEP 0001 (2000), 031; JHEP 0002 (2000) 015; D. Bailin, G. V. Kraniotis and A. Love, Phys. Lett. B 502, 209 (2001); D. Berenstein, V. Jejjala and R. G. Leigh, Phys. Rev. Lett. 88, 071602 (2002); D. Bailin, G. V. Kraniotis and A. Love, [arXiv:hep-th/0108127]; L. L. Everett, G. L. Kane, S. F. King, S. Rigolin and L. T. Wang, [arXiv:hep-ph/0203129]; L. F. Alday and G. Aldazabal, [arXiv:hep-th/0203129].

[2] G. Aldazabal, L. E. Ibáñez, F. Quevedo, A. M. Uranga, JHEP 0008 (2000) 002.
[3] M. Cvetic, A. Uranga and J. Wang, Nucl. Phys. B 595, 63 (2001), [arXiv:hep-th/0010091].

[4] L. E. Ibáñez, F. Marchesano and R. Rabadán, JHEP 0111 (2001) 002 [arXiv:hep-th/0105155].

[5] R. Blumenhagen, L. Goerlich, B. Kors, D. Lust, JHEP 0010, 006 (2000); Fortsch. Phys. 49, 591 (2001); G. Aldazabal, S. Franco, L. E. Ibáñez, R. Rabadán, A. M. Uranga, J. Math. Phys. 42, 3103 (2001); G. Aldazábal, S. Franco, L. E. Ibáñez, R. Rabadán, A. M. Uranga, JHEP 0102, 047 (2001); R. Blumenhagen, B. Kors, D. Lust, JHEP 0102, 030 (2001); M. Cvetic, G. Shiu, A. M. Uranga, Phys. Rev. Lett. 87, 201801 (2001); Nucl. Phys. B 615, 3 (2001); [arXiv:hep-th/0111179]; R. Blumenhagen, B. Kors, D. Lüst, T. Ott, Nucl.Phys. B 616 (2001) 3-33; and [arXiv:hep-th/0112015]; D. Bailin, G. V. Kraniotis, A. Love, Phys.Lett.B530 (2002) 202, Phys.Lett.B547 (2002),43, [arXiv:hep-th/0210219]; D. Bailin, [arXiv:hep-th/0210227]; H. Kataoka and M. Shinojo, [arXiv:hep-th/0112247]; G. Honecker, [arXiv:hep-th/0201037]; Fortsch.Phys.50:896-902.2002; S. Forste, G. Honecker, R. Schreyer, JHEP 0106 (2001) 004; C. Kokorelis, [arXiv:hep-th/0203187]; D. Cremades, L. E. Ibáñez, F. Marchesano, [arXiv:hep-th/0201205]; [arXiv:hep-th/0203160].

[6] D. Cremades, L. E. Ibáñez, and F. Marchesano, [arXiv:hep-th/0205074].

[7] D. M. Ghilencea, L. E. Ibáñez, N. Irge, F. Quevedo, JHEP 0208 (2002) 016 [hep-ph/0205083].

[8] D. M. Ghilencea, [arXiv:hep-ph/0208205], to appear in Nuclear Physics B.

[9] A.M. Uranga, [arXiv:hep-th/0208014].

[10] D. Ghilencea, G. G. Ross, Phys. Lett. B 480 (2000) 355 [arXiv:hep-ph/0001143]; Nucl. Phys. B 606 (2001) 101 [arXiv:hep-ph/0102306]; Phys. Lett. B 442 (1998) 165 [arXiv:hep-ph/9809217].

[11] D. M. Ghilencea and H. P. Nilles, J. Phys. G 28 (2002) 2475 [arXiv:hep-ph/0204261].

[12] K. Hagiwara et al., Phys. Rev. D 66 (2002) 010001.

[13] C. A. Scrucca, M. Serone and M. Trapletti, Nucl. Phys. B 635 (2002) 33.

[14] A. Leike, Phys. Rept. 317 (1999) 143 [arXiv:hep-ph/9805494].

[15] J. Erler and P. Langacker, Phys. Rev. Lett. 84 (2000) 212; Phys. Lett. B 456 (1999) 68; P. Langacker, in Proc. of the APS/DPF/DPB Summer Study on the Future of Particle Physics (Snowmass 2001) ed. R. Davidson and C. Quigg, [arXiv:hep-ph/0110120].

[16] D. Ghilencea, S. Groot-Nibbelink, Nucl. Phys. B 641 (2002) 35 [arXiv:hep-th/0204094].

[17] D. Cremades, L. E. Ibáñez, and F. Marchesano, [arXiv:hep-ph/0212064].