SUPERSYMMETRY AND ELECTROWEAK BREAKING BY EXTRA DIMENSIONS

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(Received June 10, 2002)

Dedicated to Stefan Pokorski on his 60th birthday

We discuss the phenomenological consequences of large extra dimensions concerning both supersymmetry and electroweak symmetry breaking. We consider separately the fundamental scenarios where this can happen. In particular cases where only the gravitational sector can propagate in the bulk of the large extra dimensions, and cases with longitudinal dimensions where all gauge and matter fields propagate. We briefly comment on the string realization of these ideas and finally present a possible scenario where electroweak breaking is triggered by the Hosotani mechanism and thus a finite Higgs mass does not require supersymmetry.

PACS numbers: 04.50.+h, 11.25.Mj, 11.30.Qc

1. Introduction

Electroweak symmetry breaking is one of the main issues in contemporary particle physics. Its implementation in a perturbative quantum field theory has led to the notion of spontaneous gauge symmetry breaking, and in particular to that of the Higgs mechanism, requiring the existence of the Higgs boson. This in turn has both experimental and theoretical consequences. On the experimental side, the Higgs boson is the missing ingredient in the Standard Model of strong and electroweak interactions and its

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detection the main goal in present and future high energy particle accelerators. On the theoretical side, the presence of the Higgs boson generates an inherent hierarchy problem in the Standard Model and has motivated the introduction of supersymmetry.

Supersymmetry is a symmetry between bosons and fermions with the same mass and internal quantum numbers. In particular the minimal Supersymmetric Extension of the Standard Model (MSSM) implies that every ordinary particle has an associated superpartner that should show up in direct or indirect experimental searches. In fact supersymmetric searches are one of the main experimental goals in future colliders. However, supersymmetry is not an exact symmetry of Nature since there are no scalar particles degenerate in mass with the Standard Model leptons and quarks. Thus, the mechanism of supersymmetry breaking is another open problem in particle physics.

The possibility of large extra dimensions [1] and low string (quantum gravity) scale [2–4] opened after the introduction of duality in string theory. This was the case in the strong coupling regime of the heterotic string, or M-theory [5], where the eleventh dimension can be large and controls the size of the M-theory scale. The other example is that of type I strings. There the Standard Model fields can live on a Dp-brane, a $(p+1)$-dimensional hypersurface embedded into the ten-dimensional space-time, where open strings end. In this scenario, there are p-3 extra dimensions where gauge and/or matter fields propagate and whose size can be as large as $\sim 1/\text{TeV}$ [6]. On the other hand, the remaining 9-p extra dimensions transverse to the p-brane can be much larger, up to a millimeter, allowing to lower the string scale down to TeV energies. In both cases, the presence of large extra dimensions naturally appear in the theory and can help in understanding both problems of supersymmetry and electroweak symmetry breaking.

In this review we will concentrate on the phenomenological consequences of large extra dimensions, concerning both the issues of supersymmetry and electroweak symmetry breaking. As we will see, they lead to distinct scenarios that can be tested in future accelerators. A positive signal would uncover extra dimensions and would hint for the physics responsible of reconciling gravity and quantum mechanics. The scenario where extra dimensions are populated only by the gravitational sector is considered in Section 2. Section 3 contains the case where Standard Model fields live in the extra dimensions. The status of embedding these scenarios in M-theory or explicit string models is briefly summarized in Section 4. The presence of large extra dimensions allows for a non-supersymmetric solution to the hierarchy problem if the electroweak symmetry breaking proceeds through a Hosotani mechanism [7]. This issue is discussed in Section 5.

2. Supersymmetry breaking transverse to the branes

We consider an effective five-dimensional supergravity theory with cutoff at the scale M_5 and compactified to four dimensions on the orbifold S^1/Z_2. The gravitational sector is propagating in the bulk of the fifth dimension while gauge and matter fields are localized in the walls. This theory has been recently given an off-shell formulation [8] and used for phenomenological purposes to break supersymmetry by the Scherk–Schwarz (SS) mechanism [9] and/or by gaugino condensation on the hidden wall [10–16].

A natural breaking of supersymmetry in this theory is by boundary conditions, or Scherk–Schwarz mechanism, based on the $\mathrm{SU}(2)_R$ symmetry of the $N=2$ supersymmetry algebra. The gravitino mass eigenvalues for the Kaluza–Klein (KK) modes are $m_{3/2}^{(n)}=(n+\omega)/\rho$, where ω is the SS parameter and ρ is the radius of S^1. This mechanism has recently [16] been given an interpretation in terms of Wilson line breaking with $\omega=\rho V_5^2$, where $\vec{V}_5$ are the auxiliary fields that gauge $\mathrm{SU}(2)_R$ in the off-shell formulation of $N=2$ supergravity in 5D. Therefore, the tree-level potential for V_5^2 is flat, reminiscent of no-scale models of supergravity, and its one-loop effective potential is given by

$$V_{\mathrm{eff}}(\omega)=\frac{3}{34\pi^6\rho^4}\left[Li_5\left(\mathrm{e}^{2i\pi\omega}\right)+\text{ h.c.}\right]\,. \tag{2.1}$$

The potential (2.1) has a maximum at $\omega=0$ and a global minimum at $\omega=1/2$. Therefore, the zero mode gravitino mass is $m_{3/2}=1/2\rho$ as first obtained in Ref. [11].

Another possible way of supersymmetry breaking is by gaugino condensation on the hidden brane (wall), say at $x^5=0$, upon confinement of the corresponding gauge group. As we will see, this breaking is equivalent to the previous Scherk–Schwarz breaking. The physical picture is that a condensate $\langle\lambda\lambda\rangle$ develops at a scale Λ_c where the gauge coupling of the hidden gauge group becomes strong. This phenomenon can be described by introducing a chiral supermultiplet whose Vacuum Expectation Value (VEV) reproduces the condensate, and a non-perturbative superpotential $W\propto\langle\lambda\lambda\rangle$ that contributes a brane term to the 5D Lagrangian, as [14]

$$\mathcal{L}_{\mathrm{brane}}^{3/2}=\tfrac{1}{2}W\delta(x_5)\bar{\psi}_\mu\gamma^{\mu\nu}\psi_\nu\,, \tag{2.2}$$

where ψ_μ is the 4D Majorana gravitino.

The term (2.2) introduces a "democratic" mass matrix for the gravitino KK modes that interferes with the diagonal compactification and SS masses. Rediagonalization now yields for the gravitino KK mass eigenstates the values $m_{3/2}^{(n)}=(n+\Delta_{3/2})/\rho$, where [16, 17] $\Delta_{3/2}(\omega,W)$ is a function of the auxiliary field V_5^2 through the SS parameter ω, and of the non-perturbative superpotential W. Even if W is a parameter whose value has been fixed by the dynamics of gaugino condensation, ω is undetermined at the tree-level and hence $\Delta_{3/2}$ is a flat direction. This field can be fixed by introducing one-loop corrections. The effective potential is now

$$V_{\rm eff}(\omega) = \frac{3}{34\pi^6\rho^4}\left[Li_5\left(e^{2i\pi\Delta_{3/2}}\right) + \text{h.c.}\right], \tag{2.3}$$

whose minimization yields $\Delta_{3/2} = 1/2$ and hence for the gravitino zero mode $m_{3/2} = 1/2\rho$. This indicates that gaugino condensation is equivalent to the SS mechanism in M-theory, as anticipated in Ref. [11].

The communication of supersymmetry breaking to the matter sector, localized on the visible brane, is expected to proceed by radiative gravitational corrections. This calculation was performed in Refs. [11,18]. For the scalar partners of the Standard Model fermions one finds

$$m_0 \simeq m_{3/2}^2/M_p\,, \tag{2.4}$$

which yields $m_0 \simeq 1$ TeV for $\rho^{-1} \simeq 10^{11}$ GeV. For the fermionic partners of the Standard Model gauge bosons the one-loop result is finite [11] $m_\lambda \propto m_{3/2}^3/M_p^2$, which is exceedingly small for phenomenological purposes, with a proportionality constant which can be zero depending on the regularization procedure [18]. In any case, higher-loop corrections should be tiny for phenomenology and we conclude that the minimal scenario has to be enlarged to allow for an extra source of gaugino masses.

3. Supersymmetry breaking parallel to the branes

We consider here the simplest example of a five dimensional theory compactified on the orbifold $\mathcal{M}_4 \times S^1/Z_2$, as in the previous section, but with the difference that now also gauge and matter fields can propagate in the bulk of the extra dimension. Of course, in this case, the typical radius R of the extra dimension must not exceed $\sim 1/\text{TeV}$ to agree with electroweak precision physics.

3.1. Supersymmetry breaking

The simplest mechanism of supersymmetry breaking is by boundary conditions (Scherk–Schwarz compactification). In fact, the SS parameter ω can be given a similar interpretation as in the previous section just by changing the compactification radius ρ to R, *i.e.* $\omega = RV_5^2$. In this case, depending on the matter content propagating in the bulk, one-loop corrections fix ω either to 0 or to 1/2 [16].

However, in the case of a hidden brane with supersymmetry breaking there can be an interplay between both phenomena similar to that existing in M-theory for the gravitino. In particular if the gauginos and hyperscalars couple to supersymmetry breaking as [16]

$$\mathcal{L}^{1/2}_{\rm brane} = \tfrac{1}{2}\delta(x_5) W \bar{\lambda}\gamma^5\lambda\,, \tag{3.1}$$

where λ are the 4D Majorana gauginos, and similarly for hyperscalars, the KK mass eigenvalues are $m_{0,1/2} = (n+\Delta)/R$, where the shift $\Delta(\omega, W)$ is again a function of V_5^2 through the SS parameter. Minimization of the effective potential yields, depending on the matter content of the bulk and the value of W, any possible value of ω between 0 and $1/2$ and correspondingly any value of Δ [16]. To be as general as possible we will then consider Δ as a free parameter that plays the role of the SS parameter. Then the supersymmetry breaking gaugino and scalar masses for zero-modes are

$$M_{1/2} = M_0 = \frac{\Delta}{R} \tag{3.2}$$

provided they propagate in the bulk. In view of the problem with gaugino masses for vector multiplets localized on the branes, that we pointed out in the previous section, we will henceforth assume that those vector multiplets are bulk fields. On the other hand, localized scalars get their masses via gauge and Yukawa interactions from bulk fields.

The gauge interactions couple the fields in the boundary to the KK-towers of gauge bosons and gauginos. At the one-loop, we find that the boundary scalars get a mass given by [19],

$$\Delta_g m_i^2 = \frac{g^2 C(R_i)}{4\pi^4}\left[m^2(0) - m^2(\Delta)\right]\,, \tag{3.3}$$

where R_i is the representation of the gauge group under which the boundary field transforms, and $m^2(q)$ is defined by ($z \equiv {\rm e}^{i2\pi q}$)

$$m^2(q) = \frac{1}{2R^2}\left(Li_3(z) + Li_3(1/z)\right)\,. \tag{3.4}$$

The boundary field can also have Yukawa couplings to an $N=1$ chiral supermultiplet that consists in the KK-towers of a complex scalar and a bispinor. In this case, we find that the scalar field of the boundary gets a mass given by [19]

$$\Delta_{\rm Y} m_i^2 = \frac{Y^2}{16\pi^4}\left[m^2(2\Delta) + m^2(0) - 2m^2(\Delta)\right]\,, \tag{3.5}$$

where Y is the Yukawa coupling between the bulk and boundary fields. Finally, we have calculated the contribution of the KK-towers to a scalar trilinear coupling, A, between two boundary fields, Q and U, and one field in the bulk. This contribution arises from gaugino loops and gives

$$\Delta A = \frac{Y g^2 T_Q^a T_U^a}{8\pi^3} A(\Delta)\,, \tag{3.6}$$

where

$$A(\Delta) = i\left[Li_2(\mathrm{e}^{i2\pi\Delta}) - Li_2(\mathrm{e}^{-i2\pi\Delta})\right]/R\,, \tag{3.7}$$

and T_R^a is the generator of the gauge group in the representation of R.

The resulting supersymmetric spectrum predicted by this kind of scenarios is characteristic. One of its main properties is that supersymmetry breaking is carried by gauge and Yukawa interactions, and it automatically solves the supersymmetric flavor problem.

3.2. Electroweak symmetry breaking

It has been recognized that in models with extra dimensions electroweak breaking is triggered at one-loop if the top quark is localized on the brane. In that case one gets the additional bonus that it yields a finite mass term. This can be understood since, even if the 5D theory is non-renormalizable, the power-law divergence is canceled by supersymmetry. After subtracting the 5D part it remains a finite piece that is cutoff by $1/R$, as it happens in field theory at finite temperature for thermal masses. A very simple example was worked out in detail in Ref. [20] where quark and lepton SU(2) singlets propagate in the bulk while all doublets (including the Higgs sector of the MSSM) are localized in the branes. In this way one can construct the MSSM superpotential localized on the physical brane and one expects that the top-quark Yukawa coupling will dominate over the other Yukawa and gauge couplings and trigger electroweak symmetry breaking.

The Higgs potential along the direction of the neutral components of the fields H_2 and H_1 contains mass terms $m_i^2|H_i|^2$, $i=1,2$, as well as the mixing mass $m_3^2 H_1 H_2 +$ h.c. The supersymmetric tree-level relations $m_1 = m_2 = \mu$ and $m_3^2 = 0$, where μ is the supersymmetric Higgsino mass term, hold. These relations are spoiled by radiative corrections which provide contributions to all the above parameters. These corrections are driven by the $\mathrm{SU}(2)_L \times \mathrm{U}(1)_\mathrm{Y}$ gauge couplings g and g', and by the top and bottom Yukawa couplings, defined as:

$$v\, h_\mathrm{t} = m_\mathrm{t}\sqrt{1+t_\beta^2}\,/\,t_\beta\,, \quad v\, h_\mathrm{b} = m_\mathrm{b}\sqrt{1+t_\beta^2}\,, \tag{3.8}$$

where $t_\beta \equiv \tan\beta \equiv v_2/v_1$, $v_i = \langle H_i \rangle$ are the vacuum expectation values of the Higgs fields, $v = \sqrt{v_1^2+v_2^2} = 174.1$ GeV, and m_t and m_b are the top and bottom running masses. Notice that $h_{\rm b}$ can become important only for large values of t_β, as those that will be found by minimization of the one-loop effective potential. We will only consider the leading radiative corrections.

All radiative corrections to the potential parameters depend on $1/R$ and Δ. In particular the one-loop radiative corrections to the mass of any scalar localized on the brane were computed in (3.3) and (3.5). A simple application to the Higgs mass terms m_1^2 and m_2^2 yields:

$$\begin{aligned} \Delta m_2^2 &= \mu^2 - \frac{6h_{\rm t}^2 - 3g^2}{32\pi^4}\frac{f(\Delta)}{R^2}\,, \\ \Delta m_1^2 &= \mu^2 - \frac{6h_{\rm b}^2 - 3g^2}{32\pi^4}\frac{f(\Delta)}{R^2}\,, \end{aligned} \tag{3.9}$$

where the function $f(\Delta)$ is defined as $[r = \exp(2\pi i\Delta)]$

$$f(\Delta) = 2\zeta(3) - [Li_3(r) + Li_3(1/r)]\ . \tag{3.10}$$

The mass term m_3^2 is generated by the one-loop diagram exchanging KK-modes of gauginos, $\lambda^{(n)}$, and localized Higgsinos, $\widetilde{H}_{1,2}$. The resulting contribution is given by

$$m_3^2 = \mu\frac{3\,g^2}{512\pi^2}\frac{\Delta}{R}\left[i\,Li_2(r) - i\,Li_2(1/r)\right]\ . \tag{3.11}$$

Notice that, for the particular case $\Delta = 1/2$ $(r=-1)$, $m_3^2 = 0$ reflecting the fact that the gauginos $\lambda^{(n)}$ are, in that case, Dirac fermions.

We can see that the negative contributions to the Higgs mass in (3.9) trigger electroweak symmetry breaking. The detailed predictions for the Higgs mass spectrum are of course model dependent but can always be described from the point of view of the MSSM parameter space. In particular for the model described in Ref. [20] the typical predictions are large $\tan\beta$, more precisely $\tan\beta \sim m_{\rm t}/m_{\rm b}$, and light pseudoscalar m_A (since m_3^2 only gets a radiative contribution), and therefore one gets a Higgs spectrum corresponding to these values.

4. String and M-theory implementation

So far the above analysis was performed at the level of an effective field theory. This is expected to arise as a low energy limit of M-theory or type I string vacua. The particularity of compactifications down to four dimensions of the Hořava–Witten M-theory vacua is the presence of one (the eleventh) dimension where no gauge degrees of freedom propagate. It is a natural framework for implementing the ideas discussed in Section 2 [11, 12]. In more realistic situations, the analysis is somehow more involved [13] because of the fact that the internal space does not factorize into a product of a segment by a Calabi–Yau threefold [2]. These compactifications include either standard [2] or non-standard [21] embedding of the spin connection in the gauge connection. Analysis of the resulting field theory in five dimensions and of the possible role of different fields in mediating supersymmetry breaking have been considered by Stefan Pokorski and collaborators [22].

Type II compactifications allow, on the other hand, to consider the corresponding supersymmetry breaking at the perturbative string theory level. The hope being that existing conformal field theory techniques would in such cases provide useful tools to explore the explicit dependence of the resulting effective Lagrangian parameters on the details of the string vacuum.

In such constructions closed strings describe gravity while gauge interactions are described by open strings with ends bounded to propagate on Dp-branes. The six internal compact dimensions split into the (p-3) longitudinal and the (9-p) transverse ones. Because gravitational and gauge interactions appear at different orders in string perturbation theory the string scale can be lowered to the TeV scale, of the same order of magnitude as the inverse longitudinal dimensions, in which case the transverse dimensions should be much larger, as large as the submillimeter scale, and sensitive to gravitational experiments. In this kind of scenarios the Standard Model fields are described by open strings with ends at (p+1)-hypersurfaces embedded in the 10D space-time, or Dp-branes. The p-3 compact dimensions are forming in the simplest cases tori or orbifolds.

While the brane states are "localized" inside the bulk, there is the possibility to have some states localized inside the brane itself. They appear located along its intersection with other branes. For instance, the case in Section 2 corresponds to D3-branes localized at the boundary of an S^1/Z_2 segment, while in Section 3, we add D4-branes spanning this segment and identify instead the gauge fields of the Standard Model with the lowest excitations of open strings propagating on the D4-branes.

Supersymmetry can be broken along a direction S^1/Z_2 transverse to 3-branes. It can then be shown that on one end of the segment, one half of the bulk supersymmetry is realized with the appearance of the associated orientifold planes $O3$, while anti-orientifold planes $\bar{O}3$ carrying opposite Ramond–Ramond (RR) charges and conserving the other half of the bulk supersymmetry are located on the other end of the segment. Cancellation of Neuveu–Schwarz (NS–NS) tadpoles, necessary to obtain a four dimensional flat Minkowski space, require to add an appropriate number of pairs of branes–antibranes. There are two cases:

- In the one case, corresponding to $D3$ on top of $O3$ and $\bar{D}3$ on top of $\bar{O}3$, only the bulk states feel the supersymmetry breaking at tree-level [23].
- The other case with $\bar{D}3$ on top of $O3$ and $D3$ on top of $\bar{O}3$, leaving non-vanishing RR local tadpoles, the massless states on each of the branes are no more degenerate between fermions and bosons [24]. In fact, it was found that the supersymmetry is non-linearly realized. This is similar to the situation with NS–NS tadpoles [25].

Explicit models have been constructed in this way and some of their main features are still under study.

5. Higgs mass in non-supersymmetric scenarios

One of the main motivations for considering the quantum gravity scale to lie in the TeV range is to provide an alternative to supersymmetry when dealing with the problem of gauge hierarchy. It is then important to consider the fate of the Higgs mass in explicit realizations of this scenario.

We consider first a simple case where the whole one-loop effective potential of a scalar field can be computed. We suppose d (large) extra dimensions compactified on orthogonal circles with radii $R_i > 1$ (in units of the string length $l_{\rm s} \equiv M_{\rm s}^{-1}$) with $i = 1, \ldots, d$. An interesting situation is provided by a class of models where a non-vanishing VEV for a scalar (Higgs) field ϕ results in shifting the mass of each KK excitation by a constant $a(\phi)$:

$$M_{\vec{m}}^2 = \sum_{i=1}^{d} \left[\frac{m_i + a_i(\phi)}{R_i}\right]^2 , \tag{5.1}$$

where $\vec{m} = \{m_1, \cdots, m_d\}$ with m_i integers. Such mass shifts arise for instance in the presence of Wilson lines, $a_i = q \oint (dy^i)/(2\pi) g A_i$, where A_i is the internal component of a gauge field with gauge coupling g and q is the charge of the given state under the corresponding generator. A straightforward computation shows that the ϕ-dependent part of the one-loop effective potential is given by [26]:

$$V_{\rm eff} = -{\rm Tr}(-)^F \, \frac{\prod_{i=1}^{d} R_i}{32\,\pi^{\frac{4-d}{2}}} \, \sum_{\vec{n}} {\rm e}^{2\pi i \sum_i n_i a_i} \int_0^\infty dl \; l^{\frac{2+d}{2}} f_{\rm s}(l) \; {\rm e}^{-\pi^2 l \sum_i n_i^2 R_i^2} \tag{5.2}$$

where $F = 0, 1$ for bosons and fermions, respectively. We have included a regulating function $f_{\rm s}(l)$ which contains for example the effects of string oscillators. To understand its role we will consider the two limits $R_i > 1$ and $R_i \ll 1$. In the first case only the $l \to 0$ region contributes to the integral.

This means that the effective potential receives sizable contributions only from the infrared (field theory) degrees of freedom. In this limit we would have $f_{\rm s}(l) \to 1$. For example, in the string model considered in [27]:

$$f_{\rm s}(l) = \left[\frac{1}{4l}\frac{\theta_2}{\eta^3}(il+\frac{1}{2})\right]^4 \to 1 \qquad \text{for} \qquad l \to 0, \tag{5.3}$$

and the field theory result is finite and given by:

$$V_{\rm eff}(\phi) = -{\rm Tr}(-)^F \; \frac{\Gamma(\frac{4+d}{2})}{32\pi^{\frac{12+d}{2}}} \prod_{i=1}^{d} R_i \sum_{\vec{n}\neq\vec{0}} \frac{{\rm e}^{2\pi i \sum_i n_i a_i(\phi)}}{\left[\sum_i n_i^2 R_i^2\right]^{\frac{4+d}{2}}} \,. \tag{5.4}$$

As a result of the Taylor expansion around $a_i = 0$, we are able to extract the finite one-loop contribution to the coefficient of the term of the potential quadratic in the Higgs field. It is given by a loop factor times the compactification scale [26]. For instance, in the case of $d = 1$ dimension, one obtains $\mu^2 \sim g^2/R^2$ up to a proportionality constant which is calculable in the effective field theory.

On the other hand, if we consider $R_i \to 0$, which by T-duality corresponds to taking the extra dimensions as transverse and very large, the one-loop effective potential receives contributions from the whole tower of string oscillators as appearing in $f_{\rm s}(l)$ leading to squared masses given by a loop factor times $M_{\rm s}^2$:

$$\mu^2 = \pm\varepsilon^2 g^2 M_{\rm s}^2 \,. \tag{5.5}$$

The precise numerical coefficient ε^2 is sensitive to details of the considered string model. The sign has been found in constructing string examples to be given by the difference between the number of light fermions and bosons.

We turn now to the realistic case of the Standard Model Higgs field. We first focus on the case with radii $R_i > 1$ *i.e.* with longitudinal directions where all fields of the Standard Model propagate. A finite result, due to the mass spectrum described above with a constant shift for all KK states, can be obtained if the Higgs field is identified with the internal component of a gauge field extending the Standard Model in higher dimensions. The minimal extension is U(3)× U(3). Unfortunately, in the case of one extra dimension, the tree-level quartic interaction term is absent leading to an unacceptably small Higgs mass ($\sim$ 50 GeV). Therefore, we are led to consider $d > 1$ which in turn leads generically to extra Higgs fields corresponding to the different internal components of the gauge fields. Moreover, unless the quartic term is absent, in which case there is no improvement compared to $d = 1$, the Higgs fields are not flat directions (Wilson lines) and the computation of the one-loop effective potential above does not apply. However, it is possible to show that the squared masses of the Higgs fields can be extracted from the above formulae. An explicit computation was performed for a compactification on a T^2/Z_2 orbifold of a six-dimensional gauge theory U(3)$\times$ U(3) with massless matter fields transforming only in the representations of the Standard Model [26].

In the case where $R_i \to 0$, the Higgs mass is given by (5.5). Moreover, in the situation where the Higgs arises from open strings ending on parallel D-branes, the Higgs quartic coupling is related to the gauge coupling. It is then possible to extract the relation $M_{\rm s} = M_h/\sqrt{2}g\varepsilon$ between the Higgs mass M_h, the string scale $M_{\rm s}$ and the parameter ε, in order to achieve the correct electroweak breaking scale [27].

One of us (MQ) would like to thank the Physics and Astronomy Department of the Johns Hopkins University, where part of this work has been done, for the hospitality extended to him as Bearden Visiting Professor. This work is supported in part by EU under RTN contracts HPRN-CT-2000-00148 and HPRN-CT-2000-00152, in part by IN2P3-CICYT contract Pth 96-3 and in part by CICYT, Spain, under contract FPA2001-1806.

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