Soft Masses in Theories with Supersymmetry Breaking by TeV-Compactification

I. Antoniadis[a], S. Dimopoulos[b], A. Pomarol[c†] and M. Quirós[d]

[a] *Centre de Physique Théorique (CNRS UMR 7644), Ecole Polytechnique, F-91128 Palaiseau*

[b] *Physics Department, Stanford University, Stanford, California 94309, USA*

[c] *TH-Division, CERN, CH-1211 Geneva 23, Switzerland*

[d] *Instituto de Estructura de la Materia, CSIC, Serrano 123, E-28006 Madrid*

Abstract

We study the sparticle spectroscopy and electroweak breaking of theories where supersymmetry is broken by compactification (Scherk-Schwarz mechanism) at a TeV. The evolution of the soft terms above the compactification scale and the resulting sparticle spectrum are very different from those of the usual MSSM and gauge mediated theories. This is traced to the softness of the Scherk-Schwarz mechanism which leads to scalar sparticle masses that are only logarithmically sensitive to the cutoff starting at two loops. As a result, squarks and sleptons are naturally an order of magnitude lighter than gauginos. In addition, the mechanism is very predictive and the sparticle spectrum depends on just two new parameters. A significant advantage of this mechanism relative to gauge mediation is that a Higgsino mass $\mu \sim M_{\mathrm{susy}}$ is automatically generated when supersymmetry is broken. Our analysis applies equally well to theories where the cutoff is near a TeV or $M_{P\ell}$ or some intermediate scale. We also use these observations to show how we may obtain compactification radii which are hierarchically larger than the fundamental cutoff scale.

[†] On leave from: IFAE, Universitat Autònoma de Barcelona, E-08193 Bellaterra, Barcelona

1. Introduction

Theories with supersymmetry softly broken at the weak scale have been the most popular approach to the hierarchy problem for the last seventeen years [1]. In spite of this, the problem of supersymmetry breaking still remains an open question. Perhaps this is not surprising since supersymmetry breaking is intimately connected to the cosmological constant problem, the most difficult problem in physics. So far there have been two scenaria that have been suggested. One is that there is some high scale physics of perhaps gravitational or GUT origin leading to the soft supersymmetry breaking terms [1], and is often referred to as the "gravity mediated scenario". The second postulates that supersymmetry breaking originates in a sector with which we share gauge interactions and is called the "gauge mediated scenario" [2]. There have been a lot of works in the literature dealing with the phenomenological consequences of either gravity or gauge mediated scenaria over the last seventeen years.

Recently a third daring possibility was suggested [3]-[6] which is based on breaking supersymmetry by compactification (Scherk-Schwarz mechanism) [7, 8], and involves new TeV-size spatial dimensions. The theoretical viability of this possibility is far from obvious, since it is embedded in theories that live in higher dimensions from the TeV to the Planck scale. In fact, large dimensions are a general prediction of any (known) perturbative description of supersymmetry breaking in string theory, which necessarily relates the breaking scale to the size of some compact dimension(s) [9].

More recently there has been a new proposal in which the hierarchy problem is solved by lowering the Planck scale down to a TeV [10, 11]. This raises the possibility that the Scherk-Schwarz mechanism for supersymmetry breaking may be embedded in theories without any severe ultraviolet (UV) problems. However, as we shall discuss in the last section, in this case one has to face the problem of the cosmological constant in the bulk which is generically much larger than the vacuum energy on the (observable) wall.

These cautionary remarks are intended to underline that there is no well established framework in which the Scherk-Schwarz mechanism for supersymmetry breaking is embedded. However, leaving apart the cosmological constant problem, it may be premature to dismiss the Scherk-Schwarz mechanism as a viable possibility for supersymmetry breaking. Moreover, as we shall show in this paper, the spectroscopy and experimental signatures of Scherk-Schwarz supersymmetry breaking (SSSB) are distinct and drastically different than either gravity or gauge mediation.

An important aspect of SSSB is that it is totally analogous to the breaking of supersymmetry via temperature, where the role of temperature is played by the inverse of the compactification radius $1/R$. This implies that supersymmetry breaking quantities are UV-insensitive. This is a consequence of the exponential Boltzmann suppression factors which suppress the contribution of any high energy level to a thermodynamic quantity. In practice it means that supersymmetry breaking physics in SSSB will only depend on physics up to the compactification scale and will not be sensitive to what happens beyond it. This is especially welcome since the theory above the compactification scale is intrinsically higher dimensional and not treatable with standard field theory tools. The beauty of the SSSB is that it is insensitive to the higher dimensional theory as far as soft terms and supersymmetry breaking parameters are concerned. Thus, in SSSB the supersymmetry breaking parameters are under better control than supersymmetry preserving quantities, such as gauge and Yukawa couplings.

An important corollary of this, pointed out in ref. [5], is that the cosmological constant does not have quadratic divergences or, equivalently, quadratic sensitivity to the UV cutoff. This is analogous to the fact that the free energy at finite temperature is proportional to T^4 and has no quadratic divergences either. Again, this follows from the exponential suppression of high energy states' contribution to the free energy. A testable consequence of this behavior is the existence of light gravitationally coupled "moduli" with $\sim sub-$

millimeter wavelengths [12, 5].

The aim of this paper [1] is to study the implications of the intrinsic softness of the SSSB mechanism for the sparticle spectroscopy. We will show that the resulting sparticle spectrum markedly differs from the gauge or gravity mediated case and leads to much larger hierarchies of the scalar and gaugino masses without any fine tuning. This will be done in Sections 3 and 4. The paper is organized as follows. In Section 2, we present our general framework and discuss the role of extra dimensions and the issue of gauge coupling unification. In Section 3, we review the method of supersymmetry breaking by SSSB compactification and compute the one-loop corrections to the soft terms. In Section 4, we study the breaking of the electroweak symmetry and discuss the resulting superparticle spectrum and its properties.[2] In Section 5, we show how the softness of SSSB can be useful in attempts to dynamically relate the compactification and fundamental scale (cutoff) in a hierarchical way. Finally, Section 6 contains some concluding remarks.

2. Large dimensions and unification

Here, we consider a supersymmetric extension of the standard model with an extra dimension y compactified on a line interval S^1/Z_2, obtained upon identification under $y \to -y$ of the points of a circle S^1 of radius $R \sim \text{TeV}^{-1}$. In general, the Z_2 discrete symmetry acts also non trivially on the 5-dimensional (5D) fields. These theories have two types of matter states: the bulk (untwisted) ones, living with the gauge fields in the 5D bulk, and the boundary (twisted) ones, that are localized at the two fixed points of the orbifold, $y = 0$ and $y = \pi R$.

---

[1]The main results contained in this work have been presented at the SUSY 98 Conference, Oxford (11-17 July 1998) [13].

[2]A previous attempt to study the phenomenology of SSSB [4] did not take into account the extreme softness of the soft breaking terms.

The massless spectrum has $N = 1$ supersymmetry in four dimensions and should contain the MSSM particles. The massive spectrum, however, forms towers of Kaluza-Klein (KK) excitations for all the fields living in the 5D bulk, with masses n/R for $n = 0, 1, \ldots$; they fall into supermultiplets of extended $N \geq 2$ supersymmetry. We will distinguish two different cases:

(a) The KK modes are organized either in $N = 4$ supermultiplets, or in $N = 2$ but are falling into appropriate group representations leading to vanishing beta-functions. For brevity, we shall refer to this case as the $N = 4$ one.

(b) The KK modes form just $N = 2$ multiplets.

On the other hand, obviously, the 4D boundary fields have no KK excitations.

Following refs. [4, 6], we will consider that in addition to gauge multiplets only the Higgs fields live in the 5D bulk, as a part of $N = 2$ vector supermultiplets or hypermultiplets. In this way, the μ-problem is automatically solved, since as we shall see in the next section, Higgsinos acquire a mass of the order of the compactification scale by the Scherk-Schwarz mechanism of supersymmetry breaking. On the other hand, quarks and leptons chiral multiplets are assumed to be localized in the 4D boundary.

In the $N = 4$ case (a), there is no contribution from the KK states to the beta-function coefficients of the gauge couplings. The only contribution arises from the zero modes and the twisted states of the 4D boundary. Therefore, the gauge couplings evolve as in the MSSM and unify at $M_{st} \simeq 10^{16}$ GeV. In general, in these models, additional constraints have to be imposed in order to avoid potential growing of the Yukawa couplings. For instance, when the quarks and leptons are also bulk fields and come from $N = 2$ hypermultiplets, this condition is automatically satisfied, since their wave function is not renormalized. Of course, in this case, special model building is needed to satisfy the condition of vanishing of the 1-loop $N = 2$ beta-functions. When quarks and leptons are twisted fields, the Yukawa coupling constraints are satisfied if for instance there are non-trivial infrared-stable fixed

points in the full theory [3].

In the generic $N = 2$ case (b), the KK states contribute to the gauge beta-function coefficients, and change the logarithmic scale dependence of the gauge coupling to a power-law running. This accelerates the gauge couplings evolution and they may unify at much lower energies [14]. Assuming that only the gauge and Higgs fields live in 5D, and that supersymmetry is broken by the Scherk-Schwarz mechanism along the orbifold compactification at the scale $M_c \equiv 1/R$, one has at one loop level

$$\frac{1}{\alpha_i(m_Z)} = \frac{1}{\alpha_{st}} + \frac{b_i^{SM}}{2\pi} \ln \frac{M_c}{m_Z} + \frac{b_i^{MSSM}}{2\pi} \ln \frac{M_{st}}{M_c} + \frac{b_i^{KK}}{2\pi} \left[\frac{M_{st}}{M_c} - 1 - \ln \frac{M_{st}}{M_c} \right] + \Delta_i \,, \quad (2.1)$$

where $b^{SM} = (41/10, -19/6, -7)$, $b^{MSSM} = (66/10, 1, -3)$ and $b^{KK} = (3/5, -3, -6)$ are respectively the beta-function coefficients of the Standard Model (SM), MSSM and KK states, while Δ_i denote additional string threshold corrections.

In the absence of threshold corrections, it turns out that the measured values of the three gauge couplings at m_Z lead to an approximate unification (within 2% taking into account the experimental errors) with

$$M_{st} \simeq 45 \text{ TeV} \qquad \alpha_{st}^{-1} \simeq 50 \,, \quad (2.2)$$

for $M_c \simeq 1$ TeV. However, as one can see from eq. (2.1), the gauge couplings acquire a power law sensitivity with respect to the UV cutoff M_{st}. As a result, the gauge coupling unification conditions are extremely sensitive to string threshold corrections [15]. This casts doubts on the significance of this calculation. It shows that unification of couplings cannot be decided by a low energy effective-theory computation; it requires a detailed knowledge of the full UV theory.

It is important to point out the relevance of the underlying string theory in both cases we discussed above, and mainly in the $N = 4$ case. Unlike gauge couplings whose quantum corrections are constrained by holomorphicity, physical amplitudes cannot in general be

---

[3]This possibility was pointed out to us by Nima Arkani-Hamed.

reliably computed beyond the compactification scale, as the effective field theory becomes non-renormalizable (higher dimensional). The full string theory is then required above M_c to fix the coefficients of the higher dimensional effective operators. However, as we will see in the next section, this "pathology" does not hold for the effective couplings that are generated after supersymmetry breaking, due to the extreme softness of the SSSB mechanism.

3.   Soft terms in Scherk-Schwarz compactification

Theories with extra dimensions, and in particular five dimensional theories, allow the use of the Scherk-Schwarz (SSSB) mechanism to break supersymmetry [7, 6]. This consists in imposing to the 5D fields a different periodicity condition for bosons and fermions under a $2\pi R$ translation of the extra dimension:

$$\Phi(x^\mu, y + 2\pi R) = e^{2\pi i q_\Phi} \Phi(x^\mu, y) , \qquad (3.1)$$

where q_Φ is the R-symmetry charge of the field Φ. Due to the periodicity condition (3.1), the fields are Fourier expanded as

$$\Phi(x^\mu, y) = \sum_{n=-\infty}^{\infty} e^{iy(n+q_\Phi)/R} \, \Phi^{(n)}(x^\mu) , \qquad (3.2)$$

where y is assumed to be compactified on the circle S^1. Reducing the theory from 5D to 4D, eq. (3.2) leads to a tower of KK states with a fermion-boson mass splitting inside each KK supermultiplet:

$$
\begin{aligned}
m_B^2 &= (n + q_B)^2 \, M_c^2 , \\
m_F^2 &= (n + q_F)^2 \, M_c^2 , \quad n = 0, \pm 1, \pm 2, ... ,
\end{aligned}
\qquad (3.3)
$$

where $M_c \equiv 1/R$, and q_B and q_F are the charges of the bosons and fermions, respectively.

In orbifold compactifications, the expansion (3.2) is truncated, since only the invariant states under the orbifold group remain in the theory. For example in S^1/Z_2 compactifications, the Z_2 parity, $y \to -y$, acts on the KK-states as $\Phi^{(n)} \to \Phi^{(-n)}$. Therefore the Z_2 projects the KK-tower into Z_2-even ($\Phi_+^{(n)}$) or Z_2-odd ($\Phi_-^{(n)}$) states:

$$\Phi_+^{(n)} = \Phi^{(n)} + \Phi^{(-n)}, \qquad n = 0, 1, 2, \dots,$$

$$\Phi_-^{(n)} = \Phi^{(n)} - \Phi^{(-n)}, \qquad n = 1, 2, \dots. \tag{3.4}$$

Massless fields arise in the theory only if either q_B or q_F are zero. We are interested in the $q_B = 0$ case, in which only the vector and scalar boson $n = 0$ states remain in the massless spectrum of the theory, and can be associated with the bosonic sector of the SM. Although vector bosons remain massless because of the gauge symmetry, the $n = 0$ complex scalar field, ϕ, will get a mass at the one-loop level due to the breaking of supersymmetry by the SSSB mechanism. Each level of KK excitations contribute to the scalar mass. This is a crucial difference with respect to 4D theories with softly broken supersymmetry.

The one-loop m_ϕ^2 induced by a tower of KK with a mass splitting (3.3) can be obtained from the effective potential $V(\phi)$:

$$m_\phi^2 = \left. \frac{dV(\phi)}{d|\phi|^2} \right|_{\phi=0}, \tag{3.5}$$

with $V(\phi)$ given by

$$V(\phi) = \frac{1}{2} \text{Tr} \sum_{n=-\infty}^{\infty} \int \frac{d^4p}{(2\pi)^4} \ln \left[\frac{p^2 + (n + q_B)^2 M_c^2 + M^2(\phi)}{p^2 + (n + q_F)^2 M_c^2 + M^2(\phi)} \right], \tag{3.6}$$

where the trace is over the degrees of freedom of the KK tower and $M^2(\phi)$ is the ϕ-dependent mass of the KK states. From eq. (3.5), we obtain

$$m_\phi^2 = \frac{1}{2} \text{Tr} \left. \frac{dM^2(\phi)}{d|\phi|^2} \right|_{\phi=0} \sum_{n=-\infty}^{\infty} \Pi_n(0), \tag{3.7}$$

$$\Pi_n(0) = \int \frac{d^4p}{(2\pi)^4} \left[\frac{1}{p^2 + (n + q_B)^2 M_c^2} - \frac{1}{p^2 + (n + q_F)^2 M_c^2} \right]. \tag{3.8}$$

As in finite temperature calculations, we must first sum over the infinite tower of KK states and then perform the momentum integral. Eq. (3.8) thus leads to [4]

$$m_\phi^2 = \frac{1}{32\pi^4} \left[\Delta m^2(q_B) - \Delta m^2(q_F) \right] \text{Tr} \left. \frac{dM^2(\phi)}{d|\phi|^2} \right|_{\phi=0}, \tag{3.9}$$

where

$$\Delta m^2(q) = \frac{1}{2} (Li_3(z) + Li_3(1/z)) M_c^2, \tag{3.10}$$

with $z \equiv e^{i2\pi q}$ and $Li_n(z)$ are the polylogarithm functions $Li_n(z) = \sum_{k=1}^{\infty} \frac{z^k}{k^n}$. For q ranging from 0 to 1/2, $(Li_3(z) + Li_3(1/z))/2$ goes from 1.2 to -0.9.

We have just considered a theory compactified on S^1. For a compactification on the S^1/Z_2 orbifold, the contribution to m_ϕ^2 from a KK-tower must sum from $n = 0$ to ∞ (see eq. (3.4)). Nevertheless, the contribution to m_ϕ^2 of an even field (Π_n^+) can be always added to the one of an odd field (Π_n^-) such that they can be considered as the contribution of a single KK-tower where n goes from $-\infty$ to ∞ [6]:

$$\sum_{n=0}^{\infty} \Pi_n^+ + \sum_{n=1}^{\infty} \Pi_n^- = \sum_{n=-\infty}^{\infty} \Pi_n, \tag{3.11}$$

where $\Pi_{\pm n} \equiv \Pi_n^\pm$. Thus, in S^1/Z_2 orbifold compactifications, the effective number of states of the KK-tower is reduced by a half with respect to that in S^1.

The contribution (3.9) is finite and ultraviolet independent. There are different ways to understand this result. The simplest way is to notice that a theory with SSSB supersymmetry breaking keeps a clear analogy with a theory at finite temperature T (a quantum field theory with the time compactified) where M_c plays the role of T. At finite temperature, the T-dependent effective potential is not affected by the ultraviolet cutoff due to the Boltzmann suppression of the heavy states. Of course, the T-independent part of the effective potential is ultraviolet sensitive. In our case, however, $M_c = 0$ corresponds to the supersymmetric limit and therefore the corrections to $V(\phi)$ are zero. Alternatively, the

---

[4]For further details see ref. [16].

insensitivity of m_ϕ^2 to an ultraviolet scale $\Lambda \gg M_c$ (where new physics arises) can be probed by just putting an explicit ultraviolet cutoff in our theory and showing that the result (3.9) is not modified. This is in analogy with Casimir energy calculations [17]. For example, we can insert in the sum of eq. (3.7) a function $f(\Lambda, nM_c)$ that goes to zero for $nM_c \gg \Lambda$. The function f must be normalized such that $f \to 1$ for $n \to 0$. For example, we can take $f = e^{-nM_c/\Lambda}$. As expected, we find that the calculation of m_ϕ^2 gives, for $\Lambda \gg M_c$, the same result (3.9).

Using eq. (3.9) we can calculate the contribution to the soft mass of any scalar field of the theory. Let us take as an example, the model of ref. [6] in which after compactification on S^1/Z_2 a massless scalar ϕ arises from a 5D hypermultiplet. For simplicity, let us assume $q_F = 1/2$ and $q_B = 0$. The fermionic Z_2-even sector consists of two bispinors, a gaugino and the partner of ϕ, that combine with the odd sector to form two full KK-towers. We then have 4 degrees of freedom and $\mathrm{Tr}\,M^2(\phi) = 8g^2 C(\phi)|\phi|^2$ where $C(\phi)$ is the quadratic Casimir of the scalar ϕ in the corresponding gauge group $[C(\mathbf{N}) = (N^2 - 1)/(2N)$ for the fundamental representation of SU(N)]. Therefore

$$\mathrm{Tr}\,\left.\frac{dM^2(\phi)}{d|\phi|^2}\right|_{\phi=0} = 8g^2 C(\phi)\,. \tag{3.12}$$

By supersymmetry eq. (3.12) must hold both for bosons and fermions. Using eqs. (3.9) and (3.12), we obtain

$$m_\phi^2 = \frac{7g^2 C(\phi)\zeta(3)}{16\pi^4} M_c^2 \simeq 5 \times 10^{-3} M_c^2\,, \tag{3.13}$$

where for the numerical estimate we have taken $g^2 C(\phi) \sim 1$ and $\zeta(3) \simeq 1.2$. Thus, the scalar remains around an order of magnitude lighter than the compactification scale.

Let us now consider the fields living in the 4D boundary. These fields do not have associated KK excitations and are massless at tree level. Nevertheless, if they couple to the fields living in the 5D bulk, the supersymmetry breaking will be transmitted from the bulk to the boundary and, as a consequence, the scalars living in the boundary will get masses at the one-loop level. Let us consider again the case of an S^1/Z_2 orbifold. The possible

interactions between the bulk and the boundary fields can be found in ref. [18]. The gauge interactions couple the fields in the boundary to the KK-towers of the gauge boson, gaugino and the auxiliary D-field. At the one-loop, we find that the boundary scalars get a mass given by (for $q_B = 0$)

$$m_i^2 = \frac{g^2 C(R_i)}{4\pi^4} \left[\Delta m^2(0) - \Delta m^2(q_F) \right] , \tag{3.14}$$

where R_i is the representation of the gauge group under which the boundary field transforms, and $\Delta m^2(q)$ is given in eq. (3.10). The boundary field can also couple to an $N = 1$ chiral supermultiplet that consists in the KK-towers of a complex scalar, a bispinor and the auxiliary F-field. In this case we find that the scalar field of the boundary gets a mass given by

$$m_i^2 = \frac{Y^2}{16\pi^4} \left[\Delta m^2(q_B) + \Delta m^2(2q_F - q_B) - 2\Delta m^2(q_F) \right] , \tag{3.15}$$

where Y is the Yukawa coupling between the bulk and boundary fields. The first term in eq. (3.15) arises from the scalar KK-tower, the second term from the auxiliary F-field KK-tower and the third from the bispinor KK-tower. Again these contributions are also finite and ultraviolet independent.

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

$$A = 4Y g^2 T_Q^a T_U^a \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2} \sum_{n=-\infty}^{\infty} \frac{(n+q_F)M_c}{p^2 + (n+q_F)^2 M_c^2} , \tag{3.16}$$

that leads to

$$A = \frac{Y g^2 T_Q^a T_U^a}{8\pi^3} \Delta A(q_F) , \tag{3.17}$$

where

$$\Delta A(q_F) = i \left[Li_2(e^{i2\pi q_F}) - Li_2(e^{-i2\pi q_F}) \right] M_c , \tag{3.18}$$

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

Note that in contrast to the case of supersymmetric parameters, as gauge and Yukawa couplings that exhibit generically a power dependence on the UV cutoff, the soft breaking parameters are insensitive to it. This is due to the extreme softness of the SSSB mechanism which dies off exponentially in the supersymmetric limit, in analogy to the situation at finite temperature. This behavior persists, as well, for all couplings of higher dimensional operators induced by the supersymmetry breaking. In fact, it is easy to see that these couplings vanish in the supersymmetric limit and they are suppressed by powers of M_c/M_{st}.

Let us finally comment on higher loop effects. In analogy with the finite temperature calculation, we can implement some of these effects by replacing the couplings g and Y in the above calculations by the running couplings at the scale M_c, $g(M_c)$ and $Y(M_c)$. The effect of the boundary fields, however, does not have an analogy with finite temperature. The fields in the boundary contribute at the one-loop level to the wave function renormalization constant of the KK-excitations. This contribution will make the KK masses to evolve logarithmically with the renormalization scale. As a consequence, at higher-loop orders, M_c^2 must be replaced by the renormalized M_c^2 (see below eq. (5.3)).

4.   Superparticle spectrum and electroweak symmetry breaking

Sspectroscopy and the LSP

Let us apply the above calculation to our model, where gauginos and Higgsinos in the bulk are given the same boundary conditions, q_F, corresponding to a common R-symmetry charge. A more general case [6] will be treated in [16]. Since gauge and Higgs bosons live in the 5D bulk their corresponding $n = 0$ fermions, gauginos and Higgsinos, will get masses of order M_c:

$$m_\lambda \;=\; q_F \, M_c \,, \tag{4.1}$$

$$m_{\widetilde{H}} \;\; = \;\; q_F \, M_c \,. \tag{4.2}$$

Therefore the massless states of the 5D bulk correspond to the gauge and Higgs bosons of the SM. Quarks and leptons superfields reside in the 4D boundary, and then they also remain massless at the tree-level. Nevertheless, since supersymmetry is broken in the bulk, squarks and sleptons will get masses at the one-loop level through the gauge and Yukawa interactions, leaving only the fermion sector of the SM in the massless spectrum. Using eqs. (3.14) and (3.15), we obtain the squarks and sleptons masses:

$$m_{\widetilde{Q}}^2 \;\; = \;\; \left(\frac{8}{6}\alpha_3 + \frac{3}{4}\alpha_2 + \frac{1}{60}\alpha_1\right)\Delta m_g^2 + \frac{1}{2}\alpha_t\Delta m_H^2\,, \tag{4.3}$$

$$m_{\widetilde{U}}^2 \;\; = \;\; \left(\frac{8}{6}\alpha_3 + \frac{4}{15}\alpha_1\right)\Delta m_g^2 + \alpha_t\Delta m_H^2\,, \tag{4.4}$$

$$m_{\widetilde{D}}^2 \;\; = \;\; \left(\frac{8}{6}\alpha_3 + \frac{1}{15}\alpha_1\right)\Delta m_g^2\,, \tag{4.5}$$

$$m_{\widetilde{L}}^2 \;\; = \;\; \left(\frac{3}{4}\alpha_2 + \frac{3}{20}\alpha_1\right)\Delta m_g^2\,, \tag{4.6}$$

$$m_{\widetilde{E}}^2 \;\; = \;\; \frac{3}{5}\alpha_1\Delta m_g^2\,, \tag{4.7}$$

where

$$\Delta m_g^2 = \left[\Delta m^2(0) - \Delta m^2(q_F)\right]/\pi^3\,, \tag{4.8}$$

and

$$\Delta m_H^2 = \left[\Delta m^2(0) + \Delta m^2(2q_F) - 2\Delta m^2(q_F)\right]/(2\pi^3)\,, \tag{4.9}$$

with $\Delta m^2(q)$ given in eq. (3.10).

The above equation gives us a very predictive spectrum for the squarks and sleptons. It only depends on two free parameters, q_F and M_c. Notice also that the above contributions are positive as they are necessary to avoid color or charge breaking. The ratio of masses is given by (for $q_F = 1/2$)

$$10\,m_{\widetilde{Q}} \simeq 10\,m_{\widetilde{D}} \simeq 10\,m_{\widetilde{U}} \simeq 25\,m_{\widetilde{L}} \simeq 40\,m_{\widetilde{E}} \simeq m_{\lambda,\widetilde{H}}\,. \tag{4.10}$$

Therefore the gauginos and Higgsinos are the heaviest supersymmetric particles and the right-handed slepton is the lightest (LSP) one. In an R-parity conserving theory, the right-handed slepton will be stable and will cross the detector leaving an ionizing track. It can be discovered by looking at anomalous ionization energy loss, dE/dx, in the tracking detector gas [19]. The actual experimental lower bound on its mass is 82.5 GeV [19]. Cosmological arguments all but exclude charged, stable particles with masses in the ~ 100 GeV to ~ 10 TeV range [20]. A significant number of these particles will survive annihilation with their antiparticles and, at the time of Nucleosynthesis, combine with other nuclei to form "heavy" hydrogen, helium, etc., leading to heavy versions of these atoms today. Searches for such anomalous isotopes put very strong limits on their fractional abundance –as small as 10^{-30} for hydrogen. There is a variety of other possible considerations showing the implausibility of stable-charged LSPs, including interstellar calorimetry –the thermodynamics of interstellar clouds [21]– as well as neutron stars [22].

One way out of this is to ensure that the charged LSP is unstable. A simple way to accomplish this is to postulate an R-parity breaking interaction. Indeed, if R-parity is violated, the right-handed slepton will decay into SM leptons. In fact, since the R-parity violating coupling is renormalizable, the slepton is expected to decay inside the detector and can be easily discovered [23]. Another possibility to avoid a right-handed slepton LSP occurs in theories with right-handed neutrinos [24]. These must have miniscule Yukawa coupling to account for the observed smallness of the neutrino masses. Since the right-handed sneutrinos are electroweak singlets and also have miniscule Yukawa couplings, they will get tiny masses of order of the (Dirac) neutrino masses in the sub-eV range and will be the LSP. We are not aware whether such an LSP passes all the necessary cosmological safety tests, but it does not seem to us to be obviously excluded.

Electroweak Breaking

To study the breaking of the electroweak symmetry, we must analyze the soft masses of the Higgs(es). In the MSSM, we need two Higgs $SU(2)_L$-doublets, H_1 and H_2, in order to give masses to all the fermions living in the 4D boundary. After imposing the supersymmetry breaking with the SSSB compactification, these two Higgses can either arise as two (tree-level) massless states, or as a massless and a massive state. This depends on the R-charges of the Higgses, and different possibilities have been proposed in refs. [4, 6]. The simplest case corresponds to having a unique massless $SU(2)_L$-doublet, ϕ, that will be responsible for the electroweak symmetry breaking. This massless field can be either associated to one of the MSSM Higgses or to a linear combination of the two, $\phi \equiv \cos\beta H_1 + \sin\beta H_2$. The mixing angle $\tan\beta$ is model dependent. For example, for the model of ref. [6] in which ϕ arises from a 5D hypermultiplet and corresponds to a flat direction of the D-terms, one has $\cos\beta = \sin\beta$. In ref. [4] the only massless mode is the Higgs that couples to the top-quark: this would correspond to $\sin\beta = 1$. Further scenarios will be considered in ref. [16].

Here we will be only interested in knowing whether ϕ gets a vacuum expectation value and therefore breaks the electroweak symmetry. For this purpose we must calculate the quantum corrections to its mass. Using eq. (3.9), we have

$$m_\phi^2(M_c) = \left(\frac{3}{4}\alpha_2 + \frac{3}{20}\alpha_1\right)\Delta m_g^2 \,. \tag{4.11}$$

This mass is positive. Nevertheless, we must also consider the correction to the Higgs mass due to the stop. This correction arises at the two-loop level but it is important since the stops are heavier than the Higgs. This is given by

$$m_\phi^2(m_Z) \simeq m_\phi^2(M_c) - \frac{3\alpha_2 m_t^2}{8\pi m_W^2}(m_{\tilde{Q}}^2 + m_{\tilde{U}}^2)\ln\frac{M_c^2}{m_{\tilde{Q}}^2} \,. \tag{4.12}$$

As in theories of gravity or gauge-mediation supersymmetry breaking, this contribution turns the Higgs mass to negative values and triggers the breaking of the electroweak sym-

metry. Imposing the minimization condition

$$-m_\phi^2(m_Z) = \frac{m_Z^2}{2} \cos^2 2\beta \,, \tag{4.13}$$

one can derive the value of the compactification scale M_c. Taking $q_F = 1/2$, $m_t \simeq 175$ GeV and $\cos 2\beta \sim 1$, we obtain

$$M_c \sim 3.7 \text{ TeV} \,, \qquad m_{\widetilde{Q}} \sim 400 \text{ GeV} \,, \qquad m_{\widetilde{E}} \sim 100 \text{ GeV} \,. \tag{4.14}$$

We must notice that the above prediction is quite sensitive to the values of m_t and α_3 that have large experimental uncertainties. We find some values for these parameters for which the two terms of the RHS of eq. (4.12) approximately cancel out, and consequently the supersymmetric spectrum turns to be much heavier.

Contrast with the MSSM and Gauge Mediation

Sparticle spectroscopy in SSSB is strikingly different from that of more familiar gauge mediated and MSSM. All fermionic sparticles are at least an order of magnitude heavier than the bosonic ones, a smoking gun for this framework. In particular, the gauginos and higgsinos are ~ 40 times heavier than the right-handed sleptons. A consequence of this, following from the present lower limit of 82.5 GeV on the mass of the right-handed sleptons, is that the compactification scale as well as the gaugino and higgsino masses must be no less than 3 TeV. Consequently, the KK excitations of ordinary particles would not be accessible at LHC.

Another point of contrast is the ease with which SSSB generates a mass μ for the higgsinos. This occurs as an integral part of the SSSB and naturally accounts for the equality of the μ term and the supersymmetry breaking scales. In contrast, in gauge mediation one needs to work hard to accomplish this task [25].

Finally there are some similarities. As in gauge-mediated theories, the breaking of

supersymmetry is communicated to the squarks by the gauge interactions. Therefore the theory does not have dangerous flavor violating interactions.

5. Dynamical determination of the compactification radius

In this section we determine the value of the radius by minimizing the vacuum energy with respect to the corresponding modulus field, or equivalently with respect to the compactification scale M_c. As we mentioned already in the introduction, the Higgs contribution at the electroweak breaking minimum (4.13), being proportional to $m_H^4(M_c)$, is negligible compared to the direct contribution computed in ref. [5]. The latter reads (in the large radius limit):

$$
\begin{aligned}
E &= \frac{1}{2}\mathrm{Str} \int \frac{d^4p}{(2\pi)^4} \ln\left\{ p^2\left(1 - \gamma_T \frac{\alpha}{2\pi}\ln\frac{p^2}{\Lambda^2}\right) + M^2 \right\} + \cdots \\
&= \sum_i \eta_i \left(1 + 4\gamma_i \frac{\alpha_i}{2\pi}\ln\frac{M_c}{\Lambda}\right) M_c^4 + \cdots ,
\end{aligned}
\tag{5.1}
$$

where Λ is the cutoff scale, and all couplings in (5.1) are considered at the scale Λ. In this context the cutoff Λ is the scale at which the matching with the fundamental (string) theory is done. Scale independence of the effective action guarantees independence of the effective theory with respect to the choice of the scale Λ. For practical purposes it is customary to take $\Lambda = M_{st}$, where the boundary conditions are provided from the underlying theory. The two terms inside the bracket in the second line of (5.1) correspond, respectively, to the 1-loop and the dominant (logarithmic) two-loop contribution due to the wave-function renormalization of the i-th bulk mode, which is coupled to the massless (twisted) fields in the boundary with coupling α_i. This coupling denotes generically either the gauge or the Yukawa couplings between fields in the bulk and in the boundary. Since only even fields couple to the boundary these interactions are $N = 1$ supersymmetric. For those fields in the bulk without gauge and Yukawa interactions with the boundary (as e.g. the gravitational and moduli multiplets) $\alpha_i \equiv 0$. The dots stand for the remaining subdominant

(non logarithmic) two-loop contribution, as well as for higher loops. Note that the two-loop contribution of only bulk fields has no logarithmic dependence in M_c in analogy with finite temperature, while the two-loop contribution of only boundary fields vanishes due to supersymmetry. γ_i is a positive numerical coefficient (the eigenvalue of γ_T for the i-th bulk field) coming from the one-loop integration over the boundary states; its sign is always positive as boundary fields can never be gauge bosons in the present context.

By imposing the Λ-independence of (5.1) we can deduce the β-functions for η_i. To lowest order they are given by:

$$\beta_i \equiv \Lambda \frac{d\eta_i}{d\Lambda} = 4\eta_i\gamma_i\frac{\alpha_i}{2\pi} \, , \tag{5.2}$$

whose formal solution can be written as

$$\eta_i(M_c) = \eta_i(\Lambda) \exp\left\{ 4\gamma_i \int_0^t \frac{\alpha_i(t')}{2\pi} dt' \right\} \, , \tag{5.3}$$

where $t \equiv \ln(M_c/\Lambda)$. The logarithmic dependence of the two-loop vacuum energy and the corresponding β-functions (5.2) follow from the one-loop running of all untwisted masses $M_i^2 = \eta_i^{1/2} M_c^2$ due to the wave function renormalization of bulk fields from the massless twisted loops. As a result, the logarithms appearing in the two-loop expression (5.1) can be absorbed in the (one-loop) renormalized masses $M_i(M_c)$.

The two-loop result (5.1) can be resummed to all-loop in the leading-log approximation by the improved vacuum energy:

$$E = \eta(M_c) \, M_c^4 \, , \tag{5.4}$$

where $\eta(M_c) = \sum_i \eta_i(M_c)$ and $\eta_i(M_c)$ is defined in (5.3). Minimization of eq. (5.4) with respect to M_c leads to

$$M_c\frac{dE}{dM_c} = 4M_c^4\left[\eta(M_c) + \Gamma(M_c)\right] = 0 \, , \tag{5.5}$$

where $\Gamma = \sum_i \eta_i\gamma_i\alpha_i/2\pi$. The minimum of the potential is then given by the value of M_c such that

$$\eta(M_c) + \Gamma(M_c) = 0 \, , \tag{5.6}$$

and

$$\Gamma(M_c) + \frac{1}{4} M_c \frac{d\Gamma(M_c)}{dM_c} > 0 \,, \tag{5.7}$$

which is the condition for the extremal (5.6) to be a minimum.

This phenomenon, i.e. the appearance of a minimum by radiative corrections, has been long ago known as dimensional transmutation [26, 27], as one dimensionless parameter, η, is traded for the VEV of a field, M_c. The physical picture by which M_c does acquire a VEV is then similar to radiative breaking in field theory. We start running the η_i-parameters at the scale $\Lambda = M_{st}$ where the fundamental theory gives us the boundary values of all couplings $\alpha_i(M_{st})$. At the boundary, $\eta(M_{st})$ and $\Gamma(M_{st})$ should not satisfy eq. (5.6). As we go down with the energy the quantity $\eta + \Gamma$ should approach zero and, at a given scale M_c, it should change sign. However, taking into account from eq. (5.3) that $\eta_i(M_c)$ is a monotonically decreasing (increasing) function for $\eta_i(\Lambda) > 0$ ($\eta_i(\Lambda) < 0$), it follows that some $\eta_i(M_c)$ are required to be positive and some $\eta_i(M_c)$ should be negative. On the other hand, notice that η_i are numerical factors depending on the R-charges used to break supersymmetry. In particular, the contribution to η from a single bosonic and fermionic degree of freedom is given by:

$$\Delta\eta = -\frac{3}{128\pi^6} \left[Li_5(e^{2i\pi q_B}) - Li_5(e^{2i\pi q_F}) + h.c. \right] \,. \tag{5.8}$$

The expression (5.8) is negative for $q_B = 0$ and $q_F = 1/2$, which means that if the fermion number operator $(-1)^F$ is used for the Scherk-Schwarz breaking the condition (5.6) for radiative determination of the compactification radius is never realized. However other R-symmetries (as e.g. the $SU(2)_R$ of $N = 2$ supersymmetry) might provide different signs for different sectors and yield the necessary conditions for radiative breaking.

We can now expand eq. (5.6) to lowest order and find an approximated solution for the non-trivial minimum as

$$M_c = \exp\left\{ -\frac{1}{4} \left(\frac{\eta(\Lambda)}{\Gamma(\Lambda)} + 1 + \mathcal{O}(\hbar) \right) \right\} \Lambda \,. \tag{5.9}$$

For $\Lambda = M_c$ the solution (5.9) satisfies trivially eq. (5.6) to lowest order, as it should, while for $\Lambda = M_{st}$, M_c can be hierarchically smaller than the string scale, depending on the particular string model and on the value of the gauge couplings α_i at the string (unification) scale. In the latter case the large logarithm developed by the minimum does not invalidate perturbation theory, it just reflects a bad choice of the scale and can be reabsorbed in the renormalized parameters.

A very simple example can be provided by a model where only the strong coupling is kept and all other couplings α_i (electroweak, Yukawa, gravitational,...) are neglected. Then we have a strongly coupled gauge (gluino vector multiplet) sector with $\eta_s \equiv \eta_s(M_{st}) < 0$ and a non-interacting sector (electroweak vector multiplets, Higgs and gravitational multiplets,...) with $\eta_0 \equiv \eta_0(M_{st}) > 0$. The sign of η_s is unambiguous since for vector multiplets $q_B = 0$. The sign and value of η_0 depends of course on the field content and the R-invariance used for the Scherk-Schwarz mechanism. We will consider $\rho = \eta_0/\eta_s$ as a free parameter and adopt the running of the strong coupling given in (2.1) and the string scale and unification coupling obtained in (2.2).

In the region of scales between M_{st} and M_c the running of α_i is dominated by the linear term. Neglecting the logarithmic term in (2.1) we can obtain an analytic expression for $\eta_s(M_c)$, and so for the vacuum energy as:

$$\frac{1}{\eta_s}E = \rho - \left[\frac{M_c}{M_{st}} + \frac{b_s^{KK}\alpha_{st}}{2\pi}\left(1 - \frac{M_c}{M_{st}}\right)\right]^{4\gamma_s/\left(\frac{2\pi}{\alpha_{st}} - b_s^{KK}\right)}, \tag{5.10}$$

where $b_s^{KK} = -6$ and $\gamma_s = 6$ is the contribution to the anomalous dimension of gluons from the chiral quark supermultiplets in the boundary. In Fig. 1 we have plotted the vacuum energy (5.10) as a function of M_c, in TeV units, for $\rho = 0.74$. We see that a local minimum develops around 1 TeV.

To conclude, this mechanism can therefore be used to fix the size of the dimension that breaks supersymmetry at a TeV, in either $N = 4$ case, with logarithmic unification of gauge couplings, or in the generic $N = 2$ case with power low evolution and the string scale near
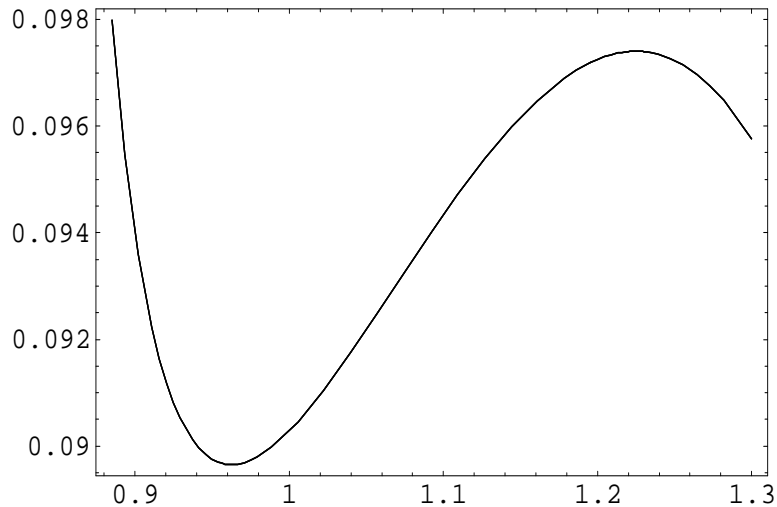
Figure 1: *Effective potential/(TeV)^4 as a function of $M_c/$ TeV in the simple model above.*

the TeV region. In this case, however, the generic bulk contribution to the vacuum energy is much bigger than M_c^4 due to the existence of n additional ultra-large dimensions of size r, that are required to account for the weakness of four-dimensional gravity:

$$E_{\text{bulk}} \sim M_c^{4+n} r^n \sim M_{st}^2 M_{P\ell}^2 \qquad \text{for} \quad M_c \sim M_{st}\,. \qquad (5.11)$$

The scaling M_c^{4+n} is a consequence of the Scherk-Schwarz breaking in $4 + n$ non-compact dimensions and can also be understood from the four-dimensional viewpoint as the multiplicity $(rM_c)^n$ of the KK-towers with respect to the n ultra-large dimensions. Since these KK-states have no standard model gauge interactions, there are no logarithmic corrections. As shown in eq. (5.11), this bulk contribution brings back essentially the problem of quadratic divergences after supersymmetry breaking and invalidates the radiative determination of the compactification scale M_c. In the context of TeV strings this problem is even worse since such a cosmological constant induces a new scale much bigger than M_{st}. One has therefore to impose the condition that this bulk contribution to the vacuum energy

vanishes. This selects out special models having equal number of bosons and fermions in the $(4 + n)$-dimensional bulk after supersymmetry breaking, level by level, at least perturbatively [28]. The next dominant contribution is then M_c^4 up to logarithms and the above mechanism of fixing M_c can be applied. Of course, the problem of determining the additional ultra-large radii r still remains open.

As a result, in both $N = 4$ and generic $N = 2$ cases, the compactification scale is determined in terms of the string scale and the unification coupling, and it can be hierarchically smaller. It is then remarkable that once the compactification scale is fixed, the phenomenology of supersymmetry breaking in both cases is very little distinct, due to the extreme softness of the Scherk-Schwarz breaking that leads to a logarithmic sensitivity of the soft terms in the string scale only at two loops. For example, from eq. (5.9) it follows that starting with a string scale $M_{st} \sim 10^{16}$ GeV, one obtains a compactification scale near the TeV region provided that $\Gamma/\eta = \mathcal{O}(10^{-2})$, which is reasonable since Γ is a two-loop correction while η is a one-loop effect.

6. Concluding remarks

The first interesting consequence of our analysis is the pattern of supersymmetry breaking. While gaugino and Higgsino masses are of the order of the compactification scale, scalar masses are generated at one loop level via gauge interactions and are naturally one order of magnitude lighter. Thus, flavor universality is guaranteed as in gauge mediated models. On the other hand, again as in gauge mediation, the stop correction to the Higgs mass-squared drives it to negative values, breaking the electroweak symmetry. The resulting spectrum consists of heavy charginos and neutralinos ($2 - 3$ TeV), squarks at $400 - 500$ GeV, and the right-handed slepton as the lightest supersymmetric particle with mass close to the electroweak scale. Moreover, the higher dimensional nature of the theory and the softness of the Scherk-Schwarz breaking replace effectively the ultraviolet cutoff with the

compactification scale, keeping the loop corrections to the Higgs mass due to the heavy gauginos small. Thus, one obtains a pattern of supersymmetry breaking with hierarchical structure, which is very different from all other scenaria. In addition, the models we study are extremely predictive, since they have no free parameters, other than a discrete option of boundary conditions, and the superparticle spectrum is fully determined.

Furthermore, this mechanism offers a possibility to determine dynamically the compactification scale by relating it to the fundamental (string) scale in a hierarchical way. One may think naively that the radius modulus dependent potential, generated by the vacuum energy, would be runaway and the extra dimension either decompactifies, or else, it shrinks to zero size in the minimum. However, in the presence of boundary (twisted) fields with gauge interactions, there are logarithmic corrections that can stabilize the radius at a nontrivial minimum. Moreover, its value has an exponential sensitivity to the coefficient of the one loop logarithm, and thus, it may be hierarchically smaller that the string scale. As a result, this mechanism can generate a very large or smaller value for the compactification scale, depending on the detailed spectrum of the model.

Acknowledgements

References

[1] S. Dimopoulos and H. Georgi, *Nucl. Phys.* **B193** (1981) 150.

[2] For a recent review see G. Giudice and R. Rattazzi, hep-ph/9801071 and references therein.

[3] I. Antoniadis, *Phys. Lett.* **B246** (1990) 377; Proc. PASCOS-91 Symposium, Boston 1991 (World Scientific, Singapore) p.718; I. Antoniadis and K. Benakli, *Phys. Lett.* **B326** (1994) 69; I. Antoniadis, K. Benakli and M. Quirós, *Phys. Lett.* **B331** (1994) 313; K. Benakli, *Phys. Lett.* **B386** (1996) 106.

[4] I. Antoniadis, C. Muñoz and M. Quirós, *Nucl. Phys.* **B397** (1993) 515.

[5] I. Antoniadis, S. Dimopoulos and G. Dvali, *Nucl. Phys.* **B516** (1998) 70.

[6] A. Pomarol and M. Quirós, hep-ph/9806263, *Phys. Lett.* **B** (1998) to appear.

[7] J. Scherk and J.H. Schwarz, *Nucl. Phys.* **B153** (1979) 61 and *Phys. Lett.* **B82** (1979) 60; E. Cremmer, J. Scherk and J.H. Schwarz, *Phys. Lett.* **B84** (1979) 83; P. Fayet, *Phys. Lett.* **B159** (1985) 121 and *Nucl. Phys.* **B263** (1986) 649.

[8] R. Rohm, *Nucl. Phys.* **B237** (1984) 553; C. Kounnas and M. Porrati, *Nucl. Phys.* **B310** (1988) 355; S. Ferrara, C. Kounnas, M. Porrati and F. Zwirner, *Nucl. Phys.* **B318** (1989) 75; C. Kounnas and B. Rostand, *Nucl. Phys.* **B341** (1990) 641; E. Kiritsis, C. Kounnas, P.M. Petropoulos and J. Rizos, *Phys. Lett.* **B385** (1996) 87.

[9] I. Antoniadis, C. Bachas, D. Lewellen and T. Tomaras, Phys. Lett. B 207 (1988) 441.

[10] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, *Phys. Lett.* **B429** (1998) 263 and hep-ph/9807344; I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. Dvali, *Phys. Lett.* **B436** (1998) 257.

[11] J. D. Lykken, *Phys. Rev.* **D54** (1996) 3693; G. Shiu and S.-H.H. Tye, hep-th/9805157.

[12] S. Ferrara, C. Kounnas and F. Zwirner, *Nucl. Phys.* **B429** (1994) 589.

[13] M. Quirós, talk given at the SUSY 98 Conference, Oxford, England, 11-17 July 1998, http://hepnts1.rl.ac.uk/SUSY98/.

[14] K.R. Dienes, E. Dudas and T. Gherghetta, *Phys. Lett.* **B436** (1998) 55 and hep-ph/9806292.

[15] D. Ghilencea and G.G. Ross, hep-ph/9809217.

[16] A. Delgado, A. Pomarol and M. Quirós, in preparation.

[17] H.B.G. Casimir, *Proc. Kon. Nederl. Akad. Wetensch.* **51** (1948) 793.

[18] E.A. Mirabelli and M. Peskin, *Phys. Rev.* **D58** (1998) 065002.

[19] See, for example, The OPAL Collaboration, *Phys. Lett.* **B433** (1998) 195.

[20] A. De Rujula, S.L. Glashow, U. Sarid, *Nucl. Phys.* **B333** (1990) 173; S. Dimopoulos, D. Eichler, R. Esmailzadeh and G.D. Starkman, *Phys. Rev.* **D41** (1990) 2388.

[21] R.S. Chivukula, A.G. Cohen, S. Dimopoulos and T.P. Walker *Phys. Rev. Lett.* **65** (1990) 957.

[22] A. Gould, B.T. Draine, R.W. Romani and S. Nussinov, *Phys. Lett.* **B238** (1990) 337.

[23] See, for example, The ALEPH Collaboration, *Eur. Phys. J.* **C4** (1998) 433.

[24] N.V. Krasnikov, *Pis'ma v ZhETF* **67** (1998) 727.

[25] G. Dvali, G.F. Giudice and A. Pomarol, *Nucl. Phys.* **B478** (1996) 31.

[26] S. Coleman and E. Weinberg, *Phys. Rev.* **D7** (1973) 1888.

[27] C. Kounnas, F. Zwirner and I. Pavel, *Phys. Lett.* **B335** (1994) 403 and references therein.

[28] S. Kachru, J. Kumar and E. Silverstein, hep-th/9807076; J.A. Harvey, hep-th/9807213; G. Shiu and S.-H.H. Tye, hep-th/9808095; S. Kachru, and E. Silverstein, hep-th/9810129.