Abstract

We present a five-dimensional model compactified on an interval where supersymmetry is broken by the Scherk-Schwarz mechanism. The gauge sector propagates in the bulk, two Higgs hypermultiplets are quasilocalized, and quark and lepton multiplets localized, in one of the boundaries. The effective four-dimensional theory is the MSSM with very heavy gauginos, heavy squarks and light sleptons and Higgsinos. The soft tree-level squared masses of the Higgs sector can be negative and they can (partially) cancel the positive one-loop contributions from the gauge sector. Electroweak symmetry breaking can then comfortably be triggered by two-loop radiative corrections from the top-stop sector. The fine tuning required to obtain the electroweak scale is found to be much smaller than in the MSSM, with essentially no fine-tuning for few TeV gaugino masses. All bounds from direct Higgs searches at LEP and from electroweak precision observables can be satisfied. The lightest supersymmetric particle is a (Higgsino-like) neutralino that can accommodate the abundance of Dark Matter consistently with recent WMAP observations.
1 INTRODUCTION

Experiments are about to probe the physics which is responsible for electroweak symmetry breaking (EWSB), and to hopefully shed some light on one of the biggest open questions which arise within the Standard Model: the origin and size of the electroweak scale. Supersymmetry (SUSY) has long been a very promising candidate to provide a satisfactory explanation of EWSB, as loops of the top-stop sector naturally drive the Higgs squared mass to negative values, enforcing the Higgs to acquire a nontrivial vacuum expectation value (VEV). Moreover these quantum corrections are cutoff at the mass of the stop itself, thereby explaining the smallness of the electroweak scale as long as the stop mass does not exceed several TeV. In fact understanding the mechanism that triggers supersymmetry breaking is one of the main issues in supersymmetric theories and it should determine the phenomenology of supersymmetric particles at future high-energy colliders as the LHC.

On the other hand the existence of extra dimensions is a general prediction of fundamental (string) theories that aim to unify all interactions, including gravity, and provide a consistent quantum description of them. In particular if the radii of the extra dimensions are as large as the 1/TeV scale \( R \), matter can propagate in the bulk and the very existence of extra dimensions can provide new mechanisms for supersymmetry and electroweak breaking. It has also been pointed out that extra dimensions can help to suppress the dangerous flavour violating interactions of SUSY breaking, as long as quark and lepton matter is localized on a supersymmetry preserving brane, as we will assume in this paper. Finally this type of models where matter fields are localized on 3-branes while the gauge sector propagates in the bulk of extra dimensions generically appears in intersecting brane constructions.

An attractive way of breaking supersymmetry (genuine to theories with extra dimensions) is the Scherk-Schwarz (SS) mechanism that makes use of twisted boundary conditions (BC’s). In five and six-dimensional supersymmetric theories the \( SU(2)_R \) invariance of the supersymmetry algebra can be used to break supersymmetry and hence its breaking should primarily be felt by \( SU(2)_R \) doublets (gauginos in vector multiplets and Higgs bosons in hypermultiplets). For definiteness we will consider a five-dimensional (5D) model where the fifth dimension is compactified on an interval of length \( \pi R \) (compactification scale \( M_c \equiv 1/R \)). States propagating in the bulk of the extra dimension break supersymmetry due to their BC’s at the endpoints of the interval. While each of the BC’s preserves an

\[^{4}\text{We will often use, unless explicitly stated, units where } R = 1.\]
$N = 1$ subgroup of the 5D $N = 2$ SUSY, they need not coincide and hence SUSY can be broken nonlocally. This particular way of breaking SUSY forbids any explicit local soft breaking terms in the 5D action. This improves the UV sensitivity, as quantum corrections are cutoff at the scale $1/R$ due to non-locality \cite{6}. Furthermore superfields localized on one of the boundaries only feel supersymmetry breaking at the loop level.

In this paper we will consider the natural scenario where gauge multiplets propagate in the bulk of the extra dimension (which is assumed to be compactified at or above the TeV scale), while quark and lepton superfields are localized towards one of the boundaries \cite{7}. In this way one can obtain a reliable superpartner spectrum where gauginos get tree level masses of the order the compactification scale, while squark and slepton masses are radiatively generated. Furthermore all (5D) massless fields are flavour blind and dangerous flavour nondiagonal interactions are mediated only by fields with (5D) masses of the order of the cutoff, $\Lambda$, and hence they are suppressed as $\exp(-\Lambda \pi R)$ \cite{3}.

Since the top-quark is localized, the stop mass is generated at one-loop and EWSB should be triggered at two-loop. A detailed discussion of this phenomenon can be found in Refs. \cite{8} where it was shown that the one-loop positive contribution to the squared Higgs mass from the gauge sector cannot be canceled by the two-loop negative contribution from the top-stop sector. It was concluded that in SM-like models with only one light Higgs and all quark and lepton fields localized in one of the boundaries EWSB does not take place (within the uncertainties of the two-loop calculation). A possible way out to this problem would be to somewhat delocalize the left-handed and/or the right-handed top quark multiplet \cite{8, 9, 10, 11}. In this case the corresponding scalar quarks feel SS supersymmetry breaking at the tree-level and so EWSB proceeds at one-loop. In these models the degree of delocalization/quasilocalization of fields $\phi$ is controled by the bulk masses $M_\phi$. Then depending on these masses FCNC \cite{12} are possible by the tree level exchange of Kaluza-Klein (KK) modes of the gauge bosons and a careful choice of masses has to be done to avoid them. Furthermore as it was pointed out in Ref. \cite{13}, in models with only one Higgs hypermultiplet propagating in the bulk, quadratically divergent Fayet-Iliopoulos (FI) terms are generated at one loop. Although consistent with both gauge symmetry and supersymmetry, these terms introduce a quadratic sensitivity of the Higgs mass to the UV cutoff. Finally even in models without quadratically divergent FI terms, in the presence of bulk masses for hypermultiplets linearly divergent FI terms $\sim M_\phi \Lambda$ may be generated unless special conditions on the mass matrix are met.
In this article we propose a different solution to the EWSB problem in models where all quark and lepton superfields are localized on one of the boundaries and the Higgs hypermultiplets have a bulk mass $M \sim M_c$. As it is well known [14] this can lead to a localization of the wavefunction of the lightest mode towards one of the branes (the lightest mode thus becomes quasilocalized) while all the higher modes typically become very heavy and decouple. Moreover one can set up a well defined expansion in powers of $\epsilon = \exp(-M \pi R)$ which can be carried out to arbitrary high orders. As the strictly localized limit (where no tree level soft terms can appear) corresponds to $O(\epsilon^0)$ one expects some tree level soft masses in the Higgs sector of the order $M\epsilon$. Due to the exponential dependence these can be naturally of the same order as bulk loop corrections if the compactification scale is taken to be in the TeV region. EWSB in this model is favored by two facts:

- For a region of the SS parameter space the tree-level soft masses can be tachyonic [14] and can then totally or partially compensate for the positive contribution from gauge one-loop radiative corrections. Under these circumstances EWSB will proceed in a fairly natural fashion triggered by the two-loop radiative corrections from the top-stop sector.

- In a model with two Higgses the condition for EWSB does not necessarily imply that one of the Higgs masses becomes negative. In some cases even if the (negative) two-loop correction is not able to overcome the positive tree-level and one-loop contributions to the soft masses, EWSB can proceed.

Our model contains two light Higgs doublets, as the MSSM, and so neither FI quadratic nor linear divergences will be generated by radiative corrections. The low energy theory is the MSSM with a peculiar spectrum of supersymmetric particles generated by the SS supersymmetry breaking. At the tree-level the corresponding supergravity theory would be a no-scale model [16] and thus no anomaly mediated supersymmetry breaking [3] will appear [19, 20]. Moreover due to the smallness of the SUSY breaking scale and the extreme softness of the SS breaking the usual fine-tunings of the MSSM can be avoided or at least, to a large extent, alleviated. Finally the lightest supersymmetric particle (LSP) of the model is a neutralino that is a good candidate to Dark Matter. To the best of our knowledge this is the first time that the MSSM with EWSB is obtained from an extra dimensional model with SS supersymmetry breaking and all matter localized in a boundary.

[5] For an application of this quasilocalization effect to matter fields and flavor physics see [11, 16, 17].
Note that SS breaking clearly distinguishes our model from those with similar bulk field content but with localized SUSY breaking proceeding at the distant brane, as for instance in Refs. [21, 22]. In particular, in the context of gaugino mediation [22], the compactification scale is generally very high (grand unification scale), while in our case it will turn out to be a few TeV. Although we are giving up MSSM-like high-scale unification, power law running of gauge couplings make unification at a much lower scale possible [23]. Furthermore, based on earlier work [24] it has been pointed out in [25] that this running should rather be interpreted as power-law threshold corrections which are exactly calculable due to the bulk $N = 2$ supersymmetry, thus opening up the fascinating possibility to construct extremely predictive models of grand unification.

The outline of the paper is as follows. In section 2 the model will be introduced and the mass eigenvalues and eigenstates, as well as the tree-level effective lagrangian, computed. A (moderate) $\mu$ problem is pointed out and a possible dynamical solution is outlined. In section 3 the conditions for EWSB are presented. We will establish that EWSB will take place radiatively: when there is a (partial) cancellation between the positive one-loop gauge corrections and the (negative) tree-level masses EWSB is triggered by the two-loop corrections induced by the top-stop sector. The degree of fine-tuning is analyzed and proven to be much less acute than in the MSSM. In section 4 numerical solutions are presented for a generic example and the typical supersymmetric spectra are depicted. All bounds, from direct searches at LEP on Higgs masses and from indirect electroweak precision observables, can be satisfied for compactification scales $M_c > \sim 6$ TeV. We have also studied the constraints from the requirement that the LSP annihilates at a rate consistent with recent WMAP data which leads to very heavy ($\gtrsim 20$ TeV) gauginos and almost Dirac light Higgsinos, still consistent with all experimental data. Finally in section 5 we present our conclusions.

2 The Higgs sector and tree level soft terms

The Higgs field is a 5D hypermultiplet which is a doublet under $SU(2)_W \otimes SU(2)_H$ where $SU(2)_H$ is a global symmetry introduced to account for two Higgs hypermultiplets. Although we will break the latter symmetry by both bulk and brane mass terms, it is useful to establish a covariant notation. We assume a flat extra dimension with coordinate $y$ parametrizing the interval $0 \leq y \leq \pi$. We will work with 4D superfields [26, 27, 28]. The hypermultiplet can be written as two left handed chiral superfields $\mathcal{H}, \mathcal{H}^c$ as

$$\mathbb{H}^{a,i} = (\mathcal{H}, \mathcal{H}^c)^{a,i}. \quad (2.1)$$
Here $i$ and $a$ are the $SU(2)_W$ and $SU(2)_H$ indices respectively. Note that the hypercharge assignment $Y = +\frac{1}{2}$ for $H$ implies $Y = +\frac{1}{2}$ for $H$ and $Y = -\frac{1}{2}$ for $H^c$. The bulk Lagrangian for the Higgs hypermultiplet reads:

$$\mathcal{L}_{\text{Higgs}} = \int d^4\theta \frac{T + \bar{T}}{2} \left\{ \mathcal{H} \exp(T_a V^a) \mathcal{H} + \mathcal{H}^c \exp(-T_a V^a) \bar{\mathcal{H}}^c \right\}$$

$$- \int d^2\theta \left\{ \mathcal{H}^c (\partial_y - M \mathcal{T}) \mathcal{H} + \text{h.c.} \right\} \quad (2.2)$$

where the mass matrix $M$ is hermitian and in general nondiagonal in $SU(2)_H$. We can parametrize it as

$$M = M' + M \vec{p} \cdot \vec{\sigma}. \quad (2.3)$$

where $M'$ and $M$ are arbitrary masses and $\vec{p}$ is a unit vector in $su(2)_H$. The radion field $T$ will be taken nondynamical,

$$T = R + 2 \omega \theta^2. \quad (2.4)$$

Its scalar component parametrizes the size of the extra dimension and a non-zero $\omega$ implements the SS breaking [27, 29] (See also Ref. [30]). The BC's for the fields $\mathcal{H}, \mathcal{H}^c$ are determined by the variational principle. The boundary Lagrangian is taken to be

$$\mathcal{L}_f^{\text{Higgs}} = \int d^2\theta \frac{1}{2} \left( \mathcal{H}^c [1 + \bar{s}_f \cdot \vec{\sigma}] \mathcal{H} + \text{h.c.} \right) \quad (2.5)$$

at the boundary at $y = y_f$ ($f = 0, \pi$), and $\bar{s}_f$ is again a unit vector in $su(2)_H$. The form of Eq. (2.5) ensures consistent BC’s by applying the variational principle to the bulk+brane system defined in Eqs. (2.2) and (2.5), taking into account the boundary terms which come from the variation of the term in Eq. (2.2) containing the derivative with respect to $y$.

The superfield BC’s following from this procedure are

$$(1 + \bar{s}_f \cdot \vec{\sigma}) \mathcal{H} = 0, \quad \mathcal{H}^c (1 - \bar{s}_f \cdot \vec{\sigma}) = 0 \quad (2.6)$$

at the corresponding boundaries. Note that the matrices acting on the fields in Eq. (2.6) take the form of projectors such that some linear combinations of $\mathcal{H}_a^a$ ($\mathcal{H}^a_0$) are set to zero at the boundary, while the BC’s of the “orthogonal fields” remain undetermined at this level. By means of a global $SU(2)_H$ rotation we can always achieve $\bar{s}_0 = (0, 0, -1)$ which

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{$^6$}Eq. (2.5) is a sufficiently general boundary Lagrangian for our purposes. A more detailed discussion of the most general boundary Lagrangian for the system Eq. (2.2) can be found in [31].
will prove it to be a convenient choice for us. With this convention, the BC’s at \( y = 0 \) read:

\[
H^2 = 0, \quad H^c_1 = 0. \tag{2.7}
\]

Note that this particular boundary condition ensures the absence of potentially hazardous quadratically divergent FI terms as the two chiral superfields \( H^1 \) and \( H^c_2 \), which do not vanish at \( y = 0 \) carry opposite hypercharge. These even fields \(^7\) can be used to write Yukawa superpotentials at this boundary for the up and down sectors respectively:

\[
W = \lambda_u H^1(x,0)Q(x)U(x) + \lambda_d H^c_2(x,0)Q(x)D(x), \tag{2.8}
\]

where the 5D Yukawa couplings \( \lambda_u \) and \( \lambda_d \) have mass dimension \(-\frac{1}{2}\).

The BC’s for the odd Higgs scalars \( H^2 \) and \( H^c_1 \) are given by the \( \theta = 0 \) component of Eq. (2.6), while those of the even Higgs scalars \( H^1 \) and \( H^c_2 \) follow from computing the bulk equations of motion (EOM) for the odd auxiliary fields and imposing both the scalar and auxiliary component of the BC’s Eq. (2.6). At \( y = 0 \) we find

\[
H^2 = 0, \quad (\partial_y - M' + c_0 M)H^1 = 0, \tag{2.9}
\]

\[
H^c_1 = 0, \quad (\partial_y + M' + c_0 M)H^c_2 = 0, \tag{2.10}
\]

where we define

\[
c_f = \vec{s}_f \cdot \vec{p}, \quad (f = 0, \pi). \tag{2.11}
\]

At \( y = \pi \), the equations take a similar form, replacing \( H^{1,2} \) and \( H^c_{1,2} \) by the corresponding linear combinations. Although our boundary conditions avoid the generation of quadratically divergent FI terms, there are linearly divergent contributions \(^8\) going as \( \sim \Lambda M' \). To further reduce UV sensitivity in our model, we will demand that \( M' = 0 \) although these terms are much less dangerous than the quadratically divergent ones.

We will assume that the \( F \)-term of the radion gets a VEV and triggers SS-SUSY breaking \(^{27,29}\). The BC’s determine the mass spectrum through three SUSY preserving parameters: the angle \( \tilde{\omega} \), defined by

\[
\cos(2\pi\tilde{\omega}) = \vec{s}_0 \cdot \vec{s}_\pi, \tag{2.12}
\]

\(^7\)In analogy to the orbifold language, we will refer to the fields which do not vanish at \( y = 0 \) as “even” (at \( y = 0 \)). Likewise the fields of Eq. (2.6) and Eq. (2.7) will be called “odd” fields at \( y = 0 \).

\(^8\)These contributions can be seen to have their origin in the sign difference between Eq. (2.9) and Eq. (2.10).
as well as the quantities $c_0$ and $c_\pi$ defined in Eq. (2.11).\textsuperscript{9} Furthermore there is only one SUSY breaking parameter, the SS twist $\omega \equiv |F_T/2|$.

The mass eigenvalues for the Higgs scalars are determined by the zeroes of the equation\textsuperscript{15}

$$\left( \cos(\Omega\pi R) - \frac{c_0 M}{\Omega} \sin(\Omega\pi R) \right) \left( \cos(\Omega\pi R) + \frac{c_\pi M}{\Omega} \sin(\Omega\pi R) \right) = \cos^2(\omega \pm \tilde{\omega}) \pi. \quad (2.13)$$

where $\Omega^2 = m^2 - M^2$. For fermions (Higgsinos) we simply have to set $\omega = 0$ in Eq. (2.13). A detailed discussion of the properties of Eq. (2.13) can be found in Refs. \textsuperscript{15, 31}. Here we will only consider an interesting limit, that of quasilocalized fields. By assuming that $Mc_0 > 0$, for

$$\epsilon \equiv \exp(-\pi c_0 MR) \ll 1 \quad (2.14)$$

there are two 4D modes $H_\pm(x)$ whose wavefunctions localize towards the boundary at $y = 0$. They have masses

$$m_\pm^2/M^2 = (1 - \epsilon_0^2) + 4\epsilon_0^2 \left( 1 - \frac{2\cos^2(\omega \pm \tilde{\omega}) \pi}{1 + c_\pi/c_0} \right) \epsilon^2 + O(\epsilon^4). \quad (2.15)$$

There might also be two modes localizing at $y = \pi$ which can be made heavy \textsuperscript{31}. From now on we will only keep the two lightest modes, which will make up the MSSM Higgs sector. The corresponding Higgsinos have Dirac masses given by Eq. (2.15) with $\omega = 0$. For the hyperscalars, the 5D fields can be approximated as\textsuperscript{10}

$$H^1(x, y) = \sqrt{c_0 M} \exp(-c_0 MRy) [H_+(x) + H_-(x)] + O(\epsilon) \quad (2.16)$$

$$H^c_2(x, y) = \sqrt{c_0 M} \exp(-c_0 MRy) [\tilde{H}_-(x) - \tilde{H}_+(x)] + O(\epsilon) \quad (2.17)$$

while the dependence of $H^2(x, y)$ and $H^c_1(x, y)$ on $H_\pm(x)$ is only of $O(\epsilon)$. From Eq. (2.8) we can see that the states

$$H_u(x) \equiv \frac{1}{\sqrt{2}} (H_+ + H_-) \quad (2.18)$$

\textsuperscript{9}Due to their geometric interpretation as angles between vectors, these three parameters are not completely independent. For fixed $\tilde{\omega}$, $c_0$ and $c_\pi$ must lie within an elliptic disc \textsuperscript{15}. For $\tilde{\omega} = 0$ (1/2) this disc degenerates to the line $c_0 = c_\pi (-c_\pi)$.

\textsuperscript{10}More precisely the corrections to Eq. (2.16) and Eq. (2.17) are $O(\epsilon^2 y^2 / \pi)$, so for $y = 0$ the suppression is actually $O(\epsilon^2)$ while at $y = \pi$ it is $O(\epsilon)$. 

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can be identified with the MSSM up and down type Higgses respectively. Note that the dimensionless 4D Yukawa couplings are given by $y_{u,d} = \sqrt{2} c_0 M \lambda_{u,d}$.

We can now easily calculate the tree-level (soft) masses in the lagrangian

$$L_{\text{mass}} = -(\mu^2 + m^2_{H_u}) |H_u|^2 - (\mu^2 + m^2_{H_d}) |H_d|^2 + m^2_3 (H_u \cdot H_d + \text{h.c.})$$

(2.20)

where $\mu$ is the Dirac mass of Higgsinos. The masses in Eq. (2.20) take the following general form

$$\mu^2 = (s_0^2 + a_1 \epsilon^2) M^2 , \quad s_0^2 \equiv 1 - c_0^2 ,$$

(2.21)

$$m^2_{H_u} = m^2_{H_d} = a_2 \epsilon^2 M^2 ,$$

(2.22)

$$m^2_3 = a_3 M^2 \epsilon^2 .$$

(2.23)

where the coefficients $a_i$ are $O(1)$ numbers which depend on the BC parameters and all higher order corrections are $O(\epsilon^4)$. Typically we would like all these masses to be of $O(m_Z)$ for EWSB to occur without too much fine-tuning. We are thus forced to choose $s_0 = O(\epsilon)$; for the geometry of the BC’s this means that the angle between the vectors $\vec{s}_0$ and $\vec{p}$ is very small and as a consequence $c_0 \simeq 1$ and $c_\pi \simeq \cos(2\pi \tilde{\omega})$.

The smallness of $s_0$ ($|s_0| \sim 1\%$ of the theoretically allowed region $|s_0| \leq 1$) gives rise to a $\mu$-problem in our model at the percent level. The fact that the $\mu$-term and the soft terms arise at different orders in the $\epsilon$ expansion can be traced back to the following fact. Notice that both boundary and bulk mass matrices preserve $U(1)_H$ subgroups of the global $SU(2)_H$, generated by $\vec{s}_f \cdot \vec{\sigma}$ and $\vec{p} \cdot \vec{\sigma}$ respectively. For $s_0 = \pm \vec{p}$ (corresponding to $s_0 = 0$) the surviving $U(1)$ at $y = 0$ and the $U(1)$ in the bulk coincide, this symmetry being broken only by the mismatched $U(1)$ at $y = \pi$. The zero modes feel this breaking through their wavefunctions, which are however suppressed at $y = \pi$ as $\sim \epsilon$. Hence we expect $\mu \sim \epsilon^2$ when $s_0 = 0$ as it can be checked from the $\epsilon$ expansion of fermionic mass eigenvalues.

When $s_0 \neq 0$, the breaking of the $U(1)$ at $y = 0$ is really felt to $O(1)$ as it occurs even

\[11\] For a typical compactification scale $M_c \equiv 1/R \sim 5 - 10$ TeV, Eqs. (2.22) and (2.23) would require $M \gtrsim M_c$, giving $\epsilon \sim 10^{-2}$. The fact that $M \sim M_c$ is a generic prediction in this class of models and might indicate that $M$ plays some role in the stabilization of $R$, as for instance in\[10\] and\[32\].
for infinitesimally small $y > 0$ and hence the $\mu$-term is unsuppressed. On the other hand
supersymmetry is broken à la Scherk-Schwarz, which can be interpreted as a mismatch of
the surviving boundary $U(1)_R$ subgroups of the $N = 2 SU(2)_R$ automorphism group in the
bulk. Again, the zero-mode wave-functions feel this only to $\mathcal{O}(\epsilon)$, and the corresponding
soft terms are suppressed as $m^2 \sim \epsilon^2$.

One could think of a dynamical solution to the $\mu$ problem in the following way. Assume
the relation $s_0 = 0$ to be exact at the 5D cutoff scale $\Lambda$. If the resulting $U(1)$ symmetry at
$y = 0$ is only broken by the VEV of a localized SM singlet field $S$ coupling as
\begin{equation}
\Delta W = \Lambda^{-1} S(x) H^1(x, 0) \mathcal{H}_2(x, 0),
\end{equation}
the $\mu$-term will be directly proportional to $\delta = \langle S \rangle / \Lambda$. The quantity $\delta$ will be small if
the $U(1)$ breaks at a lower scale, i.e. $\langle S \rangle \ll \Lambda$. In fact, one can see that the backreaction
of the new dynamically generated $\mu$-term on the boundary conditions gives $s_0 \simeq \delta$. Note
that this mechanism is the 5D version of the 4D NMSSM where a singlet is coupled to the
Higgs superfields as the term $SH_u H_d$ in the superpotential. While assuming $O(1)$ Yukawa
couplings, in 4D the VEV $\delta$ is constrained by EWSB to be $\delta \sim m_Z / \Lambda$, in our 5D theory
the much milder constraint $\delta \sim m_Z / M_c$ holds. For the purpose of this paper we will simply
assume $s_0$ to be a small quantity.

In the approximation of small $s_0$ the soft terms become
\begin{equation}
m^2_{H_u} = m^2_{H_d} = 4M^2 \sin^2(\pi \omega)(1 - \tan^2(\pi \bar{\omega})) \epsilon^2
\end{equation}
\begin{equation}
m^2_3 = 4M^2 \sin(2\pi \omega) \tan(\pi \bar{\omega}) \epsilon^2
\end{equation}
while the $\mu$-term is given by
\begin{equation}
\mu^2 = s_0^2 M^2 + \mathcal{O}(s_0^2 \epsilon^2).
\end{equation}

Finally, we would like to comment on the quartic $D$-term potential. As it has been
shown in Ref. [21], for a gauged localized chiral multiplet the tree-level quartic potential
becomes proportional to $\delta(0)$. In other words this quartic potential appears to violate SUSY
by a huge amount. It was further shown in [21] that due to a trilinear interaction of the
scalar with the adjoint chiral multiplet $\Sigma$, the $N = 2$ partner of the bulk vector multiplet,
these singularities cancel in all physical processes. In our setup the Higgs is quasi-localized

\footnote{The factor $\Lambda^{-1}$ has been introduced to render the coupling dimensionless, and we assume it to be of $\mathcal{O}(1)$.}
and all delta functions are regularized with width $M^{-1/2}$, causing the $D$-term to diverge as $\sim M$. As it will be shown in [31], one can calculate the effect of $\Sigma$ to the quartic potential exactly by solving its 5D EOM. This can be interpreted diagrammatically as integrating out the heavy KK modes of $\Sigma$, generating an effective $H^4$ interaction from the trilinear $H^2\Sigma$ coupling. Adding the latter to the explicit $D$-term we get for the neutral components of the Higgs doublets

$$V_{\text{quartic}} = \frac{g^2 + g'^2}{8} (|H_u|^2 - |H_d|^2)^2 + O(\epsilon^2). \quad (2.28)$$

where $g$ and $g'$ are the 4D $SU(2) \otimes U(1)_Y$ gauge couplings. The leading contribution is thus precisely the MSSM one, making explicit the cancellation of the divergent contributions in the strictly localized limit which corresponds to $\epsilon \to 0$.

### 3 ELECTROWEAK SYMMETRY BREAKING

In this section we will investigate in some detail the possibility of EWSB. The conditions for EWSB and stability of the flat $|H_u| = \pm |H_d|$ directions

$$(\mu^2 + m_{H_u}^2)(\mu^2 + m_{H_d}^2) < m_3^4$$

$$2\mu^2 + m_{H_u}^2 + m_{H_d}^2 > 2|m_3| \quad (3.1)$$

are incompatible with the tree-level induced SS supersymmetry breaking where $m_{H_u}^2 = m_{H_d}^2$. In this way EWSB should proceed radiatively and we must incorporate radiative corrections to the Higgs potential. As matter is strictly localized and Higgses are quasi-localized, SUSY breaking will predominantly be mediated by one-loop gaugino loops that provide a (positive) contribution to the squared masses of squarks, sleptons and Higgses.

In particular the squark masses will be dominated by the contribution from the gluinos which is given by [7]

$$\Delta m_{\tilde{t}, \tilde{b}}^2 = \frac{2 g_3^2}{3\pi^2} M_c^2 f(\omega) \quad (3.2)$$

where the function $f(\omega)$ is defined by

$$f(\omega) = \sum_{k=1}^{\infty} \frac{\sin(\pi k \omega)^2}{k^3}, \quad (3.3)$$
while electroweak gauginos provide a radiative correction to the slepton and Higgs masses as

$$\Delta^{(1)} m^2_{H_u} = \Delta^{(1)} m^2_{H_d} = \frac{3g^2 + g'^2}{8\pi^2} M_c^2 f(\omega)$$  \hspace{1cm} (3.4)$$

Furthermore there is a sizable two-loop contribution to the Higgs soft mass terms, as well as to the quartic coupling, coming from top-stop loops with the one-loop generated squark masses given by Eq. (3.2). This contribution can be estimated in the large logarithm approximation by just plugging the one-loop squark masses in the one-loop effective potential generated by the top-stop sector [7]. The goodness of this approximation has been shown in Ref. [8] where a rigorous two-loop calculation of the effective potential has been performed. For the sake of this paper, where EWSB will not be marginal (as we will see later) it is enough to consider the effective potential in the large logarithm approximation, which yields the two-loop corrections to the Higgs masses

$$\Delta^{(2)} m^2_{H_u} = \frac{3y_t^2}{8\pi^2} \Delta m^2_t \log \frac{\Delta m^2_t}{Q^2}$$,  \hspace{1cm} (3.5)$$

$$\Delta^{(2)} m^2_{H_d} = \frac{3y_b^2}{8\pi^2} \Delta m^2_b \log \frac{\Delta m^2_b}{Q^2}$$,  \hspace{1cm} (3.6)$$

where the renormalization scale should be fixed to the scale of SUSY breaking, i.e. the gaugino mass $\omega M_c$ [7]. Notice that the corrections from the bottom sector are also considered, which would only be relevant for large values of $\tan \beta$.

A word has to be said about the bulk Higgs-Higgsino one-loop contribution to the soft masses. The reason we did neglect them with respect to the one-loop gauge contribution (and even the leading two-loop one) above is that they are strongly suppressed due to their quasi-localization. The leading $O(\epsilon)$ corrections come from the Higgs-Higgsino loop contribution to the stop mass. They are proportional to the tree level soft Higgs mass $m^2_{H_u} \sim M^2 \epsilon^2$ and hence suppressed as $\epsilon^2 \log \epsilon$ with respect to the gluon-gluino contribution of Eq. (3.2). We will typically find values of $\epsilon \sim 10^{-2}$ and thus these corrections are really subleading. In principle we could easily incorporate in our analysis the radiative corrections to $m^2_3$ as calculated in [10]. However for most of the part of parameter space we are interested in, this is only a tiny correction to the tree level value, Eq. (2.27), and we will neglect it in our analysis.

Finally, the leading two-loop corrections to the quartic self coupling of $H_u$ and $H_d$ in
the potential

\[ \Delta V_{\text{quartic}} = \Delta \gamma_u |H_u|^4 + \Delta \gamma_d |H_d|^4 \]  

are given by

\[ \Delta \gamma_u = \frac{3y_t^4}{16\pi^2} \log \frac{\Delta m_t^2 + m_t^2}{m_t^2}, \]  

\[ \Delta \gamma_d = \frac{3y_b^4}{16\pi^2} \log \frac{\Delta m_b^2 + m_b^2}{m_b^2}. \]  

where \( m_t \) and \( m_b \) are the top and bottom quark masses respectively.

Electroweak symmetry breaking can now occur in our model in a very peculiar and interesting way. In fact the tree-level squared soft masses \( m_{H_u,H_d}^2 \) given in Eq. (2.25) are suppressed by the factor \( \epsilon^2 \) and therefore, for values of \( M \sim M_c \) they can be comparable in size to the one-loop gauge corrections \( \Delta^{(1)} m_{H_u,H_d}^2 \) given by Eq. (3.4). Furthermore the tree-level masses \( m_{H_u,H_d}^2 \) are negative for values of \( \tilde{\omega} > 1/4 \) and then there can be a (total or partial) cancellation between the tree-level and one-loop contributions to the Higgs masses. Under extreme conditions they can even cancel, \( m_{H_u,H_d}^2 + \Delta^{(1)} m_{H_u,H_d}^2 \simeq 0 \), in which case the negative two-loop corrections \( \Delta^{(2)} m_{H_u}^2 \) will easily trigger EWSB. On the other hand in the limit of exact localization of the Higgs fields \( \epsilon \to 0 \) the tree-level masses will vanish and the one-loop gauge and two-loop top-stop corrections have to compete, which will make the EWSB marginal, as pointed out in Ref. [8]. Similarly for \( \tilde{\omega} \leq 1/4 \) the tree level masses \( m_{H_u,H_d}^2 \) are positive definite making it again difficult, although not impossible as we argued in section 1 EWSB. These simple arguments prove that there is a wide region in the space of parameters \( (\omega, \tilde{\omega}, \epsilon) \) where EWSB easily happens without any fine-tuning of these parameters. Of course EWSB also depends on the Higgsino mass \( \mu \) and on the compactification scale \( M_c \) (or equivalently on the gluino mass as it happens in the MSSM) and we will be concerned about the possible fine-tuning in those mass parameters.

It is easy to check that, due to the smallness of the SUSY breaking scale which will be in the TeV region, as well as the extreme softness of the SS mechanism, the usual fine-tuning problems of the MSSM can almost entirely be avoided. To see this consider the \( Z \) mass from the minimization conditions of the potential in the limit \( 1 \ll \tan^2 \beta \ll m_t^2/m_b^2 \)

\[ \frac{m_Z^2}{2} = -(\mu^2 + m_{H_u}^2 + \Delta^{(1)} m_{H_u}^2 + \Delta^{(2)} m_{H_u}^2). \]  

(3.10)
As it is intuitively clear, essentially no fine tuning is necessary if we can make EWSB to work with all terms in Eq. (3.10) roughly of electroweak size. Let us quantify a little further this statement by considering the sensitivity with respect to the fundamental parameters $M_i$

$$\Delta_{M_i} = \left| \frac{M_i^2}{m_Z^2} \frac{\partial m_Z^2}{\partial M_i^2} \right|$$

(3.11)

where $M_i = \mu, m_{H_u}, M_c$. In terms of these fundamental parameters Eq. (3.10) can be rewritten as

$$m_Z^2 = -2\mu^2 - 2m_{H_u}^2 - \kappa M_c^2$$

(3.12)

where typically $\kappa \sim 10^{-3}$, and the corresponding sensitivity parameters are given by

$$\Delta_\mu = \frac{2\mu^2}{m_Z^2}$$

$$\Delta_{M_c} = |\kappa| \frac{M_c^2}{m_Z^2}$$

$$\Delta_{M_{H_u}} = |1 + \Delta_\mu + \text{sign}(\kappa) \Delta_{M_c}|$$

(3.13)

In Fig. 1 we plot the three sensitivity parameters in (3.13) for the model, that we will present in section 4, corresponding to $\omega = 0.45$, $\bar{\omega} = 0.35$ and $M = 1.65 M_c$. This model gives a viable spectrum and it is consistent with all electroweak precision observables for $M_c > \sim 6.5$ TeV. As one sees from Fig. 1 and Eq. (3.13) the largest sensitivity appears to be with respect to the parameter $m_{H_u}$. In fact for $M_c = 6.6$ TeV the required amount of fine-tuning is $\sim 10\%$ while for larger values of $M_c$ the fine-tuning naturally increases quadratically. Thus for instance for $M_c = 10$ TeV the fine-tuning is $\sim 4\%$

We can now compare this situation with the one in the MSSM. The gluino mass for a given value of $M_c$ is $M_3 = \omega M_c$ so that in our example, for $M_c \sim 10$ TeV we have $M_3 \sim 5$ TeV. In the MSSM the Z mass squared is proportional to $M_3^2$ for the same reason as in our model, but with a much larger coefficient $O(1)$ due to large logarithms $\log m_Z/m_{\text{GUT}}$. A gluino of mass a few TeV in the MSSM will require a (tan $\beta$ dependent) fine-tuning as large as 0.01%. A careful treatment of the fine tuning issues related to the gluino mass can be found in Ref. [35]. This back-of-an-envelope calculation just wanted to stress the rough differences between our mechanism of EWSB and typical results in the MSSM.

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13 We are defining our fundamental parameters such that the sensitivity on them is really a measure of fine-tuning in the sense of Ref. [34].
Figure 1: The sensitivity parameters in Eq. (3.13) as functions of $M_c$ in TeV for the case $\omega = 0.45$, $\tilde{\omega} = 0.35$ and $M = 1.65M_c$. From top to bottom the lines are: $\Delta m_{H_u}$ (blue line), $\Delta M_c$ (green line) and $\Delta \mu$ (red line).

4 Supersymmetric Spectra and Dark Matter

We will now calculate the Higgs and superpartner spectra for some specific values of the parameters. We would like to plot our predictions as functions of $M_c$ with all other parameters ($\omega, \tilde{\omega}, M$) fixed. Because of the exponential dependence of the tree level soft masses it will prove convenient to trade $M$ by $\epsilon$ (which provides a fixed ratio of $M/M_c$) when varying over $M_c$ in order to avoid excessively large or small masses.

The parameters $\omega$ and $\tilde{\omega}$ give $O(1)$ coefficients in the soft parameters. Their possible values can be further restricted by demanding that the right-handed slepton mass $m_{\tilde{\nu}_R}$ be above the mass of the lightest neutralino, as there are strong constraints on charged stable particles and we would like the lightest neutralino to be the lightest supersymmetric particle (LSP) and a Dark Matter candidate. For the nature of the latter notice that gaugino masses are given by $\omega M_c$ while Higgsino masses are essentially controlled by the $\mu$-parameter. We thus expect the neutralino to be almost pure Higgsino with a mass basically given by $\mu$. On the other hand the right handed slepton mass is radiatively generated and proportional to $g' M_c$. The size of the $\mu$ term is determined by the minimization conditions and will increase $\sim M_c$ for large $M_c$ (as it has to compensate the negative radiative corrections to $m_{H_u}^2$). However the tree level soft mass terms Eq. (2.25) increase for smaller $\tilde{\omega}$ which in turn allows for a smaller $\mu$. The requirement that the neutralino be lighter than the charged sleptons thus favours the region $\omega > \tilde{\omega}$.

We then solve the minimization conditions for EWSB which will give us two predictions, $\tan \beta$ and $\mu$ as functions of the only left free parameter, $M_c$. Then all masses will become
functions of $M_c$. In particular in the Higgs sector all masses are obtained from the effective potential where the one-loop corrections to the quartic couplings are included. The mass of the SM-like Higgs is then computed with radiative corrections to the quartic couplings considered at the one-loop level. It is well known that including just the one-loop effective potential overestimates somehow the Higgs masses and improving the effective potential by an RGE resummation of leading logarithms provides more realistic results. In this paper we will nevertheless be content by evaluating masses in the one-loop approximation. The squark and slepton masses are dominated by the gaugino loop contribution and hence grow approximately linearly with $M_c$. We find \[ (m_{\tilde{q}_L}, m_{\tilde{u}_R}, m_{\tilde{d}_R}, m_{\tilde{e}_L}, m_{\tilde{e}_R}) = (0.110, 0.103, 0.102, 0.042, 0.025) \sqrt{f(\omega)M_c} \] (4.1) where the function $f(\omega)$ is given in Eq. (3.3)\footnote{Numerically $f(\omega) \lesssim 1$ for the values of $\omega$ we will be interested in.}.

On the other hand the gauginos have a mass given by
\[ M_{1/2} = \omega M_c, \] (4.2)
and the Higgsinos, charginos and neutralinos, a mass approximately equal to $\mu$, $m_{\tilde{\chi}^\pm} \simeq m_{\tilde{\chi}^0} \simeq \mu$. They are quasi-degenerate in mass and its mass difference can be given to a very good approximation (for $\mu < 0$) by \[ \Delta m_{\tilde{\chi}} \equiv m_{\tilde{\chi}^\pm} - m_{\tilde{\chi}^0} \simeq (0.35 + 0.65 \sin 2\beta) \frac{m_W}{M_{1/2}} \] (4.3)
which means that typically e.g. for $M_c \sim 10$ TeV, $\Delta m_{\tilde{\chi}} \sim 1$ GeV. The phenomenology for Tevatron and $e^+e^-$ colliders of models where charginos and neutralinos are quasi-degenerate in mass was worked out in Refs. \[38\]. The most critical ingredients in the phenomenology of these models are the lifetime and decay modes of $\tilde{\chi}^\pm$ which in turn depend almost entirely on $\Delta m_{\tilde{\chi}}$. Conventional detection of sparticles is difficult since the decay products ($\tilde{\chi}^\pm \rightarrow \tilde{\chi}^0 \pi^\pm, \tilde{\chi}^0 \ell^\pm \nu_\ell, \ldots$) are very soft and alternative signals must be considered \[38\].

We will now consider in detail a typical example that will be solved numerically and we will plot all the predictions of the model as functions of $M_c$. We choose $\omega = 0.45$, $\bar{\omega} = 0.35$ and $M = 1.65M_c$ as in the previous example of Fig. \[1\] where the fine-tuning in these models is exemplified. The results are shown in Fig. \[2\]. The SM-like Higgs mass easily satisfies the experimental bound \[ m_{h^0} > 114.5 \text{ GeV} \] for $M_c > 6.5$ TeV. The LSP is
Figure 2: Predictions for the case $\omega = 0.45$, $\tilde{\omega} = 0.35$, $M = 1.65M_c$ (as in Fig. 1) as a function of the compactification scale. Upper left panel: $\tan \beta$. Upper right panel: the SM-like Higgs mass $m_h$. Lower left panel, from top to bottom the lines correspond to the masses of: left-handed sleptons $m_{\tilde{\ell}_L}$ (green line), heavy neutral Higgs (with a mass approximately equal to the pseudoscalar mass) $m_H \simeq m_A$ (magenta line), right-handed sleptons $m_{\tilde{e}_R}$ and neutralinos $m_{\chi^0} \simeq \mu$ (red line). Lower right panel: the squark masses $m_{\tilde{q}}$. All masses are in TeV.

the Higgsino-like with mass $\sim \mu$. Electroweak precision observables also put lower bounds on $M_c$ (see e.g. Ref. [12]). For the particularly chosen model the $\chi^2(M_c)$ distribution has a minimum around $M_c \simeq 10.5$ TeV and one deduces $M_c > 4.9$ TeV at 95% c.l.

Finally in the considered class of models where the neutralino is the LSP and $R$-parity is conserved the lightest neutralino is the candidate to Cold Dark Matter. In fact the prediction of $\Omega_{\chi^0 h^2}$ can be obtained using the DarkSUSY package [39] and can also be approximated by the expression [40]

$$\Omega_{\chi^0 h^2} \simeq 0.09 (\mu/\text{TeV})^2$$

In the particular model of Fig. 2 the prediction of $\Omega_{\chi^0 h^2}$ is given in Fig. 3.

Recent WMAP results [41] imply that $0.114 < \Omega_{\chi^0 h^2} < 0.134$. As one can see from
Fig. 3 this range in $\Omega_{\tilde{\chi}^0} h^2$ points towards the range $15 \text{ TeV} < M_c < 53 \text{ TeV}$. Then for a value of $M_c \sim 50 \text{ TeV}$ the density of Dark Matter agrees with the recent results obtained from WMAP. Notice that for such large values of $M_c$ the neutralinos are almost Dirac particles. However the non-Diracity is spoiled by $O(m_W^2/M_1) \sim 300 \text{ MeV}$ which is enough to avoid the strong limits on Dirac fermions that put a lower bound on the non-Diracity around $100 \text{ KeV}$ [42]. On the other hand the WMAP range for $M_c$ implies, in the gravitational sector, gravitino masses $m_3/2 > \sim 10 \text{ TeV}$ (depending on the value of the SS parameter $\omega$) are such that gravitinos decay early enough to avoid cosmological troubles and thus solving the longstanding cosmological gravitino problem [43].

5 Conclusions

In this paper we have analyzed electroweak symmetry breaking in a five-dimensional model where supersymmetry is broken by Scherk-Schwarz boundary conditions and quark and lepton superfields are localized at one of the boundaries. The gauge sector propagates in the bulk and thus gauginos receive tree-level masses from the Scherk-Schwarz mechanism (they are the heaviest supersymmetric particles) while squarks and sleptons acquire one-loop supersymmetry breaking masses from the bulk (they can be heavy but lighter than gauginos). The Higgs squared masses receive positive one-loop contributions from the gauge sector and negative two-loop contributions from the top-stop sector, the latter applying both to Higgses which belong to localized multiplets on the boundary and to zero modes of Higgses which belong to hypermultiplets propagating in the bulk. Under these

\footnote{Of course, such large values of $M_c$ require a fine tuning $< 1\%$, see section.}
circumstances negative two-loop corrections have to compete with positive one-loop effects and therefore electroweak breaking is marginal if not impossible.

If Higgses propagate in hypermultiplets in the bulk, but if they are quasi-localized by a supersymmetric mass, they feel the Scherk-Schwarz supersymmetry breaking as they are bulk fields but, on the other hand, their mass is not only controlled by the compactification radius $1/R$ but also by the localizing mass $M$. In fact in the “localization limit”, where $\epsilon = \exp(-\pi MR) \ll 1$, the squared masses can be comparables in size with the radiatively generated ones. Furthermore those squared tree-level masses can be physical (positive), tachyonic (negative) or even zero. The situation concerning electroweak symmetry breaking is thus very peculiar and interesting:

- If the tree-level masses are physical (or zero) electroweak breaking should be triggered at two-loop, which makes it marginal as we already pointed out. Notice that the case of localized Higgses corresponds to the limit $\epsilon \to 0$ where the tree-level soft masses are zero while a finite supersymmetric mass ($\mu$ term) may remain.

- If tree-level masses are tachyonic the conditions for electroweak symmetry breaking with stable $D$-flat directions are incompatible to each other, since the generated tree-level masses of the two Higgses are equal. However the introduction of radiative corrections, that discriminate between the $H_u$ and $H_d$ masses through the corresponding Yukawa couplings can trigger electroweak symmetry breaking.

Therefore electroweak breaking is neither purely triggered by tree-level masses nor by radiative corrections but both effects are needed: we could dub it as *tree-level assisted electroweak radiative breaking*.

The main features of these models can be summarized as follows:

- No quadratically divergent Fayet-Iliopoulos terms do appear so the Higgs mass is one-loop finite.

- Gauginos are the heaviest supersymmetric particles (they are in the TeV or multi-TeV region). Supersymmetry breaking is mediated by gauginos to the observable sector and flavor-changing neutral currents are naturally suppressed. Models are of “no-scale” type and then no anomaly mediated supersymmetry breaking occurs.

- Squarks and sleptons acquire radiative masses from loops of gluinos and electroweak gauginos, respectively. Their masses are then suppressed with respect to those of gauginos by loop factors. Furthermore, there is a striking prediction for the ratios of sfermion masses.
Eq. (4.1). Note that similar relations are known from gauge mediation models [44] (see Ref. [45] for a review). There however scalar masses are generated at the two loop level and hence different ratios apply 16.

- Due to the smallness of the supersymmetry breaking scale (in the TeV region) and the extreme softness of the Scherk-Schwarz mechanism the fine-tuning problems of the MSSM can almost entirely be avoided. For instance in our model gluinos around 3 TeV mass require a modest 10% fine-tuning.

- Higgsinos are the lightest supersymmetric particles (with a mass in the sub-TeV region). Charged and neutral Higgsinos are almost degenerate with mass splittings \( \lesssim 1 \) GeV.

- The lightest supersymmetric particle is a neutralino which is a good candidate to Dark Matter if its mass is around the TeV. This would require multi-TeV gluinos (and gravitinos) that would in turn require less than 1% fine-tuning. The gravitinos will decay early enough to avoid any cosmological problems.

The phenomenology of these models is also very peculiar. Since gauginos are superheavy they might not be detected at LHC or ILC. However since squarks are much lighter than gluinos the latter could easily decay into squarks and quarks: gluinos are then short-lived and thus they do not generate any cosmological problems. The effective theory below the TeV scale is thus a two Higgs doublet model with degenerate corresponding neutral and charged Higgsinos and (left and right-handed) sleptons. Even if Kaluza-Klein excitations might be too heavy for discovery at LHC there is a smoking gun in this model: the mass ratio between different supersymmetric particles is fixed by the relation of Eq. (4.1). The phenomenology of these models is very sensitive to the mass difference between charginos and neutralinos, which can be an indirect measure of the Kaluza-Klein masses.

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16 Depending on the size of the messenger scale and other details of the model, these relations can receive important corrections from RG running. In our case we expect RG effects to be small, as the high scale \((M_c)\) is only about two orders of magnitude above the low scale \((m_Z)\).
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