Naturally split supersymmetry

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ABSTRACT: Nonobservation of superparticles till date, new Higgs mass limits from the CMS and ATLAS experiments, WMAP constraints on relic density, various other low energy data, and the naturalness consideration, all considered simultaneously imply a paradigm shift of supersymmetric model building. In this paper we perform, for the first time, a detailed numerical study of brane-world induced supersymmetry breaking for both minimal and next-to-minimal scenarios. We observe that a naturally hierarchical spectrum emerges through an interplay of bulk, brane-localized and quasi-localized fields, which can gain more relevance in the subsequent phases of the LHC run.

Contents

1 Introduction

With no sign of supersymmetry at the CERN Large Hadron Collider (LHC) so far, even after the accumulation of ~ 5/fb data in the CMS and ATLAS experiments each, it is time to reflect on those supersymmetric models which (i) can evade easy detection at the early LHC run at 7 TeV [1], (ii) can solve problems related to large flavor changing neutral currents and CP violation [2], (iii) can give sufficient relic abundance of dark matter consistent with the WMAP data, and (iv) can still manifest in a later phase of LHC at 14 TeV with more luminosity. A minimal supersymmetric model (MSSM) spectrum like the following can do the job: light Higgsinos (around a TeV), and heavy other superpartners (few to several TeV squarks/sleptons, with a relatively light stop, and super-heavy gauginos). How natural is such a spectrum? Although a small Higgsino mixing parameter μ is encouraging from the naturalness consideration, it still requires fine-tuning to keep the quantum correction to the Higgs soft mass under control. A generic expression for this correction is given by $\Delta m^2 \sim (c/16\pi^2)m_{\tilde{q}}^2 \ln(M_S/M_Z)$, where c is an order one coefficient for third generation and small for the first two generation matter fields, M_S is the messenger scale at which supersymmetry is broken. The gluino contributes at the two loop level, so the naturalness sensitivity to gluino mass is small. Admittedly, the LHC data could not so far directly constrain the third generation squarks/sleptons, but in most of the mediation mechanisms the scalar masses of different generations are related. As LHC gradually pushes $m_{\tilde{q}}$ to higher values, naturalness would prefer a relatively low M_S (than the usual high scales preferred by gravity or even by gauge mediation). Here we take up a class of 5d scenarios introduced some years back [3] where supersymmetry breaking proceeds via Scherk-Schwarz (SS) mechanism [4–10] attributing improved naturalness. However, nonobservation of the Higgs boson to date and the WMAP relic density abundance cannot be simultaneously explained within this context, and additionally, the superparticle spectra are pushed beyond

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the reach of LHC. We incorporate a few conceptual inputs to resurrect a theoretically well-motivated framework that can address all the current issues. Here gauge fields propagate in the bulk and some (or all) matter fields are localized (with the Higgs quasi-localized) at one of the branes. Supersymmetry is broken in the bulk by SS mechanism through twisted boundary conditions, or equivalently, by the vacuum expectation value (vev) of a radion living in the bulk [11–13]. We get a naturally split spectrum where the bulk gauginos are $\mathcal{O}(10)\,\mathrm{TeV}$, while brane-localized squarks/sleptons' masses are loop suppressed. The soft masses are gennerated at the scale M_S itself, and $M_S \sim \mathcal{O}(10)\,\mathrm{TeV}$ implies a gain of a factor of ~ 7 compared to mSUGRA in the naturalness parameter [14–16]. We scan over a wide range of the model parameters to make our key observations as model independent as possible. Adding an extra gauge singlet superfield, quasi-localized near a brane, helps recover some parameter space lost earlier to collider and cosmological data, and produce a lighter spectrum with a possibility of enhanced visibility at a later phase of the LHC run.

2 Supersymmetry breaking and soft scalar masses

A 5d $N = 1$ vector supermultiplet can be decomposed from a 4d perspective into a vector multiplet $\mathcal{V}(x, y) \supset A_\mu(x, y), \lambda_1(x, y)$ and a chiral multiplet $\Phi(x, y) \supset \phi(x, y), \lambda_2(x, y)$ in the adjoint representation of the gauge group. Here, A_μ is the 4-vector gauge field, $\lambda_i (i = 1, 2)$ are gauginos, and $\phi \equiv (\Sigma + iA_5)/2$, where Σ is the 5d real scalar and A_5 is the 5th component of the 5-vector gauge field. The metric is given by $ds^2 = \eta_{\mu\nu} dx^\mu dx^\nu + R^2 dy^2$, when the 5th coordinate is compactified on S^1/Z_2 with a radius R. The gauge invariant action of bulk vector superfields coupled to a radion is given by [11–13]

$$S_{\mathrm{gauge}}^5 = \int d^4x\,dy\left[\frac{1}{4g_5^2}\int d^2\theta\,T\,W^\alpha W_\alpha + \mathrm{h.c.} + \int d^4\theta\,\frac{2}{g_5^2}\,\frac{1}{T+\bar{T}}\left(\partial_y\mathcal{V} - \frac{1}{\sqrt{2}}(\Phi+\overline{\Phi})\right)^2\right],$$
(2.1)

where $W^\alpha(x, y)$ is the field strength chiral superfield corresponding to $\mathcal{V}(x, y)$. We can write $\langle T \rangle = R + \theta^2 2\omega$, where ω is the supersymmetry breaking parameter. The mass spectrum of the component fields is given by

$$\mathcal{L}_{\mathrm{gauge}} = \frac{1}{R}\omega\lambda^{1\,(0)}\lambda^{1\,(0)} + \frac{n^2}{R^2}\left(A_\mu^{(n)} A^{\mu\,(n)} + |\Sigma^{(n)}|^2\right) + \frac{1}{R}\left(\lambda^{1\,(n)}\,\lambda^{2\,(n)}\right)\begin{pmatrix}\omega & n\\ n & \omega\end{pmatrix}\begin{pmatrix}\lambda^{1\,(n)}\\ \lambda^{2\,(n)}\end{pmatrix}.$$
(2.2)

Thus at the zero mode level we have a superfield $\mathcal{V} \supset (A_\mu, \lambda_1)$ whose gauge component remains massless while its gaugino acquires a Majorana mass ω/R, where the supersymmetry breaking parameter ω can be viewed as a twist in the $SU(2)_R$ space of which (λ_1, λ_2) is a doublet. Each Kaluza-Klein (KK) mode consists of massive gauge bosons $A_\mu^{(n)}$ and a real scalar $\Sigma^{(n)}$ each having masses of the order of n^2/R^2 (the other real component is eaten up by the KK gauge boson of the same level). Besides, there are two towers of Majorana fermions $(\lambda^{1\,(n)} \pm \lambda^{2\,(n)})$ with masses $|n \pm \omega|/R$. The masses of the brane-localized $(y = 0)$ squarks/sleptons are vanishing at tree level, and are generated at one-loop by gauge interactions [6–10],

$$m_{\widetilde{\varphi}}^2 = \frac{g^2 C_2(\widetilde{\varphi})}{4\pi^4}\left[\Delta m^2(0) - \Delta m^2(\omega)\right],$$
(2.3)

where $\Delta m^2(z) \equiv \frac{1}{2R^2}[Li_3(e^{i2\pi z}) + Li_3(e^{-i2\pi z})]$, with $Li_n(x) \equiv \sum_{k=1}^{\infty} x^k/k^n$. Here, $C_2(\widetilde{\varphi})$ is the quadratic Casimir of the $\widetilde{\varphi}$-representation under the SM gauge group. It is important to note that if the Higgs fields are localized, they receive only positive contributions from the gauge multiplets.

3 Electroweak Symmetry Breaking (EWSB)

The Higgs soft masses also receive brane-localized top-stop (bottom-sbottom) loop contributions, given by [3]

$$m^2_{H_u} = \frac{3y_t^2}{8\pi^2} m_{\tilde{t}}^2 \log \frac{m_{\tilde{t}}^2 R^2}{\omega}, \qquad m^2_{H_d} = \frac{3y_b^2}{8\pi^2} m_{\tilde{b}}^2 \log \frac{m_{\tilde{b}}^2 R^2}{\omega}. \qquad (3.1)$$

This contributions in eq. (3.1) can by itself trigger EWSB, but being a two-loop effect (since $m_{\tilde{t},\tilde{b}}$ are generated at one-loop) finds it hard to overcome a much larger one-loop positive contribution to $m^2_{H_u}$ as given by eq. (2.3). A resolution to this is to keep the H_u and H_d hypermultiplets quasi-localized near the $y = 0$ brane [3]. The advantage of quasi-localization is two-fold: (i) a bulk tachyonic mass can be generated using boundary conditions, and (ii) its mass is controlled by the supersymmetric mass M (and not $1/R$) by which quasi-localization occurs, involving a suppression factor $\epsilon = \exp(-\pi MR)$. As a result, the bulk tachyonic mass and the one-loop mass of eq. (2.3) can be of the same order, and a cancellation between them allows the two-loop contribution of eq. (3.1) dominate and trigger EWSB. The up- and down-type Higgs hypermultiplets form a doublet of a $SU(2)_H$ global symmetry of the Lagrangian. To generate a tachyonic mass one imposes suitable boundary conditions which create a twist ($\tilde{\omega}$) in that basis. The action of the bulk Higgs hypermultiplets coupled to the bulk vector and radion superfields can be written as [3],

$$S^5_{\text{Higgs}} = \int d^4x \, dy \bigg[\int d^4\theta \, \frac{T + \bar{T}}{2} \{ \bar{\mathcal{H}} \, e^{(\tau_a \mathcal{V}^a)} \, \mathcal{H} + \mathcal{H}^c \, e^{(-\tau_a \mathcal{V}^a)} \, \bar{\mathcal{H}}^c \}$$
$$- \int d^2\theta \bigg\{ \mathcal{H}^c \bigg(\partial_y - \mathcal{M}T - \frac{1}{\sqrt{2}} \Phi \bigg) \mathcal{H} + \delta(y - f) \frac{1}{2} \mathcal{H}^c [1 + \vec{s}_f \cdot \vec{\sigma}] \mathcal{H} + \text{h.c.} \bigg\} \bigg], \qquad (3.2)$$

with hypermultiplet indices suppressed. The mass matrix $\mathcal{M}$ is hermitian and non-diagonal in $SU(2)_H$ basis, given by

$$\mathcal{M} = M' + M \, p^\alpha \sigma_\alpha = a_0/R + (a/R) \, p^\alpha \sigma_\alpha, \qquad (3.3)$$

where α in the $SU(2)_H$ index, and a_0 and a are dimensionless order one coefficients. Here $\vec{s}$ and $\vec{p}$ are unit vectors in the $SU(2)_H$ space, and $(1 \pm \vec{s}_f \cdot \vec{\sigma})$ projects out a linear combination of the two $SU(2)_H$ doublet whose wave function goes to zero at the boundary. A misalignment between $\vec{s}_0$ and $\vec{s}_\pi$ causes different field combinations to survive at the two boundaries and creates a supersymmetry preserving twist angle $\tilde{\omega}$, given by $\cos(2\pi\tilde{\omega}) = \vec{s}_0 \cdot \vec{s}_\pi$.

The bulk mass term M' in eq. (3.3) was set to zero in [3] to avoid the occurence of linearly divergent ($\sim M'\Lambda$) Fayet-Iliopoulos (FI) term. Since 5d theories are inherently non-renormalizable and the cutoff in our kind of scenario is rather low, we consider putting $a_0 = 0$ is unnecessarily over-restrictive. We relax this constraint and turn on a small value
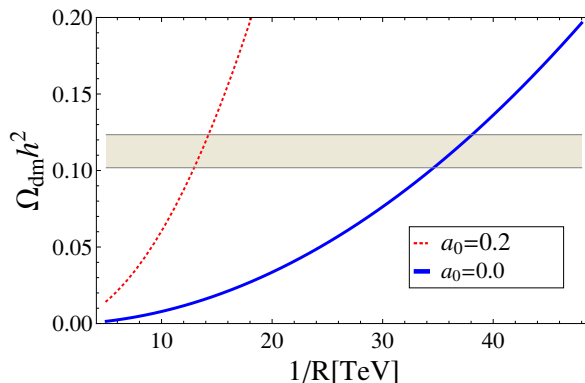
Figure 1. The dark matter density for $a = 1.65$, $\omega = 0.45$ and $\tilde{\omega} = 0.35$. The shaded region corresponds to the 3σ allowed region from WMAP [18].

of a_0 to allow the most general form of the bulk mass. We shall highlight its advantages in this paper. The soft masses of the quasi-localized up/down-type Higgses can be written as

$$m^2_{H_{u/d}} \sim M^2 \sin^2(\pi\omega)\big(1 - \tan^2(\pi\tilde{\omega})\big)\, \epsilon^2_{\mp}\,, \tag{3.4}$$

where $\epsilon_{\mp} = e^{-\pi(a\mp a_0)} \ll 1$. For $\tilde{\omega} > 1/4$ it is possible to get a tachyonic soft mass-square, while for $\epsilon \sim 10^{-2}$ the tachyonic terms can effectively cancel the positive contribution from the gaugino loops of eq. (2.3). Note that to arrange such a cancellation we simply have to put $a \sim \mathcal{O}(1)$, thus we do not pay any serious fine-tuning price.

4 The parameter space of the model

In figure 1 we demonstrate that with $a_0 \neq 0$ the relic density attains the WMAP allowed value for a relatively smaller value of R^{-1}. A nonzero a_0 increases the value of μ obtained from potential minimization. The lightest supersymmetric particle (LSP) in our model is *always* a Higgsino, and when the dark matter is Higgsino dominated it turns out that $\Omega_{\rm DM}h^2 \simeq 0.09(\mu/{\rm TeV})^2$ for $\mu \gg M_Z$ [17]. Consistency with the WMAP data [18] thus allows a lighter spectrum for $a_0 \neq 0$.

In figure 2 we display the lower limit of R^{-1} as a function of a_0 considering *all* data, especially the WMAP relic density abundance ($0.1018 < \Omega_{\rm DM}h^2 < 0.1234$) [18], the Higgs mass limits from CMS and ATLAS experiments ($115 < m_h < 127\,{\rm GeV}$) [19, 20], and lower limits on squarks/slepton masses set by Tevatron and LHC [21, 22]. For numerical estimates we have used the code microMEGAS [23, 24]. When all the three generation matter fields are brane-localized, the lower limit on R^{-1} is around $35\,{\rm TeV}$, which was $50\,{\rm TeV}$ for $a_0 = 0$. The main source of this constraint is the tension between the compatibility of EWSB occurrence and the allowed range of m_h, which tends to make the stau lighter than the Higgsino. However, if we keep Q_3 and t_R localized at $y = 0$ brane but allow all other matter fields travel in the bulk, then a stop (not a stau) becomes the next-to-lightest
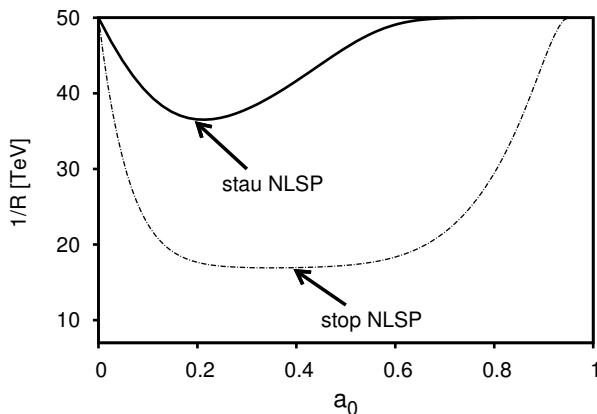
Figure 2. The lower limit of R^{-1} from all data for two different scenarios.
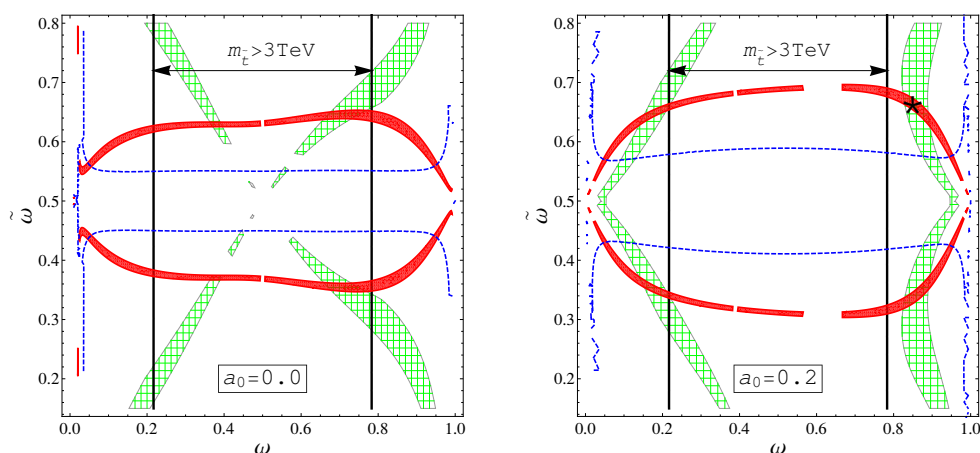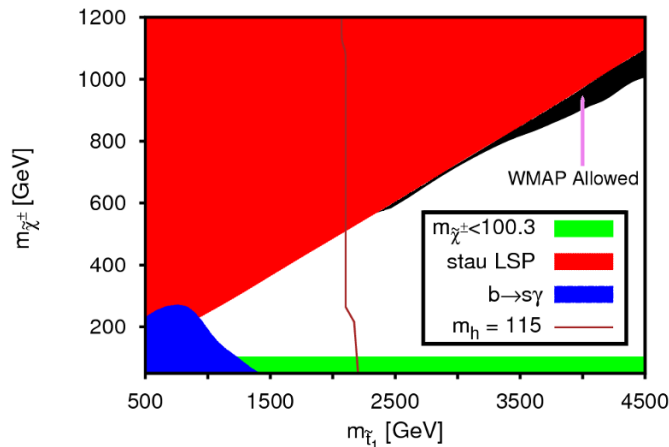

Figure 3. Allowed/disallowed zone in the twist parameters space for $1/R = 40\,\text{TeV}$ and $a = 1.65$. The green checkered region is compatible with EWSB and $115 < m_h < 127\,\text{GeV}$. The red shaded region is allowed by WMAP relic density. In between the dotted lines the stop becomes lighter than the lightest neutralino. For $a_0 = 0.2$ the region marked $(*)$ on the upper right corner maps to the parameter space where large charged tracks may be expected (see text).

supersymmetric particle (NLSP).[1] In this stop NLSP case, as we see from figure 2, the WMAP constraint gets relaxed and the lower limit on R^{-1} comes down to 16 TeV.

In figure 3 we show the constraints in the plane of the twist parameters ω and $\widetilde{\omega}$. The red shaded patches are regions where our predicted relic density is consistent with WMAP data. A nonvanishing a_0 shifts the overlap of these patches with the green chequered zone

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[1]Nonuniversal localization of fermions in the bulk, motivated for explaining the fermion mass hierarchy [25], generally leads to dangerous FCNC and CP violating operators induced by tree level flavor violating couplings of KK-gluons with the SM fermions. For different localizations of the first two families the bound on the compactification scale arising from ΔM_K and ϵ_K is quite strong ($1/R > 5000\,\text{TeV}$) [26]. Since in our case the first two families reside in bulk having identical 5d bulk masses, and only the third family quarks are brane-bound, the corresponding operators are CKM suppressed and the bound is much weaker ($1/R > 4\,\text{TeV}$) [26]. Also, the first two generation squarks for this case are too heavy to create any flavor problem at one-loop level.

Figure 4. Allowed/disallowed zone in the lightest stop- and chargino-mass plane. Only the black region is compatible with all data including WMAP.
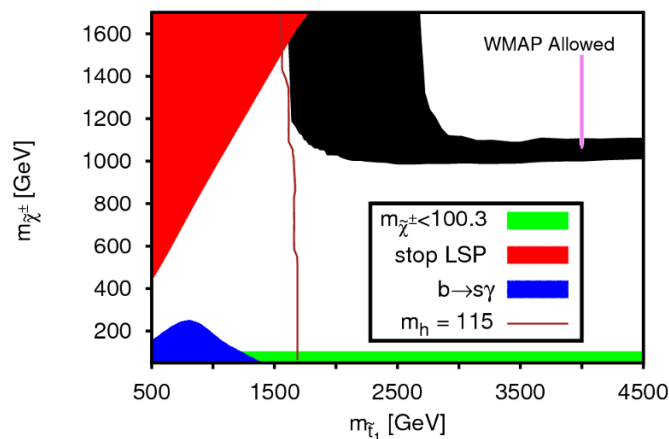
Figure 5. Same as figure 4, but only Q_3 and t_R are brane-localized, i.e. when the stop is the NLSP.

(simultaneously satisfied by EWSB and the new Higgs mass limits) to a region where the lighter stop weighs around 2 TeV.

In figures 4 and 5, we plot the constraints in the parameter space of the lighter stop mass (lightest colored sparticle) and the lighter chargino mass, when all the model parameters of the theory have been summed over in appropriate ranges. In figure 4 all matter superfields are brane-localized, whereas in figure 5 only Q_3 and t_R are brane-localized. In both cases $\tan\beta$ obtained from potential minimization varies between 3 and 15, and the trilinear coupling A_t is loop suppressed. Being almost Higgsino-like, the lighter chargino and the lightest neutralino are highly degenerate $\sim \mu$, the degeneracy being mildly lifted by radiative corrections. A substantial part of the parameter space in figure 4 is disfavored by a stau becoming an LSP. In figure 5, however, where the stop is lighter than the stau, a substantial part of the lost region is recovered. We see that a stop mass as light as 1.6 TeV is allowed in figure 5, the main constraint on it coming from the Higgs mass lower limit.

There is a substantial increase in the allowed territory (the black shaded region) which satisfies all data mentioned earlier and also the measurement of $(g-2)_\mu$ [27]. The blue shaded region in both figures 4 and 5 is excluded by $b \to s\gamma$ at 3σ [28]. To make all these plots as model independent as possible we have integrated over the model parameters over the following range: $1/R \supset [0.5 : 50]$ TeV, $\omega \supset [0 : 1]$, $\tilde\omega \supset [0 : 1]$, $a \supset [1 : 2]$ and $a_0 \supset [0 : 1]$. The lighter spectrum of figure 5 mimics that of the 'partially supersymmetric model' explored in [29]. In most of the allowed region of figure 5 the fine-tuning is about 10%.

The near equality between $m_{\tilde\chi^\pm}$ and $m_{\tilde\chi^0}$ constitutes a characteristic signature of this scenario. Within the allowed region of the model parameters, for $1/R = 40$ (16) TeV, we estimate $\Delta m_\chi \equiv m_{\tilde\chi^\pm} - m_{\tilde\chi^0}$ to lie in the range of 100 to 150 (300 to 400) MeV, which correspond to decay length 1 m to 10 cm (~ 0.5 cm) [30]. It is therefore not unexpected to observe a large charged track with heavy ionization, which corresponds to the region marked (∗) in figure 3.

5 NMSSM using a quasi-localized singlet

The next-to-minimal supersymmetric models (NMSSM) offers quite a few advantages [31]: it solves the μ problem, it can hide a Higgs boson under the cover of its singlet admixture, it has a better WMAP compatibility through a mixed singlino-Higgsino dark matter, etc. We construct a brane-world NMSSM model by quasi-localizing a gauge singlet with a supersymmetric mass M, like what we did earlier for H_u, H_d hypermultiplets. We show that the tachyonic mass of the singlet scalar indeed helps to generate its vev.

Dropping the Yukawa terms we write the superpotential and the soft breaking part of the Lagrangian as,

$$W \supset \lambda \mathcal{S} \mathcal{H}_u \cdot \mathcal{H}_d + \frac{1}{3}\kappa \mathcal{S}^3, \qquad -\mathcal{L}_{\text{soft}} \supset m_S^2 |S|^2 + \left(\lambda A_\lambda H_u \cdot H_d S + \frac{1}{3}\kappa A_\kappa S^3 + \text{h.c.}\right). \quad (5.1)$$

The vev s of the singlet scalar S is given by $\langle s \rangle \simeq \frac{1}{4\kappa}\left(-A_\kappa + \sqrt{A_\kappa^2 - 8m_S^2}\right)$, when $s \gg v_u, v_d$. A nonvanishing s therefore means either $A_\kappa > m_S^2$ or $m_S^2 < 0$. Since in our scenario A_κ is very suppressed (see later), we stimulate the $m_S^2 < 0$ option from brane-world dynamics. However, such an extreme choice of parameters, namely, $s \gg (v_u, v_d)$, is meant only for simple illustration. While doing the full potential minimization we do not make any such assumption and numerically check at every stage that EWSB does occur and low energy observables are satisfied. The bottom line is that by employing a tachyonic soft mass-square we can arrange for a nonvanishing vev s which leads to consistent phenomenology.

To follow the same method for quasi-localization we employed for H_u and H_d we must introduce an $SU(2)_H$ index to describe the bulk gauge singlet hypermultiplet $\mathbb{S}$. We write the multiplet as $\mathbb{S}^\alpha = (S_i, \Psi_s, F_{s\,i})^\alpha$, by splitting the complex hypermultiplet into two real parts, using the label α for the $SU(2)_H$ index and i is the $SU(2)_R$ index. One can introduce ω and $\tilde\omega$ exactly like before. It was shown in [32], though not specifically in the NMSSM context, that for suitable boundary conditions and for $\omega = \tilde\omega = 1/2$, a tachyonic mass $m_S^2 = -4M^2 \exp(-\pi M R)$ can be generated for a singlet scalar whose wavefunction peaks

at $y = 0$. The values of A_λ and A_κ are assumed to be zero at the scale $1/R$ and their values at the weak scale can be computed from

$$\frac{dA_\lambda}{dt} = \frac{1}{16\pi^2}\left[6A_t\lambda_t^2 + 8\lambda^2 A_\lambda + 4\kappa^2 A_\kappa + 6g_2^2 M_2 + \left(\frac{6}{5}\right)g_1^2 M_1\right] \implies A_\lambda(M_W) \sim .08\,\frac{\omega}{R}\,;$$

$$\frac{dA_\kappa}{dt} = \frac{12}{16\pi^2}(\lambda^2 A_\lambda + \kappa^2 A_\kappa) \implies A_\kappa(M_W) \sim .014\,\frac{\lambda^2\omega}{R}\,. \tag{5.2}$$

From the full scalar potential minimization we fix v_u, v_d and s and we are left with seven free parameters: R^{-1}, a, a_0, ω, $\tilde{\omega}$, λ, κ. For this NMSSM case we can afford to set $a_0 = 0$. To obtain the spectrum and the various constraints we use the package NMSSMTools [33, 34] modified for split spectrum like ours and linked to microOMEGAS [23, 24]. The key features for a benchmark point $R^{-1} = 11\,\text{TeV}$, $a = 1.6$, $\omega = 0.57$, $\tilde{\omega} = 0.66$, $\lambda = 0.4$, $\kappa = 0.06$ are the following: (i) $m_{h_1} \approx 59\,\text{GeV}$ and $m_{h_2} \approx 111\,\text{GeV}$ (this can evade the LEP-2 bound), where the lighter of the two CP-even Higgs states has a 99% branching fraction of decaying into two CP-odd states with a mass $m_{h_a} \simeq 9.4\,\text{GeV}$; ($ii$) the dark matter is the lightest neutralino with mass $\approx 56\,\text{GeV}$ with a large singlino component ($\approx 0.93\tilde{S}$); (iii) $\Omega_{\tilde{\chi}_1^0}h^2 \approx 0.1$; ($iv$) $\text{Br}(b \to s\gamma) = 3.54 \cdot 10^{-4}$; ($v$) $\text{Br}(B_s \to \mu^+\mu^-) = 2.77 \cdot 10^{-9}$. Unlike in the MSSM scenario, Δm_χ is much higher here (around $116\,\text{GeV}$ for this particular case). Its signals would be similar to as expected in the 'light Higgsino-world scenario' [1] but with enhanced cross section due to larger splitting.

6 Conclusions

By the end of 2011 supersymmetric model building has entered a new era, where the conventional gravity or gauge mediation models are feeling increasingly uncomfortable. The expected nature of superparticle spectrum is hierarchical. In this paper we have done the first detailed numerical study of a general class of brane-world inspired MSSM scenario and its NMSSM extension by confronting all laboratory and cosmological data. Some characteristic signatures are also mentioned. One of the highlights of this work is an elegant implementation of NMSSM for the first time using SS mechanism by exploiting the generation of tachyonic soft mass-square using $\text{SU}(2)_H$ rotation. In spite of its hierarchy such models suffer less from naturalness problem because of the low messenger scale at which supersymmetry is broken. This class of models is likely to gain more relevance during 2012 and beyond.

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