Supersymmetry with Light Stops

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Abstract

Recent LHC data, together with the electroweak naturalness argument, suggest that the top squarks may be significantly lighter than the other sfermions. We present supersymmetric models in which such a split spectrum is obtained through “geometries”: being “close to” electroweak symmetry breaking implies being “away from” supersymmetry breaking, and vice versa. In particular, we present models in 5D warped spacetime, in which supersymmetry breaking and Higgs fields are located on the ultraviolet and infrared branes, respectively, and the top multiplets are localized to the infrared brane. The hierarchy of the Yukawa matrices can be obtained while keeping near flavor degeneracy between the first two generation sfermions, avoiding stringent constraints from flavor and CP violation. Through the AdS/CFT correspondence, the models can be interpreted as purely 4D theories in which the top and Higgs multiplets are composites of some strongly interacting sector exhibiting nontrivial dynamics at a low energy. Because of the compositeness of the Higgs and top multiplets, Landau pole constraints for the Higgs and top couplings apply only up to the dynamical scale, allowing for a relatively heavy Higgs boson, including $m_h = 125$ GeV as suggested by the recent LHC data. We analyze electroweak symmetry breaking for a well-motivated subset of these models, and find that fine-tuning in electroweak symmetry breaking is indeed ameliorated. We also discuss a flat space realization of the scenario in which supersymmetry is broken by boundary conditions, with the top multiplets localized to a brane while other matter multiplets delocalized in the bulk.
1 Introduction

One of the strongest motivations for weak scale supersymmetry is the possibility of making electroweak symmetry breaking “natural,” i.e. a generic parameter region of the theory reproduces observed electroweak phenomena. With the Higgs potential

\[ V(h) = m^2 h^\dagger h + \lambda (h^\dagger h)^2 / 4, \]

the minimization of the potential leads to

\[ v \equiv \langle h \rangle = \sqrt{-2m^2 / \lambda} \]

and

\[ m_h^2 = -m^2, \quad (1) \]

where \( m_h \) is the physical Higgs boson mass. In the Standard Model (SM) a generic size of \( |m^2| \) is expected to be at a scale where the theory breaks down, while in supersymmetric models

\[ m^2 = |\mu|^2 + \tilde{m}_h^2, \quad (2) \]

where \( \mu \) and \( \tilde{m}_h^2 \) are the supersymmetric and supersymmetry-breaking masses for the Higgs field. Therefore, as long as these parameters are both of order the weak scale, the theory can naturally accommodate electroweak symmetry breaking.

Improved experimental constraints over the past decades, however, have cast doubt on this simple picture. In softly broken supersymmetric theories, supersymmetry-breaking masses are affected by each other through renormalization group evolution; in particular, \( \tilde{m}_h \) receives a contribution

\[ \delta \tilde{m}_h^2 \approx -\frac{3m_t^2}{4\pi^2 v^2} m_t^2 \ln \frac{M_{\text{mess}}}{m_t}, \quad (3) \]

where \( m_t \) and \( m_t \) are the top quark and squark masses, and \( M_{\text{mess}} \) the scale at which supersymmetry breaking masses are generated. (Here, we have ignored possible scalar trilinear interactions and set the left- and right-handed squark masses equal, for simplicity.) Requiring no more fine-tuning than \( \Delta \), Eqs. (1) and (2) lead to

\[ m_t \lesssim 420 \text{ GeV} \left( \frac{m_h}{125 \text{ GeV}} \right) \left( \frac{20\%}{\Delta^{-1}} \right)^{1/2} \left( \frac{3}{\ln \frac{M_{\text{mess}}}{m_t}} \right)^{1/2}. \quad (4) \]

On the other hand, recent observations at the LHC indicate:

- Generic lower bounds on the first two generation squark masses are about 1 TeV [1].
- There are hints of the SM-like Higgs boson with \( m_h \approx 125 \text{ GeV} \) [2].

Therefore, if the hints for the Higgs boson mass are true, then it strongly suggests that the squark masses have a nontrivial flavor structure, i.e. top squarks (stops) are light. \(^1\)

\(^1\)One way of avoiding this conclusion is to invoke a significant mixing of the Higgs field with another scalar field; see [3]. In general, mixing of the SM-like Higgs field with another field can weaken the naive constraint, Eq. (4), obtained in the decoupling regime (at the cost of moderate cancellation in a scalar mass-squared eigenvalue). Another possibility is to have a relatively compressed superparticle spectrum, in particular a small mass splitting between the squarks and the lightest neutralino, in which case the lower bound on the (light generation) squark masses becomes weaker.
The above observation has significant implications on an underlying model of supersymmetry breaking. This is especially because many existing models, including minimal supergravity, gauge mediation, and anomaly mediation, invoke flavor universality to avoid stringent constraints from the absence of large flavor violating processes. On the other hand, it has been realized that naturalness itself allows sfermions other than the stops (and the left-handed sbottom) to be significantly heavier than the value suggested by Eq. (4) [4, 5, 6, 7, 8, 9]. In this paper, we study a simple, general framework in which such superparticle spectra with light stops are obtained naturally.

One strategy to yield such light stop spectra is to arrange the theory in such a way that being “away” from electroweak symmetry breaking necessarily means being “close” to supersymmetry breaking, and vice versa. This makes the lighter generations (particles feeling smaller effects from electroweak symmetry breaking) obtain larger supersymmetry breaking masses, e.g. of order a few TeV, while keeping stops light. Strong constraints from flavor violation still require the first two generation sfermions to be flavor universal, but this can be achieved if these generations are both strongly localized to the supersymmetry breaking “site,” and if mediation of supersymmetry breaking there is flavor universal. The setup described here is depicted schematically in Fig. 1.

A simple way to realize the above setup is through geometry. Suppose there is an extra dimension compactified on an interval, of which the Higgs and supersymmetry breaking fields $h$ and $X$ are localized at the opposite ends. The SM gauge, quark, and lepton multiplets propagate in the bulk. Now, if two generations are localized towards the “$X$ brane” and (at least the quark doublet and up-type quark of) the other generation is localized towards the “$h$ brane,” then it
explains the (anti-)correlation between the spectrum of SM matter and its superpartners—the hierarchy of the Yukawa couplings are generated through the wavefunction overlap of SM matter with the $h$ brane, while only the first two generation sfermions obtain significant supersymmetry breaking masses through interactions with the $X$ brane.

Another manifestation of this is through dynamics—the “dimension” separating two breakings in Fig. 1 may be generated effectively as a result of strong (quasi-)conformal dynamics. Suppose there are elementary as well as composite sectors. In this case, particles in each sector interact with significant strength, but interactions involving both elementary and composite particles are suppressed by higher dimensions of composite fields. This can therefore be used to realize our setup, for example, by considering $X$ and $h$ to be elementary and composite fields, respectively. The SM matter fields are mixture of elementary and composite states—two generations being mostly elementary while the other mostly composite. In this way, the required pattern for the sfermion masses, as well as the hierarchical structure of the Yukawa couplings, are obtained. In fact, this picture can be related with the geometric picture described above. Since the strong, composite sector exhibits (approximate) conformality at high energies, the dynamics is well described by a warped extra dimension, using the AdS/CFT correspondence. (For applications of this idea in other contexts, see e.g. [10, 11, 12].)

In this paper, we present a class of models formulated in warped space, which can be interpreted either as a geometric or dynamical realization described above. In the next section, we present the basic structure of the models and interpret them as composite Higgs-top models in the desert. We pay particular attention to how strong constraints from flavor violation are avoided while generating the Yukawa hierarchy. In Section 3, we analyze electroweak symmetry breaking and present sample superparticle spectra; we also give some useful formulae for the Higgs boson mass in the appendix. In section 4, we mention a realization of our scheme in a flat space extra dimension. We conclude in Section 5.

The configuration of supersymmetry breaking and matter/Higgs fields in our models is the same as that in “emergent supersymmetry” models considered before [13, 14, 15], where the masses of elementary superpartners $\tilde{m}$ are taken (much) above the scale of strong dynamics $k'$. In this picture, the quadratic divergence of the Higgs mass-squared parameter is regulated by a combination of composite Higgsinos/stops as well as higher resonances of the strong sector (Kaluza-Klein towers). Instead, our picture here is that the theory below the compositeness scale is the full supersymmetric standard model, $\tilde{m} < k'$, so that the quadratic divergence of the Higgs mass-squared is regulated by superpartner loops as in usual supersymmetric models—the strong sector simply plays a role of generating a light stop spectrum at some energy $k'$. This alleviates the problem of a potentially large $D$-term operator [14], intrinsic to the framework of Ref. [13, 14, 15].

Three interesting papers have recently considered light stops in supersymmetry [16, 17, 18],
which are related to our study here. Ref. [16] discusses supersymmetric models in which the Higgs, top, and electroweak gauge fields are (partial) composites of a strong sector that sits at the bottom edge of the conformal window. This can be viewed as an explicit 4D realization of our warped 5D setup. (This “analogy” has also been drawn in that paper.) Ref. [17] considers the scheme of flavor mediation, where supersymmetry breaking is mediated through a gauged subgroup of SM flavor symmetries, leading to degenerate light-generation sfermions with light stops. Ref. [18] discusses light stops in the context of heterotic string theory.

2 Formulation in Warped Space

2.1 The basic structure

In this section we present a class of models realizing the basic setup of Fig. 1. We formulate it in a 5D warped spacetime with the extra dimension $y$ compactified on an $S^1/Z_2$ orbifold: $0 \leq y \leq \pi R$. The spacetime metric is given by

$$ds^2 = e^{-2ky}g_{\mu\nu}dx^\mu dx^\nu + dy^2,$$

where $k$ is the AdS curvature, which is taken to be somewhat (typically a factor of a few) smaller than the 5D cutoff scale $M_*$. The 4D Planck scale, $M_{Pl}$, is given by

$$M_{Pl}^2 \approx \frac{M_5^3}{k},$$

where $M_5$ is the 5D Planck scale, and we take $k \sim M_* \sim M_5 \sim M_{Pl}$. For now, we take the size of the extra dimension $R$ to be a free parameter, satisfying $kR > 1$. If we choose $kR \sim 10$, the TeV scale is generated by the AdS warp factor: $k' \equiv ke^{-\pi kR} \sim \text{TeV}$ [19].

We consider that the SM gauge supermultiplets $\{V_A, \Sigma_A\} (A = 1, 2, 3)$ as well as matter supermultiplets $\{\Psi_i, \Psi^c_i\}$ ($\Psi = Q, U, D, L, E$ with $i = 1, 2, 3$ the generation index) propagate in the 5D bulk. (Here, we have used the 4D $N = 1$ superfield notation; see e.g. [20].) Assuming the boundary conditions

$$\begin{pmatrix} V_A(+, +) \\ \Sigma_A(-, -) \end{pmatrix}, \quad \begin{pmatrix} \Psi_i(+, +) \\ \Psi^c_i(-, -) \end{pmatrix},$$

the low-energy field content below the Kaluza-Klein (KK) excitation scale $\sim k'$ is the gauge and matter fields of the Minimal Supersymmetric Standard Model (MSSM). They arise from the zero modes of $V_A$ and $\Psi_i$.

Now, suppose the supersymmetry breaking chiral superfield $X$ is localized on the ultraviolet (UV) brane at $y = 0$, while two Higgs-doublet chiral superfields $H_u$ and $H_d$ are on the infrared (IR) brane at $y = \pi R$. Then the bulk matter and gauge fields can interact with these fields through

$$\mathcal{L} = \delta(y) \left[ \int d^4\theta \sum_{\Psi} \left\{ \check{\eta}_{ij}X_{ij}^t \Psi_j + \left( \check{\zeta}_{ij}X_{ij}^t \Psi_j + \text{h.c.} \right) \right\} + \sum_A \left\{ \int d^2\theta \check{\xi}_A X W^a_A W_{Aa} + \text{h.c.} \right\} \right]$$

(7)
and
\[ \mathcal{L} = \delta(y - \pi R) e^{-3\pi kR} \int d^2\theta \left( \hat{y}^u_{ij} Q_u U_i H_u + \hat{y}^d_{ij} Q_i D_j H_d + \hat{y}^e_{ij} L_i E_j H_d \right) + \text{h.c.}, \] (8)
respectively, where \( \mathcal{W}_{A\alpha} \) are the field-strength superfields\(^2\).

In addition, we can introduce a singlet field \( S \) either in the bulk or on the \( y = \pi R \) brane with interactions
\[ \mathcal{L} = \delta(y - \pi R) e^{-3\pi kR} \left\{ \int d^2\theta \left( \hat{f} S H_d + \tilde{f}(S) \right) + \text{h.c.} \right\} \]
\[ + \delta(y) \int d^4\theta \left\{ \hat{\eta} S X^\dagger X S + (\hat{\zeta} S X S^\dagger S + \text{h.c.}) \right\}, \] (9)
where \( \tilde{f}(S) \) is a holomorphic function of \( S \), and the terms in the second line exist only if \( S \) is the bulk field, \( \{S, S^c\} \). The introduction of \( S \) allows us to accommodate a relatively heavy Higgs boson, including \( m_h = 125 \) GeV.

The Lagrangian for the free part of a bulk supermultiplet \( \{\Phi, \Phi^c\} \) is given by
\[ \mathcal{L} = e^{-2ky} \int d^4\theta \left( \Phi^\dagger \Phi + \Phi^c \Phi^{c\dagger} \right) + e^{-3ky} \left\{ \int d^2\theta \Phi^c \left( \partial_y + M_\Phi - \frac{3}{2} k \right) \Phi + \text{h.c.} \right\} \]
\[ + \delta(y) \int d^4\theta z_\Phi \Phi^\dagger \Phi, \] (10)
where we have included a UV-brane localized kinetic term \( z_\Phi (> 0) \), which plays an important role in our discussion. (A possible IR-brane localized kinetic term is irrelevant for the discussion.) There are two parameters in this Lagrangian: \( M_\Phi \) and \( z_\Phi \). The parameter \( M_\Phi \) controls the wavefunction profile of the zero mode in the bulk. For \( M_\Phi > k/2 \) \( (< k/2) \) the wavefunction of a zero mode arising from \( \Phi \) is localized to the UV (IR) brane; for \( M_\Phi = k/2 \) it is flat (see e.g. [21]). The parameter \( z_\Phi \) is important for a field with \( M_\Phi \approx k/2 \); it controls how much of the zero mode is regarded as the brane and bulk degrees of freedom. For \( z_\Phi M_\Phi \gg 1 \), the zero mode is mostly brane field-like, while for \( z_\Phi M_\Phi \ll 1 \) it is bulk field-like.

Our setup is realized by taking \( M_\Phi \gg k/2 \) and \( z_\Phi M_\Phi \gg 1 \) for the first two generations of matter while \( M_\Phi \ll k/2 \) for the third generation quark-doublet and right-handed top multiplets \( \{Q_3, Q_3^c\} \) and \( \{U_3, U_3^c\} \). This implies that the former are mostly brane field-like, while the latter are bulk fields with the wavefunctions localized to the IR brane. (In the 4D interpretation discussed in Section 2.3 these correspond to mostly elementary and composite fields, respectively.) The zero-mode wavefunctions for the other third generation multiplets \( \{D_3, D_3^c\} \), \( \{L_3, L_3^c\} \), and \( \{E_3, E_3^c\} \) are more flexible, although they are still subject to constraints from flavor physics, both to reproduce realistic Yukawa matrices and to avoid excessive supersymmetric contributions to flavor violation.

\(^2\)We adopt the definition of the delta function \( \int_0^\varepsilon \delta(y) = \int_{\pi R - \varepsilon}^{\pi R} \delta(y - \pi R) = 1 \), where \( 0 < \varepsilon < \pi R \).
More specifically, the wavefunction of the zero mode of the \{\Phi, \Phi^c\} multiplet in Eq. (10) is given by
\[ f_\Phi(y) = \frac{1}{\sqrt{2\pi + \frac{1}{2(M_\Phi - \frac{k}{2})^2} (1 - e^{-2\pi R(M_\Phi - \frac{k}{2})})}} e^{-(M_\Phi - \frac{k}{2})y} \] (11)
in the “conformal-field” basis, in which 5D scalar and fermion fields \(\phi\) and \(\psi\) are rescaled from the original component fields in \(\Phi\) as \(\phi = e^{-ky}\Phi|_0\) and \(\psi_\alpha = e^{-ky}\Phi|_0\). The low-energy 4D theory below \(\sim k'\) is obtained by integrating over \(y\) with this wavefunction. For the superpotential terms, it leads to
\[ \mathcal{L}_{4D} = \int d^2\theta \left( y^u_{ij} Q_i U_j H_u + y^d_{ij} Q_i D_j H_d + y^e_{ij} L_i E_j H_d + \lambda S H_u H_d + f(S) \right) + \text{h.c.}, \] (12)
where the 4D coupling constants (quantities without hat) are related with the 5D ones (with hat) by
\[ y^u_{ij} = \hat{y}^u_{ij} x_Q x_{U_j}, \quad y^d_{ij} = \hat{y}^d_{ij} x_Q x_{D_j}, \quad y^e_{ij} = \hat{y}^e_{ij} x_L x_{E_j}, \quad \lambda = \hat{\lambda} x_S, \] (13)
and \(f(S) = \hat{f}(x_S S)\). Here, the factors \(x_\Phi (\Phi = Q_i, U_i, D_i, L_i, E_i, S)\) are given by
\[ x_\Phi \equiv f_\Phi(\pi R) \simeq \begin{cases} \frac{1}{\sqrt{2\pi + \frac{1}{2(M_\Phi - \frac{k}{2})^2}}} e^{-\pi R M_\Phi} & \text{for } M_\Phi \gg \frac{k}{2} \\ \frac{1}{\sqrt{2\pi + \frac{1}{2(M_\Phi - \frac{k}{2})^2}}} & \text{for } M_\Phi \sim \frac{k}{2} \\ \frac{1}{\sqrt{2(\frac{k}{2} - M_\Phi)}} e^{-\pi R (\frac{k}{2} - M_\Phi)} & \text{for } M_\Phi \ll \frac{k}{2} \end{cases} \] (14)
where we have used \(z_\Phi k e^{-\pi k R} \ll 1\) for \(M_\Phi \ll k/2\), which is satisfied in the relevant parameter region considered later. The case with brane \(S\) is obtained by replacing \(x_S\) with 1.

For the supersymmetry-breaking terms, the 4D theory below \(\sim k'\) yields
\[ \mathcal{L}_{4D} = \int d^2\theta \left\{ \sum_{\Psi} \eta^\Psi_{ij} X^i X^j \Psi_j^\dagger + \eta^S X^i X^j S^\dagger S + \left( \sum_{\Psi} \zeta^\Psi_{ij} X^i \Psi_j^\dagger + \zeta^S X S^\dagger + \text{h.c.} \right) \right\} + \sum_A \left\{ \int d^2\theta \xi_A X W^\alpha_A W_{A\alpha} + \text{h.c.} \right\}, \] (15)
where
\[ \eta^\Psi_{ij} = \hat{\eta}^\Psi_{ij} r_{\Psi_i} r_{\Psi_j}, \quad \eta^S = \hat{\eta}^S r^2_S, \quad \zeta^\Psi_{ij} = \hat{\zeta}^\Psi_{ij} r_{\Psi_i} r_{\Psi_j}, \quad \zeta^S = \hat{\zeta}^S r^2_S, \quad \xi_A = \hat{\xi}_A, \] (16)
and the factors \(r_\Phi\) are given by
\[ r_\Phi \equiv f_\Phi(0) \simeq \begin{cases} \frac{1}{\sqrt{2\pi + \frac{1}{2M_\Phi}}} & \text{for } M_\Phi \gg \frac{k}{2} \\ \frac{1}{\sqrt{2\pi + \frac{1}{2M_\Phi}}} & \text{for } M_\Phi \sim \frac{k}{2} \\ \frac{1}{\sqrt{2(\frac{k}{2} - M_\Phi)}} e^{-\pi R (\frac{k}{2} - M_\Phi)} & \text{for } M_\Phi \ll \frac{k}{2} \end{cases} \] (17)
\(^3\)The expression \(M_\Phi \ll k/2\) here and after means that \(|M_\Phi| \ll k/2\) or \(M_\Phi < 0\).
The case with brane $S$ is obtained by $r_S \to 0$.

As will be discussed in Section 2.3, the models presented here can be interpreted, through the AdS/CFT correspondence, as those of composite Higgs-top in the supersymmetric desert. As such, small neutrino masses can be generated by the conventional seesaw mechanism. Specifically, we can introduce right-handed neutrino supermultiplets $\{N_i, N_i^c\}$ in the bulk, with Majorana masses and neutrino Yukawa couplings located on the UV and IR branes, respectively:

$$\mathcal{L} = \delta(y) \int d^2 \theta \frac{M_{ij}}{2} N_i N_j + \delta(y - \pi R) e^{-3\pi kR} \int d^2 \theta \hat{y}_{ij}^\nu L_i N_j H_u + \text{h.c.}$$

(18)

For $M_{N_i} \sim k/2$, this naturally generates small neutrino masses of the observed size (assuming the absence of tree-level neutrino-mass operators such as $\int d^2 \theta (L H_u)^2$ on the IR brane) [22]. Alternatively, small Dirac neutrino masses can be obtained if we prohibit the Majorana masses for $N_i$ and localize them to the UV brane [23].

### 2.2 Physics of flavor—fermions and sfermions

We now discuss the flavor structure of quarks/leptons and squarks/sleptons in more detail. Suppose that all the couplings on the UV brane are roughly of $O(1)$ in units of some messenger scale $M_{\text{mess}}$. In this case, Eqs. (16, 17) imply that the zero modes localized to the IR brane obtain only exponentially suppressed supersymmetry-breaking masses (at scale $k'$):

$$m_{\tilde{Q}_3, \tilde{U}_3}/m_{\tilde{\psi}_{1,2}} \ll 1.$$  

(19)

A main motivation to consider light stops is naturalness, Eq. (4). To keep this, we take $m_{\tilde{Q}_3, \tilde{U}_3} \lesssim (400 – 500)$ GeV (after evolving down to the weak scale). In order to satisfy constraints from flavor violation, the right-handed bottom and first two generation squark masses should be in the multi-TeV region [8, 24]. We therefore choose $M_{D_3} \gtrsim k/2$, and

$$m_{\tilde{\psi}_{1,2}} \sim m_{\tilde{b}_R} \sim \text{a few TeV},$$

(20)

$$m_{\tilde{t}_L, \tilde{t}_R} \sim m_{\tilde{b}_L} \lesssim (400 – 500) \text{ GeV}.$$  

(21)

The masses of $\tilde{L}_3$ and $\tilde{E}_3$ are less constrained, although we consider $M_{L_3, E_3} \gtrsim k/2$ in most of the paper, leading to $m_{\tilde{e}_L, \tilde{e}_R, \tilde{\nu}_e} \sim \text{a few TeV}$. With the mass splitting of Eqs. (20, 21), the hypercharge $D$-term contribution does not have a large effect on the Higgs mass-squared parameter to destabilize naturalness.

The masses of the gauginos are determined by parameters such as $\hat{\xi}_A$, $\hat{y}_{ij}^\psi$ and $z_\Phi$, which depend on a detailed mechanism generating operators in Eq. (7). Motivated by naturalness, in this paper we take

$$m_{B, \tilde{W}} \lesssim 1 \text{ TeV}, \quad m_{\tilde{g}} \sim 1 \text{ TeV}.$$  

(22)
The gluino mass, \( m_{\tilde{g}} \simeq M_3 \), is chosen so that the stops do not obtain large radiative corrections exceeding Eq. (21), and that the theory is not excluded by the LHC data: \( m_{\tilde{g}} \gtrsim 700 \text{ GeV} \) \[9\]. The above equations (20 – 22) specify the superpartner spectra we consider.\footnote{In deriving these expressions, we have ignored possible contributions to the supersymmetry breaking masses from the sector that stabilizes the radius of the extra dimension. This assumption is justified for certain ways of stabilizing the radius; see, e.g., Ref. \[15\].}

What about the flavor structure for quarks/leptons and those among the first two generations sfermions? In this paper, we consider a theory in which all the nontrivial flavor structures are generated from physics of the bulk (and on the IR brane). In the 4D “dual” picture discussed in Section 2.3, this corresponds to the setup in which the nontrivial flavor structure is generated through interactions of the elementary sector with the strongly-interacting composite sector. This implies that all the flavor violating effects are shut off in the high energy limit, giving the conditions

\[
\hat{\eta}_i^\Psi \propto \hat{\zeta}_i^\Psi \propto (z^\psi)_{ij} \tag{23}
\]

in flavor space, \( i,j = 1,2,3 \). In particular, in the field basis that \((z^\psi)_{ij} \propto \delta_{ij}\), which we can always take, \(\hat{\eta}_i^\Psi \propto \hat{\zeta}_i^\Psi \propto \delta_{ij}\). This can be achieved if the operators in Eq. (7) are generated by flavor universal dynamics, e.g. gauge mediation on the UV brane.

With the multi-TeV masses, the spectrum of the first two generation sfermions must be somewhat degenerate, to avoid stringent constraints from flavor. From Eq. (17), we find that the first two generation sfermion masses depend on \(\hat{\eta}_i^\Psi, (z^\psi)_{ij}, \text{ and } (M^\psi)_{ij}\). (Note that we take the bulk masses larger than \(k/2\) for the first two generations of matter.) In the field basis that \(\hat{\eta}_i^\Psi \) and \((z^\psi)_{ij}\) are proportional to the unit matrix, \(\hat{\eta}_i^\Psi \equiv \hat{\eta}_i^\Psi \delta_{ij}\) and \((z^\psi)_{ij} \equiv z^\psi \delta_{ij}\), the only source of flavor violation comes from \((M^\psi)_{ij}\), which we can diagonalize by field rotation in flavor space:

\[
(M^\psi)_{ij} = M^\psi \delta_{ij}. 
\]

The effects of flavor violation are then of order

\[
\frac{\Delta \tilde{m}_{ij}}{\tilde{m}_i + \tilde{m}_j} = \frac{r^\psi_i - r^\psi_j}{r^\psi_i + r^\psi_j}, \tag{24}
\]

multiplied by appropriate flavor mixing angles arising from diagonalization of the 4D Yukawa matrices. Here, \(r^\psi_i\) are given in Eq. (17). Requiring that these effects satisfy constraints from the \(K^-\bar{K}\) physics \[23\], we find, for example,

\[
z^\psi k \gtrsim \{15, 12, 4\} \quad \text{for } \tilde{m} = \{1, 4, 10\} \text{ TeV}, \tag{25}
\]

for \(M^\psi_2/k \approx 0.6\) and \(M^\psi_1/k \approx 0.7\), which produces hierarchy of \(O(0.1)\) by the difference of wavefunction profiles between \(\Psi_1\) and \(\Psi_2\). Here, \(\tilde{m}\) represents the masses of the first two generation sfermions, and we have assumed the maximal phase in the relevant matrix element. While the precise constraint on \(z^\psi\) depends on detailed modeling of flavor, we generically need
nonvanishing \( z_\psi \gtrsim O(10/k) \) in the case of the maximal phase in \( K-\bar{K} \) mixing.\(^3\)

The structure of the 4D Yukawa couplings can be read off from Eqs. (13, 14). For a field with \( M_\psi > k/2 \), we have a suppression arising from the wavefunction profile of the zero mode, \( \epsilon_\psi \equiv e^{-\pi R(M_\psi - k/2)} \), contributing to the hierarchy of the Yukawa couplings [21, 26]. In addition, fields with \( M_{\psi_{1,2}} > k/2 \) may have an additional suppression \( \epsilon \equiv 1/\sqrt{z_\psi M_*} \) if \( z_\psi k \gg 1 \). For example, if we take \( \lambda_{SH} \) (and the top Yukawa coupling) needs to be satisfied only below \( k/H \).

\[ \ast \]

In this section.) Since the \( \lambda_{SH} \) factors are omitted in each element, and \( \epsilon_\psi \ll 1 \) only if \( \pi R(M_\psi - k/2) \gg 1 \) and \( \epsilon \ll 1 \) only if \( e^{z_\psi k} \gg 1 \). Therefore, with suitable choices for \( M_\psi \), the observed pattern of the Yukawa couplings can be reproduced through physics of the bulk (i.e. the dynamics of the strong sector in the 4D picture) while keeping approximate flavor universality for the first two generation sfermion masses.

2.3 4D interpretation

Models discussed here can be interpreted as purely 4D models formulated in the conventional grand desert, using the AdS/CFT correspondence. (For discussions on this correspondence, see e.g. [10, 27].) In the 4D picture, the first two generations of matter are (mostly) elementary, while the third generation quark-doublet and right-handed top multiplets arise as composite fields of some strongly interacting sector, which exhibits nontrivial dynamics at an exponentially small scale \( \approx k' = ke^{-\pi kR} \). (We mostly consider that the right-handed bottom and third-generation lepton multiplets are elementary, although there is some flexibility on this choice.) This strong dynamics also produces \( S, H_u, \) and \( H_d \) fields, together with superpotential interactions \( W_H = \lambda S H_u H_d + f(S) \) at \( k' \). (We focus on the case of IR-brane localized \( S \) in this section.) Since the Higgs-top sector is strongly coupled at \( k' \), the Landau pole constraint for the couplings in \( W_H \) (and the top Yukawa coupling) needs to be satisfied only below \( k' \) [12], realizing the SUSY framework in Ref. [28].

Supersymmetry breaking is mediated at the scale \( M_{\text{mess}} \), giving TeV to multi-TeV masses to the elementary sfermions as well as the gauginos. The effect of supersymmetry breaking in

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\(^3\)In the 4D picture of Section 2.3, this corresponds to the case where the first two generations of matter are mostly elementary, with the contributions of the strong sector to their kinetic terms suppressed compared to those at tree level.
the composite sector is diluted by the near-conformal strong dynamics \cite{13}, as long as operators associated with this effect have large anomalous dimensions \cite{14}. This therefore yields only negligible soft masses for the composite fields at \( k' \) \cite{6}. A composite field, however, may obtain sizable supersymmetry breaking masses (only) if it mixes with an elementary state, which in the 5D picture corresponds to delocalizing the state from the IR brane.

The top Yukawa coupling is naturally large as the relevant fields are all composite. On the other hand, the Yukawa couplings for the first two generations of matter are generated through mixing of these states with fields in the composite sector, so are suppressed. The amount of suppression depends on the dimension of the mixing operator, and thus varies field by field, yielding a hierarchical pattern for the Yukawa matrices. Note that this way of dynamically generating the Yukawa hierarchy does not contradict the stringent constraints on supersymmetric flavor violation as long as supersymmetry breaking mediation at \( M_{\text{mess}} \) is flavor universal (e.g. as in the case of gauge mediation) and the contribution to the kinetic terms of the elementary fields from the strong sector is small (which corresponds to the condition in Eq. (25) in 5D). The overall picture for the 4D interpretation described here is depicted schematically in Fig. \( \text{[2]} \).

The value of the compositeness scale \( k' \) is constrained by phenomenological considerations. As in Eq. (21), we take the stops light to keep electroweak symmetry breaking natural. On the other hand, the LHC bound on the gluino mass for these values of stop masses is \( m_{\tilde{g}} \gtrsim 700 \text{ GeV} \), so that we need a little “hierarchy” between \( m_{\tilde{t}} \) and \( m_{\tilde{g}} \). Since \( m_{\tilde{t}} \) receives a positive contribution from \( m_{\tilde{g}} \) through renormalization group evolution, this bounds the scale \( k' \) from above. The precise bound is (exponentially) sensitive to the low energy parameters, but we typically find that \( k' \) must be below an intermediate scale; in particular, it cannot be at the unification scale. The value of \( k' \) is also limited from above by Landau pole considerations for the couplings in \( W_H \).

The lower bound on \( k' \) can be obtained for a fixed \( m_{\tilde{g}} \) by requiring that \( \tilde{t}_L, \tilde{t}_R, \) and \( \tilde{b}_L \) are sufficiently heavy to avoid the LHC bounds. Assuming that these states decay either into the lightest neutralino or the gravitino within the detector, which we would need anyway to avoid a strong constraint on quasi-stable stops, the masses of \( \tilde{t}_L \) and \( \tilde{b}_L \) must be larger than about 250 GeV \cite{9}. Moreover, if the neutralino to which these states decay is lighter than \( \approx 100 \text{ GeV} \) (or if they decay into the gravitino), then the mass of \( \tilde{b}_L \) must be larger than about 400 GeV \cite{29}. Since the running masses for these states, \( m_{\tilde{Q}_3} \) and \( m_{\tilde{U}_3} \), are vanishing at \( k' \) (up to small threshold

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\(^6\)In our models, supersymmetry breaking masses in the elementary sector, \( \tilde{m} \sim \text{a few} \text{ TeV} \), is smaller than the compositeness scale, \( k' \gtrsim 10 \text{ TeV} \) (see below). Therefore, the problem of a potentially large \( D \)-term operator \cite{13}, intrinsic to the framework of Refs. \cite{13, 14, 15}, does not arise, unless this operator is generated directly by the physics at \( M_{\text{mess}} \). The dilution of supersymmetry breaking effects in the composite sector has been studied explicitly in Ref. \cite{16} in a setup similar to ours, using Seiberg duality.
corrections), this limits $k'$ from below for a fixed $m_{\tilde{g}}$. In this paper, we take

$$k' \gtrsim 10 \text{ TeV},$$

so that the theory below the compositeness scale is the supersymmetric standard model with the superpartner spectrum given by Eqs. (20–22). With these values of $k'$, other lower bounds on $k'$ coming from precision electroweak measurements and flavor/CP violation induced by KK excitations are satisfied [30]. (Note that the masses of the lowest KK excitations are given by $\approx \pi k'$.)

Our models have the supersymmetric grand desert between $k'$ and $k \sim M_{\text{Pl}}$. Thus, if the strong sector respects a (global) unified symmetry, then we can discuss gauge coupling unification, along the lines of Ref. [31]. The prediction depends on the location of matter fields, especially $D_3, L_3$ and $E_3$; in the minimal case where these fields have $M_\Phi \gtrsim k/2$, the three SM gauge couplings approach at $\sim 10^{17}$ GeV, but with the precision of unification worse than that
in the SM ($\delta g^2/g^2 \approx 15\%$ at the unification scale). We do not pursue the issue of unification further in this paper.

3 Electroweak Symmetry Breaking

3.1 Overview

As outlined in Section 2.3, our theory above the compositeness scale $k'$ is $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ gauge theory that has the elementary fields $\Psi_{1,2}$ (and $D_3, L_3, E_3$) and the strongly interacting near-conformal sector. The beta functions for the gauge group are given by

$$b_A = b_A^{\text{MSSM}} - b_A^{Q_3+U_3+H_u,d} + b_{\text{CFT}},$$

where $b_{\text{CFT}}$ is the contribution from the strong sector, which corresponds to $1/g_5^2 k$ in the 5D picture, and is universal if this sector respects a (global) unified symmetry. Supersymmetry breaking masses for the elementary fields, including the gaugino masses $M_A$, are generated at $M_{\text{mess}}$, and they are evolved down to $k'$ by the renormalization group equations with Eq. (29). The composite fields appear at $k'$, which have vanishing supersymmetry breaking masses at that scale (up to small threshold corrections of $O(M_A^2/16\pi^2)$ in squared masses).

Physics of electroweak symmetry breaking is governed by the dynamics of the composite sector and the gaugino masses. At scale $k'$, the strong sector produces the superpotential

$$W_H = \lambda S H_u H_d + f(S) + \cdots,$$

for the Higgs sector, where the dots represent higher dimension terms which are generically suppressed by the warped-down cutoff scale $M'_s = M_s e^{-\pi k_R}$. In case $M'_s$ is close to the TeV scale, these higher dimension terms could affect phenomenology; for example, the term $(H_u H_d)^2$ can contribute to the Higgs boson mass [32]. Similarly, higher dimension terms in the Kähler potential may affect phenomenology; for example, the terms $S^\dagger H_u H_d$ and $S^\dagger H_u^* Q D$ can lead to a $\mu$ term and down-type quark masses if $S$ has an $F$-term expectation value.

In general, for relatively large values of $k'$ envisioned in Eq. (28), the effects of these higher dimension operators are insignificant, except possibly for light quark/lepton masses. We therefore consider only renormalizable terms in the Higgs potential. In particular, in the rest of the paper we focus on the case where $W_H$ contains only dimensionless terms in 4D, and discuss how electroweak symmetry breaking can work in our models. In doing so, we assume

$$m_{\tilde{Q}_3, \tilde{U}_3, H_u, H_d} \approx 0,$$

at $k'$, i.e. we ignore possible threshold corrections at that scale, which are highly model-dependent. (We later consider dynamics at the IR scale in which non-vanishing $m_S^2$ is generated at $k'$ to reproduce correct electroweak symmetry breaking.) This will illustrate basic features of electroweak symmetry breaking in our framework, in the minimal setup.
3.2 Higgs sector: $\kappa$SUSY

We consider a variant on the $\lambda$SUSY model \cite{Next-to-Minimal Supersymmetric Standard Model (NMSSM)}, which has the superpotential of the Next-to-Minimal Supersymmetric Standard Model (NMSSM) form:

$$W_H = \lambda S H_u H_d + \frac{\kappa}{3} S^3.$$ \hfill (32)

To distinguish from other $\lambda$SUSY studies in which the $\kappa$ term does not play a dominant role, we call this model $\kappa$SUSY. We assume that $S$, $H_u$, and $H_d$ are all localized to the IR brane, so we require $\lambda$ and $\kappa$ to be perturbative only up to the scale $k'$, which we take to be $10 - 1000$ TeV. For $k' = 10$ TeV, for example, we obtain $\lambda(M_Z) \lesssim 1.8$ for $\kappa(M_Z) = 0.7$; for $k' = 1000$ TeV, $\lambda(M_Z) \lesssim 1.1$ for $\kappa(M_Z) = 0.7$.

Because of Eq. (31), the only relevant dimensionful parameters for electroweak symmetry breaking are the gaugino masses, except possibly for the supersymmetry breaking mass for the $S$ field (which we will introduce in the next subsection). They set the scale for the soft supersymmetry breaking masses in the scalar potential

$$V = |\lambda H_u H_d + \kappa S^2|^2 + |\lambda S H_u|^2 + |\lambda S H_d|^2 + m_S^2 |S|^2 + m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + (\lambda A \lambda S H_u H_d - \frac{\kappa}{3} A \kappa S^3 + \text{h.c.})$$

$$+ \frac{g^2}{8} (H_u^\dagger \sigma^a H_u + H_d^\dagger \sigma^a H_d)^2 + \frac{g'^2}{8} (|H_u|^2 - |H_d|^2)^2,$$ \hfill (33)

through renormalization group evolution below $k'$. Successful electroweak symmetry breaking requires all the $S$, $H_u$, and $H_d$ fields to obtain vacuum expectation values, $v_s \equiv \langle S \rangle$, $v_u \equiv \langle H_u \rangle$, and $v_d \equiv \langle H_d \rangle$.

Once the singlet has a vacuum expectation value, $v_s$, we obtain $\mu = \lambda v_s$ and $B_\mu = \lambda A_\lambda v_s + \kappa \lambda v_s^2 = \mu (A_\lambda + \kappa \mu / \lambda)$, where $B_\mu$ is the holomorphic supersymmetry breaking Higgs mass-squared. We thus obtain the following Higgs mass-squared matrix (in the $h_u$-$h_d$-$s$ basis):

$$M^2_{\text{scalar}} \equiv \frac{1}{2} \frac{\partial^2 V}{\partial v_i \partial v_j} = \frac{1}{2} \left( \frac{\partial^2 V}{\partial v_i \partial v_j} - \delta_{ij} \frac{\partial V}{\partial v_i} \right) = \frac{1}{2} \times$$

$$\begin{pmatrix}
\bar{g}^2 v_u^2 + \frac{2 B_u}{\tan \beta} & (4 \lambda^2 - \bar{g}^2) v_u v_d - 2 B_\mu & 4 \mu v_u (\lambda - \frac{\kappa}{\tan \beta} - \frac{\lambda A_\lambda}{2 \mu \tan \beta}) \\
(4 \lambda^2 - \bar{g}^2) v_u v_d - 2 B_\mu & g^2 v_d^2 + 2 B_\mu \tan \beta & 4 \mu v_d (\lambda - \kappa \tan \beta - \frac{\lambda A_\lambda \tan \beta}{2 \mu}) \\
4 \mu v_u (\lambda - \frac{\kappa}{\tan \beta} - \frac{\lambda A_\lambda}{2 \mu \tan \beta}) & 4 \mu v_d (\lambda - \kappa \tan \beta - \frac{\lambda A_\lambda \tan \beta}{2 \mu}) & \frac{8 \kappa B_u}{\lambda} \left( 1 - \mu \frac{(A_\lambda + \kappa \mu)}{B_\mu} \right) \frac{\lambda^3 A_\lambda v_u v_d}{4 \kappa B_\mu}
\end{pmatrix},$$ \hfill (34)

where $\bar{g}^2 \equiv g^2 + g'^2$, and we have assumed that all three expectation values are real and nonzero. For us, the $A_\lambda$ and $A_\kappa$ terms are small because they are generated essentially only through weak renormalization group evolution below $k'$; $|A_{\lambda, \kappa}| \lesssim O(10 \text{ GeV})$. Other than contributing to $B_\mu$, they also contribute to singlet-doublet mixing and pseudoscalar masses, but we will ignore them.
Figure 3: Two representative plots of the scalar mass spectrum in \( \kappa \)SUSY. The solid (black), dot-dashed (red), and dashed (blue) lines represent the masses of the three mass eigenstates, which at small \( \lambda \) correspond to the SM, heavy-doublet, and singlet like Higgs bosons, respectively. The horizontal (yellow) line shows \( m_h = 125 \) GeV, and the dotted (violet) line is the mass of the lightest Higgs boson with singlet-doublet mixing turned off by hand. In the left figure we see that \( \lambda \)-doublet mixing is responsible for lowering the mass of the Higgs below its decoupling limit, Eq. \( \text{(35)} \), rather than doublet-singlet mixing. This is a generic feature for \( \tan \beta \sim 1 \). In the right figure, we see that as we increase \( \tan \beta \), singlet-doublet mixing sets in at lower \( \lambda \) than doublet-doublet mixing but that both are important in lowering the Higgs mass below Eq. \( \text{(35)} \).

in the following discussion on the (non-pseudo)scalar spectrum, as the result is not very sensitive to the values of such small \( A \) terms.

We now discuss important differences between \( \kappa \)SUSY and the MSSM as well as previous \( \lambda \)SUSY/NMSSM studies \cite{3, 33, 34}. They are illustrated in Fig. \( \text{3} \) where (tree-level) scalar masses are plotted as a function of \( \lambda \) for sample values of \( \tan \beta, \kappa, \mu \).

- We see that \( \kappa \) plays a crucial role in this theory because it appears in \( B_\mu \supset \kappa \mu^2 / \lambda \). It determines the degree of decoupling of the SM-like Higgs from the heavier scalars. The limit \( \kappa = 0 \) leads to nearly massless modes and is therefore unacceptable. In fact, as we shall see, we need \( \kappa \sim \lambda \) for a successful theory of electroweak symmetry breaking.

- The new quartic term \( \lambda^2 |H_uH_d|^2 \) leads to an extra doublet-doublet mixing which competes against \( B_\mu \): \( \mathcal{M}_{12}^2 = (2\lambda^2 - \bar{g}^2 / 2)v_u v_d - B_\mu \). As long as \( 2\lambda^2 v_u v_d < B_\mu + \bar{g}^2 v_u v_d / 2 \), this leads to the well-known enhancement of the Higgs mass in \( \lambda \)SUSY, see Eq. \( \text{(35)} \). However, once \( 2\lambda^2 v_u v_d > B_\mu + \bar{g}^2 v_u v_d / 2 \), the absolute magnitude of the off-diagonal term now increases with \( \lambda \) which leads to lowering of the Higgs mass through the very same term. We call this effect \( \lambda \)-doublet mixing. We find that in \( \kappa \)SUSY, this is the main effect that lowers the Higgs mass at large \( \lambda \) and small \( \tan \beta \), rather than mixing with the singlet, see Fig. \( \text{3} \). This
is different from Ref. [3], whose potential contains multiple extra free scales ($B_\mu$, the singlet mass) which are potentially large. (Their benchmark point has $B_\mu \approx 4 \mu^2 = (400 \text{ GeV})^2$. In this region, $\lambda$-doublet mixing accounts for only 15% of the lowering of the Higgs mass below its decoupling limit, Eq. (35); the rest comes from singlet-doublet mixing.) In fact, in $\kappa$SUSY, $\lambda \sim 2$ is excluded for $\mu \sim 200 \text{ GeV}$ exactly for this reason: the Higgs becomes tachyonic (i.e. the correct electroweak symmetry breaking vacuum disappears).

- The mass of the singlet-like scalar is not really a free parameter. It decouples together with the heavy Higgs ($B_\mu \to \infty$) but not independently. This kind of relation is to be expected in a model with a scale-free superpotential, with $\lambda \sim \kappa$. It is simply an accidental feature (due to the coefficient in $M^2_{33}$) that the singlet-like scalar is heavier than other scalars by a factor of a few, in the limit of no mixing and $\lambda \sim \kappa$.

- Doublet-singlet mixing now depends on a difference between $\lambda$ and $\kappa$. We find that, although the singlet-like Higgs is not very heavy, this greatly reduces mixing of the Higgs doublet component with the singlet and can lead to decoupling of the SM-like Higgs from the singlet for singlet masses as low as 400 GeV for $\tan \beta \sim 1$.

In the limit of small $\lambda$-doublet mixing ($2\lambda^2v_u v_d < B_\mu + \bar{g}^2 v_u v_d/2$) and negligible doublet-singlet mixing, the tree-level mass of the SM-like Higgs boson is given by

$$m_{h}^2 \approx M_Z^2 \cos^2 2\beta + \lambda^2 v^2 \sin^2 2\beta.$$  \hspace{1cm} (35)

In the appendix, we present analytical formulae for $m_h$ that include both large $\lambda$-doublet and singlet-doublet mixings up to second order in an expansion in $m_{h}^2/m_{H}^2$ (light Higgs mass over heavy Higgs mass); we also present the exact solution to $m_h$ in the regime where $\lambda$-doublet mixing dominates over doublet-singlet mixing as well as in the opposite case.

In Fig. 4(a), we present contours of the lightest Higgs boson mass, $m_h$, for typical values of parameters, where we have added the one-loop top-stop contribution with $m_{\tilde{t}} = 450 \text{ GeV}$ and $A_t = 0$. In Fig. 4(b), we show contours of the charged Higgs boson mass, which is given by

$$m_{H^+}^2 = \frac{2B_\mu}{\sin 2\beta} - \lambda^2 v^2 + M_W^2.$$  \hspace{1cm} (36)

In the non-decoupling region ($B_\mu/v^2 \gtrsim 1$) and for $\lambda > \sqrt{2}/\sin 2\beta$, the charged Higgs boson can become tachyonic. On the other hand, its mass cannot significantly fall below 300 GeV due to constraints from $b \to s\gamma$. This provides an important constraint on our parameter space and forces us to choose relatively low values of $\lambda \lesssim 1$.

Another potential issue is a light pseudoscalar arising from an approximate $R$ symmetry under which $S, H_u, H_d$ have a charge of 2/3. This symmetry is spontaneously broken by $v_u, v_d$ so that there is a light $R$-axion. This axion obtains a mass through loops of gauginos, mixing with
Figure 4: Left: Contours of the lightest Higgs boson mass in $\lambda$-tan $\beta$ plane. We find the expected rise of the Higgs mass with increasing $\lambda$ as well as the preference for low tan $\beta$. The $\lambda$-doublet mixing effect is apparent for large $\lambda$ where the Higgs mass quickly drops to zero. Right: Contours of the charged Higgs boson mass in the same plane. Relatively low values of $\lambda$ are forced by the constraint from $b \rightarrow s\gamma$, which requires $m_{H^+}$ not much below $\sim 300$ GeV.

other axions, such as the QCD axion, and $A_\lambda$, $A_\kappa$. In Ref. [35], it was determined that the $A$ terms provide the dominant contribution for $10^{-3} \lesssim |A_{\lambda,\kappa}|/v \ll 1$, which we satisfy. The mass of the $R$-axion due to the $A$ terms is given in terms of an expansion in $A_{\lambda,\kappa}/v$ by

$$m_{A_1}^2 \approx 9\frac{\mu}{\lambda} \left( \frac{\lambda A_\lambda \cos^2 \theta_A}{2} \sin 2\beta + \frac{\kappa A_\kappa}{3} \sin^2 \theta_A \right) + O(A_{\lambda,\kappa}^2),$$

(37)

where $\tan \theta_A = \mu/(\lambda v \sin 2\beta) + O(A_{\lambda,\kappa}/v)$. We see that the mass is the geometric mean of $\mu$ and $A_{\lambda,\kappa}$ times $O(1)$ factors. Since we generically have $|A_\lambda| > 1$ GeV, the mass is in tens of GeVs, so we are safe from the constraint from $\Upsilon$ decays. Since $\lambda, \kappa$ are $O(1)$, however, the Higgs can also decay into the $R$-axion with a large branching fraction, if this decay mode is kinematically allowed. Assuming $m_\phi = 125$ GeV, we find that this happens for $|A_\lambda| < 10$ GeV (neglecting $A_\kappa$). Depending on parameters, we can have $|A_\lambda| \gtrsim 10$ GeV, in which case decays of the lightest Higgs boson are SM-like.
3.3 Sample spectra

We here present sample parameter points in $\kappa$SUSY. To achieve successful electroweak symmetry breaking, in particular to obtain a sufficiently large $\mu = \lambda v_s$, we introduce a negative soft mass-squared for the singlet at $k'$, $m_S^2 \sim -(400 \text{ GeV})^2$. Such a term can arise naturally if there are (additional) messenger fields $f, \bar{f}$ on the IR brane which couple to the $S$ field in the superpotential $W = S f \bar{f}$ [30]. Here $f, \bar{f}$ are assumed to be SM-gauge singlets and have supersymmetric and supersymmetry breaking masses (roughly) of order $k'$: $M_f \sim \sqrt{F_f} \sim k'$. (This does not require a strong coincidence because the characteristic scale on the IR brane is $k' \sim M'_{\ast}$.)

The $A$ terms generated by $f, \bar{f}$ loops are small for $M_f \sim \sqrt{F_f}$, since both the $A$ terms and the soft mass-squared, $m_S^2$, are generated at the one-loop order.

We present two sample spectra in Figs. 5 and 6, which correspond respectively to two different choices of the compositeness scale, $k' = 10 \text{ TeV}$ and $1000 \text{ TeV}$, and will be discussed in more detail in Sections 3.3.1 and 3.3.2. The relevant parameters for electroweak symmetry breaking are $\lambda, \kappa, m_S^2$, and the electroweak gaugino masses $M_{1,2}$. (We choose $m_{H_u, H_d, \tilde{Q}_3, \tilde{U}_3}^2 \approx 0$ at $k'$, ignoring loop-suppressed threshold corrections.) The gluino mass is chosen to be small (but still allowed by the experimental constraint) to alleviate fine-tuning, and the bino is chosen to be the lightest observable-sector supersymmetric particle (LOSP). For the gluino mass we add the one-loop threshold correction, which can be as large as $\approx 20\%$ for the multi-TeV squark masses [37]. In this section, we assume that the gravitino is heavier than the LOSP, so that the bino is the lightest supersymmetric particle. This is the case for $M_{\text{mess}} \gtrsim M_{\text{Pl}}$, or for $M_{\text{mess}} \lesssim M_{\text{Pl}}$ if there is additional supersymmetry breaking that does not contribute to the MSSM superparticle masses but pushes up the gravitino mass above the LOSP mass [38]. If the gravitino is lighter than the bino, somewhat stronger bounds on the gluino mass would apply [39]. For example, if the bino decays promptly to the gravitino, then the lower bound is $m_\tilde{g} \approx 900 \text{ GeV}$, instead of $\approx 700 \text{ GeV}$.

In presenting the sample points, we also evaluate the amount of fine-tuning, adopting a conventional criterion [40]

$$\Delta = \max_{i,j} \frac{d \ln A_i}{d \ln B_j},$$

where $A_i = (m_{h_i}^2, v^2)$ and $B_j$ are UV parameters to be specified below. The $A_i$ correspond to the $(\theta_{h_u, \hat{v}_u} + \theta_{h_d, \hat{v}_d} + \theta_{h_s, \hat{v}_s})$ and $(\hat{v}_u + \hat{v}_d)/v$ directions in the three-dimensional $v_u, v_d, v_s$ space, respectively, where we define scalar mixing angles in terms of eigenvector overlap: $h = \theta_{h_u, h_d} b_u + \theta_{h_d, h_s} b_d + \theta_{h_s, s} s$. In the case of $\lambda$-doublet or singlet-doublet mixing, fine-tuning (e.g. due to stops) may be much alleviated compared to the MSSM due to level repulsion which is generated naturally through large off-diagonal elements in the mass matrix; in the case of singlet-doublet mixing, this has been analyzed in Ref. [3]. We here point out that large-mixing, natural scenarios with TeV-scale stops are typically accompanied by drastic deviations of Higgs couplings, so if the Higgs has only moderate deviations from SM cross sections and decay rates, then naturalness
generically requires light stops. In our analysis, we choose \( B_j = (\lambda, \kappa, m_\Sigma^2, M_{1,2,3}, k', y_t, \tilde{m}) \).

### 3.3.1 \( k' = 10 \text{ TeV} \)

The following considerations give a bottom-up picture of what is needed to generate a natural superpartner spectrum (in the decoupling regime) \cite{41,9} that radiatively breaks electroweak symmetry with \( k' = 10 \text{ TeV} \):

- The fine-tuning constraint \( (\Delta^{-1} \gtrsim 20\%) \) requires \(|\mu| \lesssim 210 \text{ GeV}, m_{\tilde{t}} \lesssim 410 \text{ GeV}\) (for degenerate stop masses without mixing), \( m_{\tilde{g}} \lesssim 790 \text{ GeV}\) (at the leading-log level; the actual bound is significantly weaker because of the effect of strong interactions), \( m_{\tilde{W}} \lesssim 890 \text{ GeV}, m_{\tilde{B}} \lesssim 2800 \text{ GeV}, \) and \( \tilde{m} \lesssim 4 \text{ TeV} \).

- Electroweak symmetry breaking requires \( \lambda, \kappa \sim 0.7 \) at low energies, as discussed in the last section; we also need \( m_\Sigma^2 \sim -(400 \text{ GeV})^2 \) to generate a sufficiently large \( \mu \) term.

In Fig. 5, we show a typical mass spectrum for \( k' = 10 \text{ TeV} \), where the lightest Higgs boson mass is evaluated with the one-loop top-stop contribution added. The production cross section \( \sigma(gg \to h) \) is modified relative to the SM due to non-decoupling stop contributions and \( A \) terms; this sample point has an enhancement of 13%. Unlike in the MSSM, the decay rate of the Higgs into \( bb \) is depleted in \( \lambda \text{SUSY} \) relative to the SM rate. As expanded in \( m_{hh}^2/m_{h}^2 \), the rate is given by (see the appendix)

\[
\frac{\Gamma(h \to bb)}{\Gamma_{SM}(h \to bb)} = 1 - \tan^2 \beta \frac{\lambda v \sqrt{\lambda^2 v^2 - M_Z^2}}{2B_\mu}.
\]  

(39)

For the \( k' = 10 \text{ TeV} \) spectrum, this formula gives 0.88, within 10% of the exact result, 0.96. Because of this suppression, the branching ratios into other modes are enhanced. In particular, we find that \( \text{Br}(h \to \gamma \gamma) \) is increased by 4% with respect to the SM, resulting in an enhancement of \( \sigma(gg \to h) \times \text{Br}(h \to \gamma \gamma) \) of 18%. This effect of an enhanced \( \gamma \gamma \) signal has been observed for a different parameter space of \( \lambda \text{SUSY} \) in Ref. \cite{3}; however, here the effect is not large and the decays are practically SM-like. Notice in particular the small mixing of the Higgs with the singlet as anticipated in section 3.2. For our \( k' = 10 \text{ TeV} \) point, decay rate times branching ratio of both the heavy Higgs and the singlet into \( WW \) or \( ZZ \) is four orders of magnitude below that of the SM Higgs of the same mass, which makes them invisible to SM Higgs searches. For the fine-tuning parameter, we obtain \( \Delta^{-1} = 19\% \), consistent with expectations based on the general argument.

The heavy Higgs decays into \( A_s A_s, A_s Z, \) and \( \bar{t}t \), with \( A_s \) decaying predominantly into \( b\bar{b} \). The singlet decays into \( A_s A_s, A_s Z, \bar{t}t \), and \( h^+ h^- \). Due to associated \( Z \) production, discovery of these particles may be possible at \( e^+e^- \) colliders such as the ILC/CLIC.
Figure 5: A typical mass spectrum for a compositeness scale of \( k' = 10 \text{ TeV} \). The states with mixing are labeled by their largest components. In the left diagram, the states are always ordered from heavy to light. The gluino mass of \( m_{\tilde{g}} = 946 \text{ GeV} \) in the table corresponds to \( M_3 = 801 \text{ GeV} \) at the scale \( m_{\tilde{t}} \) obtained using the MSSM renormalization group evolution. The wino is relatively heavy, which is necessary to generate a mass for the light pseudoscalar \( m_A > m_h/2 \) through the \( A_\lambda \) term, in line with recent hints of a Higgs discovery. If the wino is much lighter, the Higgs would decay almost entirely to pseudoscalars.

### 3.3.2 \( k' = 1000 \text{ TeV} \)

The fine-tuning constraint will be more severe for \( k' = 1000 \text{ TeV} \) than for \( k' = 10 \text{ TeV} \) because of the large \( \ln(k'/\text{TeV}) = 6.9 \). Performing the same bottom-up analysis as in the case of \( k' = 10 \text{ TeV} \), we find:

- The fine-tuning constraint \( (\Delta^{-1} \gtrsim 10\%) \) requires \( |\mu| \lesssim 290 \text{ GeV} \), \( m_{\tilde{t}} \lesssim 370 \text{ GeV} \) (for degenerate stop masses without mixing), \( m_{\tilde{g}} \lesssim 460 \text{ GeV} \) (again at the leading-log level), \( m_{\tilde{W}} \lesssim 800 \text{ GeV} \), \( m_{\tilde{B}} \lesssim 2500 \text{ GeV} \), and \( m_{\tilde{h}} \lesssim 3.6 \text{ TeV} \).

For \( k' = 1000 \text{ TeV} \), the theory is expected to be fine-tuned at the 10% level.

In Fig. 6, we show a typical mass spectrum for \( k' = 1000 \text{ TeV} \). We find that, as in the \( k' = 10 \text{ TeV} \) case, the phenomenology of the Higgs is mostly SM-like: the production cross section \( \sigma(gg \to h) \) is enhanced by 9% relative to the SM; Eq. (39) gives 0.92 as the decay rate of
Figure 6: A typical mass spectrum for a compositeness scale of $k' = 1000$ TeV. Definitions are as in Fig. 5. The gluino mass of $m_{\tilde{g}} = 814$ GeV corresponds to $M_3 = 710$ GeV at the scale $m_t$ obtained using the MSSM renormalization group evolution.

the Higgs to $b\bar{b}$ with respective to the SM, whereas the exact result is $0.96$. This translates into an increase of $\sigma(gg \rightarrow h) \times \text{Br}(h \rightarrow \gamma\gamma)$ of $13\%$ with respect to the SM. As for the heavy Higgs or the singlet, we again find four orders of magnitude suppression of production cross section times branching ratio into $WW$ or $ZZ$ compared to the SM Higgs with the same mass. We find fine-tuning of $\Delta^{-1} = 10\%$ for this sample point, which is in agreement with expectations.

We find that if we relax our requirement of tuning slightly, we can choose $k'$ to be much larger than 1000 TeV without conflicting with Landau pole constraints. We, however, note that two-loop stop contributions to the Higgs quartic are negative and the theory will therefore require larger $\lambda, \kappa$, so it is not obvious that this statement will hold at two loops. Using the tree-level potential as the other extreme to the one-loop potential, one finds that large $\lambda, \kappa \sim 0.8 - 0.9$ are needed to push the Higgs mass high enough and one cannot take $k'$ much higher than 1000 TeV due to Landau pole constraints. The truth is expected to lie somewhere between the tree-level and one-loop situations.
4 Flat Space Realization

We now discuss realizing our basic setup, Fig. 1, using a flat space extra dimension. An obvious way to do this is to simply turn off the curvature in models of Section 2. The analysis then goes similarly with the replacement $k' \rightarrow 1/\pi R$, except that we now do not have a large desert above the compactification scale, $1/R$, so we cannot have the high-scale see-saw mechanism or conventional gauge coupling unification.

In this section, we pursue an alternative realization, adopting supersymmetry breaking by boundary conditions associated with a compact extra dimension [42]. Our model is essentially that in Ref. [43]. Specifically, we consider an $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ gauge theory in 5D, with the extra dimension compactified on $S^1/Z_2$: $0 \leq y \leq \pi R$. We introduce three generations of matter and Higgs fields in the bulk, but localize the third-generation quark doublet, right-handed top, and Higgs multiplets to the $y = \pi R$ brane:

$$M_{Q_3, U_3, H_u, H_d} \ll -\frac{1}{\pi R}, \quad (40)$$

where $M_\Phi$ represents bulk masses as in previous sections. When supersymmetry is broken by twisted boundary conditions with twist parameter $\alpha$, we obtain

$$m_{\tilde{Q}_3, \tilde{U}_3, H_u, H_d} \ll m_{\tilde{\Psi}_{1,2}, \tilde{D}_3, \tilde{L}_3, \tilde{E}_3} = \frac{\alpha}{R}, \quad (41)$$

at the scale $1/R$, where we have taken

$$|M_{\Psi_{1,2}, D_3, L_3, E_3}| \ll \frac{1}{\pi R}. \quad (42)$$

This condition guarantees that the first two generation sfermions are nearly degenerate in mass, avoiding stringent constraints from flavor violation.

To obtain the spectrum we want, we take $\alpha/R$ to be in the multi-TeV region. For the gauge multiplets, we introduce sizable gauge kinetic terms on (one or both of) branes, which control the size of the gaugino masses:

$$M_A = \frac{\pi R g_{4,A}^2}{g_{5,A}^2} \frac{\alpha}{R}, \quad (43)$$

where $g_{5,A}$ and $g_{4,A}$ are the 5D bulk and 4D gauge couplings, respectively, with $g_{4,A}$ given by

$$\frac{1}{g_{4,A}^2} = \frac{\pi R}{g_{5,A}^2} + \frac{1}{g_0^2} + \frac{1}{g_\pi^2}, \quad (44)$$

in terms of $g_{5,A}$ and the brane-localized gauge couplings at $y = 0$ and $\pi R$, $\tilde{g}_{0,A}$ and $\tilde{g}_{\pi,A}$. We take $M_A$ to be in the sub-TeV region.
Introducing a singlet field $S$ together with the superpotential $\lambda S H_u H_d + f(S)$ on the $y = \pi R$ brane, the analysis of electroweak symmetry breaking goes as in the previous section, with the identification
\[ k' \to \frac{1}{\pi R}, \quad \tilde{m} \to \frac{\alpha}{R}. \quad (45) \]
A negative soft mass-squared for $S$ can be induced, for example, by introducing some bulk field that has a Yukawa coupling to $S$ on the $y = \pi R$ brane.

In the present model, the two circles in Fig. 1 are interpreted as the 5D bulk (left) and the $y = \pi R$ brane (right), rather than the $y = 0$ and $\pi R$ branes as in previous models. Because of Eq. (42), only a part of the Yukawa hierarchy can be explained by wavefunction profiles. With Eqs. (40, 42) the Yukawa matrices obtain the following structure from the wavefunctions:
\[
y_u \sim \begin{pmatrix} \epsilon^2 & \epsilon^2 & \epsilon \\ \epsilon^2 & \epsilon^2 & \epsilon \\ \epsilon & \epsilon & 1 \end{pmatrix}, \quad y_d \sim \begin{pmatrix} \epsilon^2 & \epsilon^2 & \epsilon^2 \\ \epsilon^2 & \epsilon^2 & \epsilon^2 \\ \epsilon & \epsilon & \epsilon \end{pmatrix}, \quad y_e \sim \begin{pmatrix} \epsilon^2 & \epsilon^2 & \epsilon^2 \\ \epsilon^2 & \epsilon^2 & \epsilon^2 \\ \epsilon & \epsilon & \epsilon \end{pmatrix}, \quad (46)
\]
where $O(1)$ factors are omitted in each element, and $\epsilon \equiv 1/\sqrt{\pi R}$ is the volume dilution factor. The structure beyond this must come from that of 5D Yukawa couplings between matter and Higgs on the $y = \pi R$ brane.

5 Conclusions

In this paper we have presented supersymmetric models in which light stops are obtained while keeping near flavor degeneracy for the first two generation sfermions. Such a spectrum is motivated by the naturalness argument together with the recent LHC data. Our construction is based on the basic picture in Fig. 1 being “close to” electroweak symmetry breaking implies being “away from” supersymmetry breaking, and vice versa. In models where the two sectors correspond to the two branes at the opposite ends of a (warped or flat) extra dimension, the desired superpartner spectra are obtained while reproducing the hierarchy in the Yukawa matrices through wavefunction profiles of the quark/lepton fields. A relatively large Higgs boson mass, including $m_h = 125$ GeV, can be easily accommodated if the scale of Kaluza-Klein excitations is low. For models in warped space, the AdS/CFT correspondence allows us to interpret them in terms of purely 4D theories in which the top and Higgs (and the left-handed bottom) multiplets are composites of some strongly interacting sector. An alternative realization of the picture in Fig. 1 is obtained by identifying the two sectors as the bulk of a flat extra dimension and a brane on its boundary, and by breaking supersymmetry by boundary conditions, which we have also discussed.

In the coming years, the LHC will be exploring the parameter regions of supersymmetric theories in which the stops (and the left-handed sbottom) are light. If electroweak symmetry
Figure 7: The Higgs boson mass as a function of $\lambda$ for fixed values of $\tan\beta$ and $\kappa$, given by the exact tree-level formula (black, solid line), the first-order (red, dot-dashed) and second-order (blue, dashed) analytical formulae in Eqs. (47, 48). On the left, the piece-wise exact analytical solutions for $\tan\beta = 1$ are also shown as magenta, dotted lines. The second-order formula gives a very good fit away from the point $\tan\beta = 1$, $\lambda = \lambda_{\text{crit}}$ where the dotted lines cross.

breaking is indeed natural in the conventional sense, the LHC will find the stops in the sub-TeV region. If not, and if the SM-like Higgs boson is confirmed with $m_h \simeq 125$ GeV, then we would be led to consider that supersymmetry is absent at low energies, or it is realized in a somewhat fine-tuned form, perhaps along the lines of scenarios considered in Refs. [44, 45, 46].

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A Analytical Formulae for the Higgs Boson Mass

We first describe the effects of $\lambda$-doublet mixing in the non-decoupling regime, $2\lambda^2 v_u v_d \sim B_\mu + v_u v_d g^2 / 2$, in terms of an expansion in $m_h^2 / m_H^2$ (light Higgs mass over heavy Higgs mass) up to second order. For this purpose, we suspend doublet-singlet mixing in this paragraph; it will be discussed below. We find that to first order in the above mentioned expansion, the light Higgs
mass is given by

\[ m_h^2 \approx M_Z^2 \cos^2 2\beta + \lambda^2 v^2 \sin^2 2\beta \left( 1 - \frac{\lambda^2 v^2 - M_Z^2}{2B_\mu} \sin 2\beta \right), \]  

(47)

where we have used the zeroth order result \( m_H^2 = 2B_\mu / \sin 2\beta \). This approximation is valid to within 10\% for \( \tan \beta \gtrsim 2 \) with \( \kappa = 1, \mu = 200 \text{ GeV} \). Performing the expansion to second order in \( m_h^2 / m_H^2 \), we obtain

\[ m_h^2 \approx m_{h,0}^2 - \frac{2B_\mu - M_Z^2}{2B_\mu - m_{h,0}^2} \sin 2\beta, \]  

(48)

where \( m_{h,0}^2 \) is given by Eq. (47). We find that this second-order expansion gives the correct Higgs mass to within 2\%, 5\%, and 10\% for \( \tan \beta > 2, 1.4, \) and 1.2, respectively, with \( \kappa = 1, \mu = 200 \text{ GeV} \). The analytical formulae are compared with the exact tree-level values in Fig. 7.

As \( \tan \beta \) approaches one, the gap between the light and heavy Higgs masses shrinks to zero at a value of \( \lambda = \lambda_{\text{crit}} \) given by \( \lambda_{\text{crit}} v^2 = B_\mu + M_Z^2 / 2 \). This kink-structure cannot be faithfully described by a perturbative expansion in \( m_h^2 / m_H^2 \). For \( \tan \beta = 1 \), Eq. (35) is an exact solution for \( \lambda < \lambda_{\text{crit}} \), while \( m_h^2 = 2B_\mu + M_Z^2 - \lambda^2 v^2 \) for \( \lambda > \lambda_{\text{crit}} \). In the case of a large Higgs mass, \( m_h^2 \gg M_Z^2 \), a useful expression is

\[ m_h^2 = \frac{1}{2} \left( \frac{B_\mu}{\sin \beta \cos \beta} - \sqrt{(2\lambda^2 v^2 \sin 2\beta - 2B_\mu)^2 + 4B_\mu^2 \cot^2 2\beta} \right). \]  

(49)

We now give an analytic formula for the correction to the Higgs mass from mixing with the singlet in the limit of negligible \( \lambda \)-doublet mixing, \( 2\lambda^2 v_u v_d \ll B_\mu + \bar{g}^2 v_u v_d / 2 \), which corresponds to the doublet-doublet decoupling regime. Performing again an expansion in \( m_h^2 / m_s^2 \), with \( m_s^2 \) the \( \mathcal{M}_{33}^2 \) entry of the scalar mass matrix, one finds to first order

\[ \delta m_h^2 = -\frac{4\mu^2 v^2}{m_s^2} \left( \lambda - \left( \kappa + \lambda \frac{A_\lambda}{2\mu} \right) \sin 2\beta \right)^2. \]  

(50)

In the limit \( A_{\lambda,\kappa} \to 0 \), this agrees with the result in Ref. [34].

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