

An attractor for natural supersymmetry

Timothy Cohen, Anson Hook, and Gonzalo Torroba

*Theory Group, SLAC National Accelerator Laboratory, 2575 Sand Hill Road,
Menlo Park, California 94025, USA*

(Received 13 June 2012; published 3 December 2012)

We propose an attractor mechanism which generates the *more minimal* supersymmetric standard model from a broad class of supersymmetry breaking boundary conditions. The hierarchies in the fermion masses and mixings are produced by the same dynamics and a natural weak scale results from gaugino mediation. These features arise from augmenting the standard model with a new $SU(3)$ gauge group under which only the third generation quarks are charged. The theory flows to a strongly interacting fixed point which induces a negative anomalous dimension for the third generation quarks and a positive anomalous dimension for the Higgs. As a result, a split-family natural spectrum and the flavor hierarchies are dynamically generated.

DOI: [10.1103/PhysRevD.86.115005](https://doi.org/10.1103/PhysRevD.86.115005)

PACS numbers: 12.60.Jv, 14.80.Ly

I. INTRODUCTION AND SUMMARY

The stability of the electroweak scale and the hierarchical structure of the fermion masses and mixing angles are two of the central mysteries of the standard model (SM). It is possible that these puzzles are explained by the same underlying mechanism. One approach is to supersymmetrize the SM and augment it with a new strongly interacting gauge theory. Supersymmetry tames the quadratically divergent contributions to the Higgs mass while the strong dynamics can yield a parametric suppression of the first and second generation Yukawa couplings. Various realizations of this possibility have been proposed so far, including single sector models [1], models of superconformal flavor [2,3], warped extra-dimensional realizations [4], and theories based on deconstruction [5].¹

For some of these constructions, the dynamics which gives rise to the flavor textures also produces an inverted squark hierarchy, where the lightest SM fermions have the heaviest sfermion partners. This provides a microscopic realization of the more minimal supersymmetric SM of Ref. [7], which was motivated by considerations of naturalness and flavor constraints. The phenomenology of these models has been studied thoroughly in e.g., Ref. [8]. Furthermore, the recent LHC bounds on first and second generation squark masses [9] together with attempts to minimize fine-tuning have reinvigorated interest in the phenomenology and collider signatures of such *natural supersymmetry* spectra [10,11].

In this work, we will present a new model to explain the flavor hierarchies which simultaneously yields the natural supersymmetry spectrum and radiative electroweak symmetry breaking (REWSB). This will be accomplished by adding a new strongly coupled conformal sector to the minimal supersymmetric standard model (MSSM). We

will show that starting from rather generic supersymmetry breaking boundary conditions (with some assumptions on certain approximate global symmetries), the infrared theory after escaping from the conformal regime is the more minimal supersymmetric SM.

The MSSM fields are weakly coupled both in the UV and in the IR. The conformal dynamics will generate order one negative anomalous dimensions for the third generation fields once the theory becomes strongly coupled.² Negative anomalous dimensions are only possible if the third generation is charged under this new gauge group—the unitarity bound on dimensions only applies to gauge invariant operators. The third generation Yukawa couplings are marginal operators in the conformal field theory (CFT). These marginal Yukawa couplings will induce a large positive anomalous dimension for the Higgs field. Hence, the remaining Yukawas become irrelevant deformations. It will be shown that this structure can lead to viable flavor hierarchies. Additionally, the strong dynamics will suppress soft masses for the third generation squarks and Higgs fields. Below the exit scale, these will be regenerated by gaugino mediation [12]. The model acts as an attractor for the more minimal supersymmetry spectrum and REWSB.

The goal of this work is to analyze the simplest realization of this mechanism and its main dynamical consequences. The gauge group is $SU(3)_{\text{CFT}} \times SU(3)_X \times SU(2)_W \times U(1)_Y$ where $SU(3)_{\text{CFT}}$ will flow to a strongly coupled fixed point, $SU(3)_X$ is weakly coupled, and $SU(2)_W \times U(1)_Y$ are as in the MSSM. The third generation quark superfields transform under $SU(3)_{\text{CFT}}$. The first and second generations transform under $SU(3)_X$. The $SU(3)$ groups are connected by bifundamental *link* fields. With

¹For some other models which connect the supersymmetry breaking spectrum and flavor, see Ref. [6].

²This differs from previously studied constructions, which relied on large positive anomalous dimensions for the first two SM generations. This can result from compositeness or localization in the IR region of a Randall-Sundrum throat.

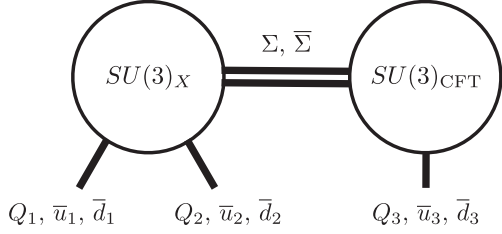


FIG. 1. The model presented here is given by an $SU(3)_X \times SU(3)_{\text{CFT}}$ quiver gauge theory. The node $SU(3)_{\text{CFT}}$ flows to an interacting fixed point and provides the necessary dynamics for generating flavor and an attractor mechanism for natural supersymmetry. $SU(3)_X$ is IR free. The bifundamental link fields Σ and $\bar{\Sigma}$ break the group to the diagonal visible $SU(3)_C$, providing an exit from the conformal regime.

this matter content, $SU(3)_{\text{CFT}}$ is in the conformal window [13]. The link fields eventually acquire a nonzero expectation value causing an exit from the conformal regime; this also breaks $SU(3)_{\text{CFT}} \times SU(3)_X \rightarrow SU(3)_C$, giving rise the visible color interactions.³ This structure is summarized in Fig. 1.

The rest of this paper is organized as follows. Sec. II describes the basic mechanism and its implications for the spectrum of soft masses and the flavor hierarchies. In Sec. III, we discuss the low energy phenomenology, including some general remarks about the spectrum and the Higgs sector. We also provide some concrete example spectra. Our conclusions and future directions are presented in Sec. IV.

II. THE MODEL

We begin by describing the model of Fig. 1. For simplicity, we will ignore the leptons, which do not affect our discussion other than to ensure anomaly cancellation. The matter content and charge assignments are given in Table I. The third generation quarks are charged under $SU(3)_{\text{CFT}}$, while the first two generations are charged under $SU(3)_X$. The bifundamental link fields are denoted by Σ and $\bar{\Sigma}$. The field A is an adjoint plus a singlet of $SU(3)_X$. The superpotential will be chosen so that the F term for A forces $\langle \Sigma \rangle = \langle \bar{\Sigma} \rangle \neq 0$. This will cause an escape from the conformal regime while also giving masses to the bifundamentals. The superpotential contains the following relevant terms:

$$W \supset Q_3 H_u \bar{u}_3 + Q_3 H_d \bar{d}_3 + \Sigma A \bar{\Sigma} + W_{\mathcal{U}(1)}. \quad (1)$$

Contractions over gauge indices are implicit. $W_{\mathcal{U}(1)}$ will be instrumental in breaking some of the Abelian symmetries

³The super top color model of Ref. [14] utilizes a similar group structure and matter content. However, unlike models of top color (see Ref. [15] for a review), the mechanism studied in this work does not utilize top condensation to break electroweak symmetry.

TABLE I. The particle content and charge assignments for the MSSM quark and CFT sectors. The subscripts denote generation assignments, and the leptons are charged as in the MSSM. The visible color gauge group is a diagonal subgroup of $SU(3)_{\text{CFT}} \times SU(3)_X$.

	$SU(3)_{\text{CFT}}$	$SU(3)_X$	$SU(2)_W$	$U(1)_Y$
Q_3	$\square$	1	$\square$	1/6
$\bar{d}_3$	$\bar{\square}$	1	1	1/3
$\bar{u}_3$	$\bar{\square}$	1	1	-2/3
H_u	1	1	$\square$	1/2
H_d	1	1	$\square$	-1/2
Σ	$\square$	$\bar{\square}$	1	0
$\bar{\Sigma}$	$\bar{\square}$	$\square$	1	0
A	1	1 + adj	1	0
$Q_{2,1}$	1	$\square$	$\square$	1/6
$\bar{d}_{2,1}$	1	$\bar{\square}$	1	1/3
$\bar{u}_{2,1}$	1	$\bar{\square}$	1	-2/3

which can spoil the desired low energy spectrum. We will discuss this term in detail below.

With this matter content, $SU(3)_{\text{CFT}}$ has five flavors and flows to a strongly interacting superconformal fixed point in the IR. The crossover scale below which this theory becomes strong is denoted by Λ_{CFT} . The remaining gauge groups are IR free and act as spectators to this strong dynamics. A crucial property of the model is that the third generation Yukawa couplings appear as relevant interactions in the CFT. The Higgs fields will then also be part of the CFT—they will receive a positive anomalous dimension. These couplings, as well as the rest of the interactions in Eq. (1), will naturally flow to order one values below Λ_{CFT} . In contrast, the remaining Yukawas will arise as irrelevant deformations, resulting in a flavor hierarchy between the third and first two generations.

If we do not add extra fields, this matter content spoils gauge coupling unification. However, there are no issues with Landau poles up to the grand unified theory scale, and one could imagine UV completing the model using full $SU(5)$ representations. We will come back to this point briefly in Sec. IV, while here we continue to focus on this minimal realization.

The energy scales in our model are as follows, see Fig. 2. At the messenger scale M , soft supersymmetry breaking operators are generated. The supersymmetry breaking mechanism and mediation can be arbitrary, up to certain assumptions on global symmetries which we describe

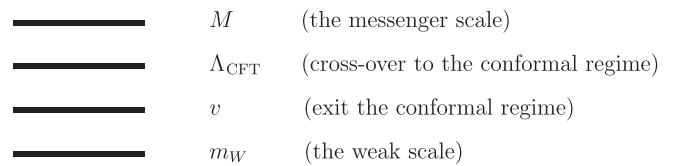


FIG. 2. The relevant scales for our model.

below. The scale M could be above or below Λ_{CFT} , but the physical soft masses should be smaller than Λ_{CFT} so that the superconformal dynamics dominate. At a scale $v < \Lambda_{\text{CFT}}$, we exit the CFT regime. This is done supersymmetrically by adding

$$W \supset -v^2 \text{Tr} A, \quad (2)$$

to Eq. (1). This new scale can be generated dynamically as explained in Ref. [16]. The link fields acquire an expectation value $\langle \Sigma \tilde{\Sigma} \rangle = v^2$, which breaks $SU(3)_{\text{CFT}} \times SU(3)_X \rightarrow SU(3)_C$. The visible gauge coupling becomes

$$\frac{1}{g_C^2} = \frac{1}{g_X^2} + \frac{1}{g_{\text{CFT}}^2}, \quad (3)$$

which is dominated by $g_X^2 \simeq g_C^2$. We assume that the exit from the conformal regime happens quickly, such that at energy scales $E < v$, a perturbative description is valid. As we show below, the weak scale $m_W < v$ is radiatively generated. We note that, in contrast with composite models, here, the MSSM fields are weakly coupled both in the UV (above Λ_{CFT}) and in the IR (below the exit scale).

A. An attractor for natural supersymmetry

In this section, we will analyze the conformal regime and how it affects the soft masses. For more details, see Ref. [17] and the references therein. We will first neglect the effects from the weakly interacting gauge groups and the first two generations. This corresponds to setting $g_{\text{SM}} \rightarrow 0$ and ignoring mixings from Yukawa couplings. We will then show that such effects amount to small finite corrections.

Our main dynamical assumption is that the fixed point is stable, which means that small perturbations of the couplings away from their fixed point value are irrelevant. Equivalently, the matrix $\partial \beta_i / \partial y_j$ must be positive definite, where y_i are the couplings of the theory and β_i are the corresponding beta functions. With this assumption, all the physical couplings flow to their fixed point values, and their higher θ components flow to zero. This can be seen by promoting the couplings to background superfields. One implication is that all soft supersymmetry breaking terms associated with relevant couplings are highly suppressed by the conformal dynamics.⁴

To understand the consequences for our model, consider a relevant superpotential interaction

$$W \supset \lambda \prod_i \Phi_i^{n_i}, \quad (4)$$

for superfields Φ_i and positive integers n_i . The physical coupling is

⁴For a model which uses this mechanism to suppress the Higgs soft mass, see Ref. [18].

$$\lambda_{\text{phys}} = \lambda \prod_i (Z_{\Phi_i}^{n_i})^{-1/2}, \quad (5)$$

where the Z_{Φ_i} are the wave function renormalizations for the superfields Φ_i and encode the soft masses as their θ^4 components. As λ_{phys} flows to its fixed point value, its θ^4 component flows to zero. Equivalently, this implies that the combination of soft masses

$$\sum_i n_i \tilde{m}_i^2, \quad (6)$$

flows to zero at the fixed point, where $\tilde{m}_i$ is the soft mass for Φ_i . Since the θ^2 component also flows to zero, the same conclusion holds for the a terms.

Similarly, promoting the gauge coupling to a superfield implies that the CFT gaugino mass and

$$\sum_r \text{dim}(r) T_r \tilde{m}_r^2 \quad (7)$$

are also suppressed by the CFT dynamics. Here, the field Φ_r has index T_r under the gauge group, e.g., $T(\square) = 1/2$, and $\text{dim}(r)$ is the dimension of Φ_r for a fixed gauge index.

As we mentioned above, we assume that the CFT is IR attractive, which means that the eigenvalues λ_i of the matrix $\partial \beta_i / \partial g_j$ are positive and, generically at a strongly coupled fixed point, order one. The previous soft parameters are then suppressed by a power law $(\frac{\mu}{\Lambda_{\text{CFT}}})^{\lambda_i}$, where μ is the renormalization group (RG) scale. This effect can be seen explicitly in weakly coupled examples such as the Wilson-Fisher fixed point. Below, we will take into account the small contributions from the perturbative SM couplings.

On the other hand, due to the nonrenormalization of conserved currents, combinations of soft masses proportional to conserved $U(1)$ symmetries,

$$\sum_i \text{dim}(i) q_i \tilde{m}_i^2 \quad (8)$$

are not renormalized by the strong dynamics. Here, q_i denotes the $U(1)$ charge. This effect will be an important constraint on the viability of obtaining the more minimal supersymmetric SM.

Our goal is to use the conformal dynamics to suppress the soft masses for the third generation squarks and Higgs fields, which is essentially the idea of conformal sequestering [19]. We must explicitly break some of the non-anomalous global $U(1)$ symmetries. Otherwise, Eq. (8) shows that they would lead to unsuppressed tachyonic soft masses. We accomplish this with the term $W_{\hat{U}(1)}$ in the superpotential of Eq. (1). As a concrete example, let us investigate a specific choice:

$$W_{\hat{U}(1)} = (Q_3 \bar{u}_3)(Q_3 \bar{d}_3). \quad (9)$$

TABLE II. The global anomaly-free $U(1)$ symmetries for the model given by Eq. (1) with the $U(1)$ breaking superpotential in Eq. (9). The charge assignments for the gauged symmetries are given in Table I.

	$U(1)_1$	$U(1)_2$	$U(1)_3$	$U(1)_R$
Q_3	1	0	0	1/2
$\bar{u}_3$	-1	-1	0	1/2
$\bar{d}_3$	-1	1	0	1/2
H_u	0	1	0	1
H_d	0	-1	0	1
Σ	0	0	1	1/3
$\bar{\Sigma}$	0	0	-1	1/3
A	0	0	0	4/3

In this theory, the superconformal R charges are uniquely determined in terms of symmetries and anomaly cancellation. The anomaly-free Abelian symmetries are given in Table II.

One can verify self-consistently that all terms in the superpotential in Eq. (1) have R charge 2 at the fixed point. The remaining first and second generation fields decouple from the strong dynamics; they are neutral under the non- R symmetries and are (approximately) free fields.

It is useful to explain in more detail how the flow to this fixed point proceeds, starting from the UV free theory. In the UV, the cubic terms in the superpotential are classically marginal, but the quartic symmetry breaking term is classically irrelevant. In terms of the canonical UV fields, $W_{\mathcal{U}(1)} = \frac{1}{M_*} (Q_3 \bar{u}_3)(Q_3 \bar{d}_3)$, where M_* is some large mass scale. First, consider the limit $M_* \rightarrow \infty$. The resulting theory is supersymmetric QCD with extra singlets and cubic superpotential deformations (SSQCD). Below the strong coupling scale for this model Λ_{SSQCD} , the theory flows to a superconformal fixed point. This CFT can be studied using a maximization [20]; we find that the superconformal R charges of Q_3 and $(\bar{u}_3, \bar{d}_3)$ are $2/3 \times \sqrt{2/7}$. The cubic interactions are then relevant perturbations of the free fixed point, driving the theory to the nontrivial SSQCD fixed point.

Next, we can add the quartic superpotential, taking $M_*/\Lambda_{\text{SSQCD}}$ large but finite. The theory first flows close to the SSQCD fixed point which we just described. According to the previous R charges, in this regime, the quartic operator $(Q_3 \bar{u}_3)(Q_3 \bar{d}_3)$ is relevant. So for any non-zero value of $M_*/\Lambda_{\text{SSQCD}}$, it will drive the theory away from the SSQCD fixed point. The crossover scale Λ_{CFT} at which such effects become important is of order $\Lambda_{\text{CFT}}^{3-4\Delta} \sim \Lambda_{\text{SSQCD}}^{4-4\Delta}/M_*$. Below this scale, the fixed point value of the quartic coupling is order one irrespective of the initial M_* , and we recover the CFT with the R charges given in Table II. We note that for larger M_* , it takes longer to flow to this fixed point; however, since the fixed point values are always order one, the $U(1)$ symmetry breaking

term will suppress the unwanted soft masses as long as $\tilde{m} < \Lambda_{\text{CFT}}$.⁵

Having analyzed the RG (RG) evolution toward our fixed point, let us return to the behavior of the soft parameters. The only combinations of soft masses which are not suppressed as the fixed point is approached are

$$\begin{aligned} & \tilde{m}_{\Sigma}^2 - \tilde{m}_{\bar{\Sigma}}^2 \\ & 2\tilde{m}_{Q_3}^2 - \tilde{m}_{\bar{u}_3}^2 - \tilde{m}_{\bar{d}_3}^2 \\ & \tilde{m}_{H_u}^2 - \tilde{m}_{H_d}^2 + \tilde{m}_{\bar{d}_3}^2 - \tilde{m}_{\bar{u}_3}^2. \end{aligned} \quad (10)$$

This implies that for arbitrary UV boundary conditions, the model does not fully sequester soft masses. However, if the supersymmetry breaking mechanism preserves approximate charge conjugation and custodial symmetries, then the contributions from Eq. (10) are negligible at the messenger scale and are not generated by the strong dynamics. This is the case in minimal gauge mediation [22], where at the messenger scale, the first difference in Eq. (10) vanishes identically, while the linear combinations in the second and third lines are much smaller than each of their respective terms. These combinations can also be suppressed by going beyond minimal gauge mediation or in gravity mediation by imposing discrete symmetries.

As we noted before, this analysis neglects effects from the weakly interacting sector of the theory. The first two generations and the SM gauginos continuously feed supersymmetry breaking contributions to the CFT fields, giving rise to *driving terms* in the beta functions for the CFT superfield couplings. However, these supersymmetry breaking effects are much smaller than the soft masses of the first two generation sfermions and gauginos, since the CFT couples to such fields only through irrelevant interactions. Specifically, they are suppressed by loop factors and by SM gauge couplings or Yukawa interactions. These corrections will be taken into account in Sec. III.

Hence, under the assumption that the supersymmetry breaking mechanism (approximately) respects the above symmetries, the strong conformal dynamics fully suppresses the soft masses of the third generation quarks and Higgs fields, up to small corrections from the weakly coupled sector. It would be interesting to modify the model to accomplish a complete sequestering without having to assume these symmetries, e.g., by adding new flavors and turning on different deformations. Some of these possibilities will be discussed briefly in Sec. IV.

Finally, at the scale ν , we exit the conformal regime. This happens in an approximately supersymmetric way and does not lead to appreciable finite corrections for the soft parameters. Therefore, the theory at energies below ν

⁵This analysis does not conflict with the results of Ref. [21] since $(Q_3 \bar{u}_3)(Q_3 \bar{d}_3)$ is not a chiral primary operator at the fixed point with superconformal R charges given in Table II. We thank Dan Green for discussions on this point.

is the weakly coupled MSSM with the soft masses for the third generation fields and Higgses suppressed with respect to the first two generations and the gauginos. The third generation squark masses are then regenerated by gaugino mediation [12], which in turn can drive the up-type Higgs soft mass squared negative. Starting from generic supersymmetry breaking mechanisms, the CFT acts an attractor for realizing the natural supersymmetry spectrum of the more minimal supersymmetric SM. The resulting phenomenology will be studied in Sec. III.

B. Generating the flavor hierarchies

Starting from the pioneering work of Nelson and Strassler [2], it has been understood how CFT dynamics can generate the flavor hierarchies at low energies from arbitrary order one Yukawas in the UV. We will explain how this works in the context of our construction. The model presented here differs from previous models of compositeness/superconformal flavor since the third generation superfields have negative anomalous dimensions.

Above the dynamical scale Λ_{CFT} , the renormalizable Yukawa couplings are

$$W \supset Y_{ij}^u Q_i H_u \bar{u}_j + Y_{ij}^d Q_i H_d \bar{d}_j + Y_{33}^u Q_3 H_u \bar{u}_3 + Y_{33}^d Q_3 H_d \bar{d}_3, \quad (11)$$

where $i, j = 1, 2$ and all the coefficients are taken to be order one. Renormalizable mixing terms between the third generation and the first two are forbidden by gauge invariance. They will be generated by irrelevant operators as we explain below.

The third generation Yukawas are relevant in the CFT regime. Below Λ_{CFT} , they flow to order one fixed values. In contrast, the first two generation Yukawas become irrelevant because the Higgs fields acquire positive anomalous dimension. For energies $v < E < \Lambda_{\text{CFT}}$, we find⁶

$$Y_{ij}^u(E) = \left(\frac{E}{\Lambda_{\text{CFT}}} \right)^{\frac{\gamma_{Q_i} + \gamma_{u_j} + \gamma_{H_u}}{2}} Y_{ij}^u(\Lambda_{\text{CFT}}), \quad (12)$$

and a similar expression for Y^d . Defining the ratio between the exit scale and dynamical scale

$$\epsilon \equiv \frac{v}{\Lambda_{\text{CFT}}}, \quad (13)$$

the suppression in the first two generation Yukawas at the exit scale becomes

$$Y_{ij}(v) = \epsilon^{\frac{\gamma_H}{2}} Y_{ij}(\Lambda_{\text{CFT}}) \ll Y_{33}(v). \quad (14)$$

We have neglected the perturbative anomalous dimensions for the first two generations. This dynamically generates a

hierarchy between the first/second and the third generation Yukawa couplings.

Next, we consider the off-diagonal Yukawa interactions between the third and first/second generations. The lowest dimension operators allowed by gauge invariance are of the form

$$W \supset \frac{1}{\Lambda_*} \bar{\Sigma} Q_3 H_u \bar{u}_{1,2} + \frac{1}{\Lambda_*} Q_{1,2} H_u \Sigma \bar{u}_3 + \dots, \quad (15)$$

where Λ_* is the scale at which these operators are generated. These lead to off-diagonal Yukawas after setting $\langle \Sigma \bar{\Sigma} \rangle = v^2$ at the exit scale. The RG evolution between Λ_{CFT} and v yields

$$Y_{i3}^u(v) = \frac{v}{\Lambda_*} \epsilon^{\frac{\gamma_{H_u} + \gamma_{Q_3} + \gamma_{\bar{\Sigma}}}{2}}, \quad Y_{3i}^u(v) = \frac{v}{\Lambda_*} \epsilon^{\frac{\gamma_{H_u} + \gamma_{u_3} + \gamma_{\Sigma}}{2}}. \quad (16)$$

Note that the theory near the UV free fixed point contains two types of classically irrelevant operators: the $U(1)$ symmetry breaking term Eq. (9) and the interactions Eq. (15). However, their IR fate is very different. As we showed before, the interaction $W_{U(1)}$ becomes relevant in the IR, driving the theory to a strongly coupled fixed point (where it becomes order one), while Eq. (15) is irrelevant along the whole flow toward the fixed point. Such irrelevant perturbations do not modify the RG flow or the suppression of soft parameters. They become marginal only after the exit of the conformal regime.

Combining these results, we obtain the following the Yukawa textures at v :

$$Y^u \sim \begin{pmatrix} \epsilon^{\frac{\gamma_{H_u}}{2}} & \epsilon^{\frac{\gamma_{H_u}}{2}} & \xi_Q \epsilon^{\frac{\gamma_{H_u}}{2}} \\ \epsilon^{\frac{\gamma_{H_u}}{2}} & \epsilon^{\frac{\gamma_{H_u}}{2}} & \xi_Q \epsilon^{\frac{\gamma_{H_u}}{2}} \\ \xi_u \epsilon^{\frac{\gamma_{H_u}}{2}} & \xi_u \epsilon^{\frac{\gamma_{H_u}}{2}} & 1 \end{pmatrix}, \quad (17)$$

with $\epsilon \ll 1$ defined in Eq. (13) and

$$\xi_Q \equiv \frac{v}{\Lambda_*} \epsilon^{\frac{\gamma_{\bar{\Sigma}} + \gamma_{Q_3}}{2}}, \quad \xi_u \equiv \frac{v}{\Lambda_*} \epsilon^{\frac{\gamma_{\Sigma} + \gamma_{u_3}}{2}}. \quad (18)$$

A similar expression holds for Y^d . Choosing Λ_* below the dynamical scale of the CFT and requiring negative $\gamma_{\bar{\Sigma}} + \gamma_{Q_3}$ and $\gamma_{\Sigma} + \gamma_{u_3}$ (as is the case in our model) gives $\xi_{Q,u} \gtrsim 1$.

For the model with the superpotential given in Eq. (9), the anomalous dimensions which determine the Yukawa couplings are $\gamma_{H_u} = \gamma_{H_d} = 1$ and $\gamma_{Q_3} + \gamma_{\bar{\Sigma}} = \gamma_{u_3} + \gamma_{\Sigma} = -3/2$. The flavor hierarchies between the third and second generations can be generated when

$$\frac{v}{\Lambda_{\text{CFT}}} \sim 10^{-4}, \quad \frac{\Lambda_*}{\Lambda_{\text{CFT}}} \sim 10^{-1} - 10^{-2}. \quad (19)$$

This model does not explain why the first generation Yukawa is smaller than the second generation one. However, this additional small Yukawa could arise by an

⁶The anomalous dimension is defined as $\Delta = 1 + \gamma/2$. For chiral primary operators, it is related to the superconformal R charge by $\gamma = 3R - 2$.

accidental degeneracy of Eq. (17), or by approximate flavor symmetries as in Ref. [23]. We have checked that by scanning over order one coefficients, we can reproduce the quark spectrum and the Cabibbo-Kobayashi-Maskawa matrix to a good approximation.

This ends the general analysis of our mechanism. The rest of the work is devoted to a study of its phenomenological consequences.

III. LOW ENERGY PHENOMENOLOGY

Having explained the main features of our mechanism, we will now analyze the properties of the spectrum and Higgs sector and the parameter ranges which lead to a realistic low energy phenomenology.

A. General properties of the spectrum

In this section, we discuss the features of the low energy spectrum in models which use the dynamics described in Sec. II. Supersymmetry breaking is communicated to the MSSM at the messenger scale M , where the operators

$$c_f^2 \int d^4\theta \frac{X^\dagger X}{M^2} \Phi_{\text{SM}}^\dagger \Phi_{\text{SM}}, \quad c_{\tilde{g}} \int d^2\theta \frac{X}{M} \mathcal{W}_\alpha \mathcal{W}^\alpha, \dots \quad (20)$$

are generated, where X is a supersymmetry breaking spurion with $\langle X \rangle \supset \theta^2 F$, Φ_{SM} is an MSSM matter superfield, $\mathcal{W}_\alpha$ is the field strength for an MSSM gauge group, and the factors of c are model dependent coefficients. These terms give sfermion and gaugino masses which are determined by the F term of X . Supersymmetry breaking is external to the dynamics described in Sec. II, and we do not constrain the soft UV boundary values, up to the assumptions on approximate symmetries required to suppress the differences given in Eq. (10).

Typically, in concrete models of supersymmetry breaking, the sfermion masses at the messenger scale are comparable for the three generations. On the other hand, sfermion and gaugino masses need not arise at the same order in F/M . This happens in many known cases. For instance, an approximate R symmetry or gaugino screening (which occurs for a wide class of gauge mediated models [24]) can lead to subleading gaugino masses. We will assume that gauginos are around the TeV scale. In principle, the sfermions can be much heavier at the messenger scale, but we do require that

$$\tilde{m}_f \sim c_f \frac{F}{M} \ll \Lambda_{\text{CFT}}, \quad (21)$$

so that the conformal dynamics will be relatively unperturbed.

Generic sfermion masses will lead to flavor changing neutral currents (FCNCs). In our setup, flavor problems can be somewhat alleviated by having heavy enough sfermions, while also imposing some degree of degeneracy

between the first two generations.⁷ In this case, $\tilde{m}_{f_{1,2}} \gtrsim \mathcal{O}(10 \text{ TeV})$ avoids dangerous FCNCs. On the other hand, there is a limit on how heavy the first two generations can be so that the third generation sfermion masses do not become tachyonic [25]. To account for this constraint, we include the dominant 2-loop contributions from the heavy states in the analysis of Sec. III C. It would also be interesting to study models where the CFT dynamics alleviates such tachyonic contributions, allowing a more complete decoupling of the first two generation sfermions.

Once we enter the conformal regime, the soft masses for the third generation sfermions and Higgs fields are renormalized by the strong dynamics as described in Sec. II, while the first two generation sfermions and gauginos are not appreciably modified. When evaluating the running of the soft parameters in the conformal regime, we must consider that gauginos and first/second generation fields are continuously feeding supersymmetry breaking contributions into the third generation and Higgs fields [17]. For most of the viable parameter space, the dominant contribution comes from the gauginos, leading to finite contributions

$$\tilde{m}_{\text{CFT}}^2 \sim \frac{g_X^2}{16\pi^2} |M_3|^2, \quad (22)$$

where M_3 is the MSSM gluino mass and g_X is the gauge coupling of the weakly interacting $SU(3)_X$.

After escaping the conformal regime, we find a soft spectrum with $\tilde{m}_{1,2} \sim \mathcal{O}$ (few TeV), $M_3 \sim \mathcal{O}(1 \text{ TeV})$, and small masses for the third generation squarks and Higgs fields. The masses for the light fields are then predominantly regenerated by gaugino mediation [12] between ν and the electroweak scale. The gauginos drive the stop mass to positive values. For $\nu \gtrsim 50 \text{ TeV}$, this makes the up-type Higgs tachyonic and triggers electroweak symmetry breaking. The RG evolution will be studied explicitly below.

B. Comments on the Higgs sector

Next, we discuss the interplay between the MSSM Higgs sector and our model. First, we consider the supersymmetric Higgs mass μ and the bilinear supersymmetry breaking Higgs mass b_μ . Our model contains a solution to the μ problem via the irrelevant interaction

$$W \supset \frac{1}{\Lambda_\mu} \Sigma \bar{\Sigma} H_u H_d. \quad (23)$$

(The tree level μ term $W \supset H_u H_d$ can be forbidden by symmetries). The operator in Eq. (23) can be generated by the same mechanism which produces the off-diagonal Yukawas in Eq. (15). This is another interesting connection

⁷This is satisfied automatically if the mediation mechanism is flavor-blind.

between flavor textures and the Higgs sector. Assuming this occurs, $\Lambda_\mu \sim \Lambda_*$, and no new scale is needed.⁸ Taking into account the CFT suppression, the μ term at the exit scale becomes

$$\mu = \left(\frac{v}{\Lambda_*} \epsilon^{\gamma_H + \gamma_\Sigma} \right) v. \quad (24)$$

In this approach, b_μ is zero at v and is generated radiatively as we run down to the weak scale. In the leading log approximation,

$$b_\mu \simeq -\frac{1}{16\pi^2} \mu \left(6g_W^2 M_2 + \frac{6}{5} g_Y^2 M_1 \right) \log\left(\frac{v}{m_W}\right), \quad (25)$$

where M_1 is the bino mass and M_2 is the wino mass. This solution to μ and b_μ can lead to REWSB.

For our model, $\gamma_H + \gamma_\Sigma = 0$. Requiring $\mu \sim 100$ GeV and using the approximate values in Eq. (19),

$$\begin{aligned} v &\sim 100 \text{ TeV}, \\ \Lambda_* &\sim 10^4\text{--}10^5 \text{ TeV}, \quad \text{and} \\ \Lambda_{\text{CFT}} &\sim 10^6 \text{ TeV}. \end{aligned} \quad (26)$$

While this is an attractive solution to the μ problem, when coupling our mechanism to a specific supersymmetry breaking model, there could be additional dynamics which explains μ/b_μ . In this case, it is not necessary to introduce Eq. (23), and the scales, Eq. (26), could be different.

We now discuss the physical Higgs mass. Below the exit scale, the gluino mass will drive the stop mass positive, which in turn contributes negatively to $\tilde{m}_{H_u}^2$. As long as the bino and wino masses are not too large, this will trigger electroweak symmetry breaking. Models with unified gauginos provide an example of successful REWSB. The down-type Higgs soft mass will be generated through a combination of competing effects from the sbottom and the heavy first/second generations (which drive it negative), and the bino and wino (which drive it positive).

Since the mechanism described in this work yields light stops and negligible a terms, there is tension with a physical Higgs mass of order 125 GeV, as currently hinted at by the LHC [26]. Thus, a realistic model must include an additional source to raise the physical Higgs mass. In the simplest version of our construction, a next to minimal supersymmetric standard model (NMSSM) type extension does not solve this problem because the CFT makes the interaction $W \supset SH_u H_d$ (with S the extra singlet in the NMSSM) irrelevant. This leads to a negligible increase in the physical Higgs mass. One option beyond singlet extensions would be to add *nondecoupling D terms* [27] below the exit scale. While we do not attempt to embed this or other mechanisms into our model, we see no

⁸One can also imagine a different discrete symmetry such that $\Sigma^3 H_u H_d$ is the lowest dimension operator which could generate an effective μ term.

fundamental obstruction. The validity of our conclusions require that this additional module does not lead to appreciable shifts for any of the soft masses.

C. An example spectrum

In order to perform a concrete analysis, we will work in the context of a model with unified gaugino masses. We will also assume that the mediation of supersymmetry breaking respects custodial symmetry and a *charge conjugation* symmetry between Q and $\bar{u}, \bar{d}$, i.e., $\tilde{m}_{Q_3}^2 = \tilde{m}_{\bar{u}_3}^2 = \tilde{m}_{\bar{d}_3}^2$. For example, both of these assumptions are well approximated by models of minimal gauge mediation. The following analysis demonstrates in a concrete setup the viability of the mechanism for splitting the third generation from the first and second. The techniques presented here can be applied to a wide class of supersymmetry breaking scenarios.

Given this framework, the spectrum is determined by choosing a gluino mass and solving the RG equations with the boundary condition at the scale v which the third generation and Higgs soft masses are given by Eq. (22). While there is an incalculable order one coefficient, we find that such effects are small in the regime of interest. If we also assume the solution to the μ problem proposed in Sec. III B, the exit scale is fixed at $v \sim \mathcal{O}(100)$ TeV. The model is then very predictive: all we need to specify are the messenger scale, gaugino, and first/second generation masses.

As an example, we find the viable spectrum presented in Table III, with first/second generation sfermion masses chosen to be 5 TeV. We have assumed an additional contribution to the Higgs quartic from a coupling g_{new} so that

$$m_Z^2 = \frac{g_Z^2}{2} (\langle H_u \rangle^2 + \langle H_d \rangle^2) \rightarrow \Xi^2 \equiv \frac{g_Z^2 + g_{\text{new}}^2}{2} (\langle H_u \rangle^2 + \langle H_d \rangle^2), \quad (27)$$

in all tree-level MSSM expressions for electroweak symmetry breaking and the Higgs sector. In our numerical analysis below, we will take $\Xi \simeq 150$ GeV. As discussed in Sec. III B, this could in principle arise from a

TABLE III. An example set of consistent parameters with the solution to the μ problem given in Eq. (24). We have assumed gaugino mass unification and to good approximation $\tilde{m}_{Q_i}^2 = \tilde{m}_{\bar{u}_i}^2 = \tilde{m}_{\bar{d}_i}^2 = \tilde{m}_i^2$ at low energies. We find that the tree-level value of the Higgs mass is $\simeq 105$ GeV which is consistent with 125 GeV when loop corrections are taken into account.

v	M_3	M_2	M_1
350 TeV	2.5 TeV	1.0 TeV	530 GeV
$\tilde{m}_3^2$	$\tilde{m}_{1,2}^2$	$\tilde{m}_{H_u}^2$	$\tilde{m}_{H_d}^2$
(1.2 TeV) ²	(5 TeV) ²	-(220 GeV) ²	(300 GeV) ²
μ	b_μ	M_A	$\tan\beta$
220 GeV	-0.030 GeV ²	135 GeV	4.2

nondecoupling D term—we are agnostic about its source and find that this leads to a small change for all the parameters except for the physical value of the CP even Higgs masses. This yields a Higgs mass of 105 GeV at tree level which (given the stop masses in Table III) will lead to a mass consistent with 125 GeV once loop corrections are taken into account. This point is also consistent with the relevant experimental bounds considered in Sec. III D below. This demonstrates the viability of our mechanism.

D. Exploring the parameter space

In this subsection, we will briefly explore the possible range of predictions for the soft mass spectrum. In order to do this, we will relax the relationship between the μ term and ν given in Eq. (24). Noting that in our concrete model the coupling $W \supset H_u H_d$ is exactly marginal at the fixed point, one can in principle generate μ and b_μ using an unrelated mechanism at scales above Λ_{CFT} . We can thus take ν and $\tan\beta$ as free parameters and explore the resultant phenomenology. In Fig. 3, we have plotted the low energy values of $\tilde{m}_{Q_3}^2 \simeq \tilde{m}_{u_3}^2 \simeq \tilde{m}_{d_3}^2$ [black, solid], m_A (with $\tan\beta = 2$) [red, dashed], and m_A (with $\tan\beta = 10$) [orange, dotted-dashed] for two choices of ν as a function of the gluino mass. The mass of the A is the only parameter with a strong dependence on $\tan\beta$. As in Sec. III C, we assume that supersymmetry breaking respects $\tilde{m}_{Q_3}^2 = \tilde{m}_{u_3}^2 = \tilde{m}_{d_3}^2$ to a good approximation.

In order to generate this plot, we use the RG equations for the MSSM to flow from ν to the weak scale including the leading 2-loop contributions from the first and second generation sparticles which we fix at 5 TeV. It is this choice which causes the third generation squarks to become tachyonic for small gluino masses. This is the excluded region plotted in opaque grey in Fig. 3. The opaque blue region is excluded due to a lack of REWSB (these conditions are unchanged from the MSSM). The light translucent green region is excluded due to the LEP bound on the A mass⁹ [28]. This constraint is cut off by the kinematic reach of LEP for the process $e^+e^- \rightarrow hA$. For $m_h = 125$ GeV (115 GeV), this implies that $m_A \gtrsim 90$ GeV ($m_A \gtrsim 100$ GeV). As a conservative estimate, we impose $m_A > 100$ GeV in Fig. 3. We have not included the bino, wino, and first/second generation soft masses in Fig. 3 since they are unaffected by our mechanism up to small effects due to off-diagonal Yukawa couplings and 2-loop diagrams.

The bounds for the $\tan\beta = 10$ case are $M_3 > 0.85(0.92)$ TeV to avoid having tachyonic squarks, $M_3 > 1.2(1.1)$ TeV for REWSB and $M_3 > 2.1(1.9)$ TeV for the A mass, given $\nu = 10^3(10^6)$ TeV. Note that the

⁹This exclusion is highly dependent on $\tan\beta$. Furthermore, one could imagine a model where H_d is not a part of the CFT. It would have a large mass, and the model would generically be in the decoupling limit.

LHC also places strong bounds on m_A from searches for ditau resonances [29]. In fact, the LHC excludes the range $120 \text{ GeV} \lesssim m_A \lesssim 220 \text{ GeV}$ for $\tan\beta = 10$ in the context of the MSSM (with $\Xi = m_Z$). We do not show these constraints in Fig. 3 since the excluded regions are for the $\tan\beta = 2$ example.

Recall that achieving a Higgs mass of 125 GeV requires physics beyond the simple model proposed here. Hence, we will only make a few comments about the mass of the Higgs. First, we note that A is light in the region of parameter space with the lightest squarks and gluino, which can have a nontrivial impact on the mass and couplings of the h . In the pure MSSM, this manifests as a dependence on both $\tan\beta$ and the Higgs mixing angle α for the Higgs couplings (for a review, see Ref. [30]). More generally, the dependence of the Higgs couplings on m_A is model dependent. It would be interesting to develop a realistic model for the Higgs sector based on our general mechanism, where this and related questions could be addressed in detail.

In generating Fig. 3, we took $\Xi = 150$ GeV, see Eq. (27); we find only mild sensitivity to the choice of Ξ . When $m_A < \Xi$, $m_h \simeq m_A \cos(2\beta)$, independent of Ξ . For the choice $\tan\beta = 10$, the one-loop corrections from the stops are approximately of the right size to generate a Higgs mass of 125 GeV in the allowed window $100 \text{ GeV} \lesssim m_A \lesssim 120 \text{ GeV}$. For larger values of m_A , the tree-level contribution to the Higgs boson mass would be set by Ξ , which could be carefully chosen to reproduce the desired result. Alternatively, one could attempt to alter the Higgs quartic with a different mechanism than the one captured by our parameter Ξ .

Since we have a splitting between the first/second and third generation squarks, we must worry about FCNC effects induced by rotating the Yukawa matrices of Eq. (17) to the physical basis. To leading order in ϵ , the relevant 1–3 and 2–3 mixing is given by $\delta_{ij} \sim \epsilon^{\gamma_H/2} \xi$, where $\delta_{ij} \equiv \tilde{m}_{ij}^2 / \max(\tilde{m}_{ii}^2, \tilde{m}_{jj}^2)$. Assuming some degree of degeneracy between the first and second generations, negligible a terms, and an absence of CP violating phases (as in minimal gauge mediation), the strongest flavor bound is from $(\delta_{13}^d)_{\text{LL=RR}} \lesssim 5 \times 10^{-3}$ [31]. There are also potential constraints from $b \rightarrow s\gamma$ and $B_s \rightarrow \mu^+ \mu^-$ which are sensitive to model dependent choices, such as details of the chargino sector. Overall, we find no impediment to accommodating these constraints in our model.

Finally, let us briefly discuss the contributions to fine-tuning which result from our mechanism. The problem of naturalness is related to the question of curvature in the symmetry breaking direction—it is a one-dimensional problem for a Higgs field H as in the standard model with a potential

$$V = m_H^2 |H|^2 + \lambda |H|^4. \quad (28)$$

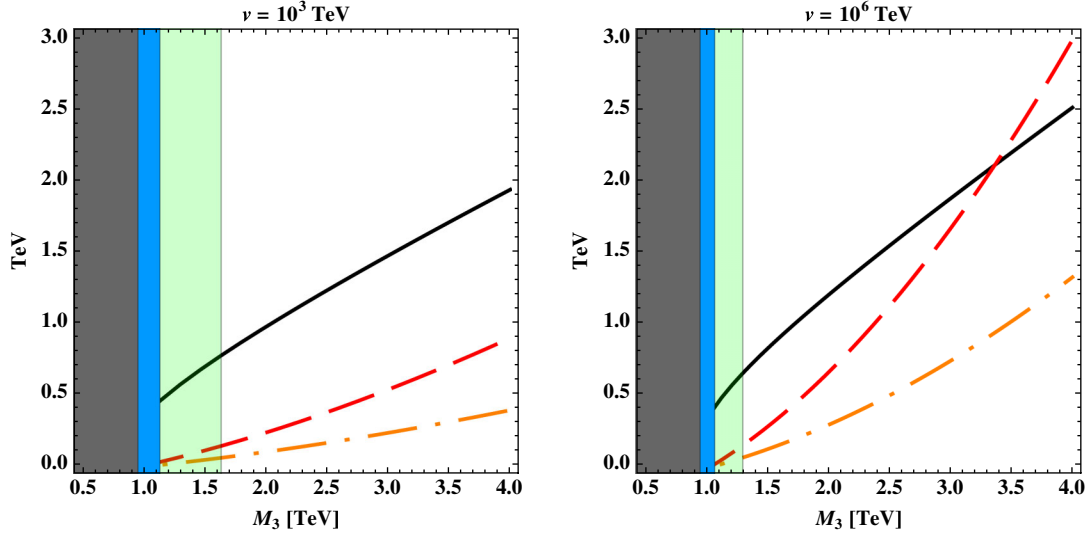


FIG. 3 (color online). Low energy spectrum for a model with unified gaugino masses, for $\nu = 10^3$ TeV (left) and $\nu = 10^6$ TeV (right). The curves represent $\tilde{m}_{Q_3} \approx \tilde{m}_{u_3} \approx \tilde{m}_{d_3}$ [black, solid], m_A (with $\tan\beta = 2$) [red, dashed], and m_A (with $\tan\beta = 10$) [orange, dotted-dashed]. (Only m_A has a strong dependence on $\tan\beta$). The first/second generation squark masses are at 5 TeV. The grey (left opaque region) is excluded due to tachyonic third generation squarks. The blue (right opaque region) is excluded by requiring radiative electroweak symmetry breaking. The green (translucent) region is excluded due to the LEP bound on the A mass. Both of these regions are plotted for the $\tan\beta = 2$ case. We fix $\Xi = 150$ GeV and for simplicity do not attempt the model dependent task of reproducing the Higgs mass for all points in this plot.

When $\langle H \rangle \neq 0$, the physical Higgs mass squared $m_h^2 = -2m_H^2$. One simple measure of fine-tuning, advocated in Refs. [11,32], is then

$$\Delta^{-1} \equiv -2 \frac{\delta m_H^2}{m_h^2} = -2 \frac{\tilde{m}_{H_u}^2}{m_h^2}, \quad (29)$$

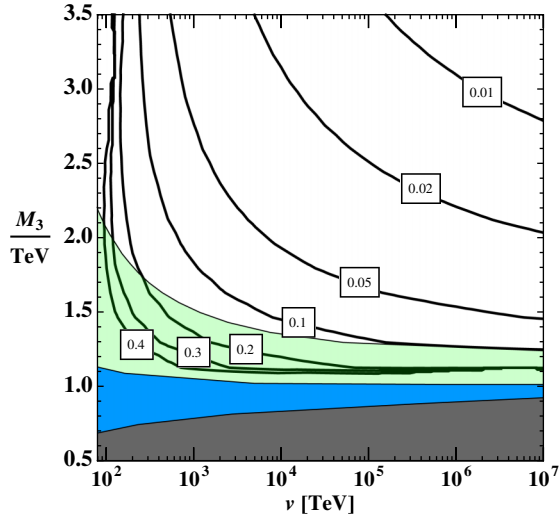


FIG. 4 (color online). We plot contours of the *fine-tuning* parameter Δ^{-1} in the M_3 versus ν plane. We make the same assumptions as in Fig. 3 with first/second generation squark masses at 5 TeV and $\tan\beta = 2$. The grey (lower opaque region) is excluded due to tachyonic third generation squarks, the blue (upper solid region) is excluded due to a lack of REWSB, and the green translucent region is excluded due to the LEP bound on the A mass.

where in the last equality, we are interested in the contribution to the Higgs soft mass in our model.

In Fig. 4, we have plotted contours of Δ^{-1} from our dynamics in the M_3 versus ν plane. The most important assumption from the point of view of fine-tuning is the gaugino mass spectrum. We have also plotted the region which is excluded due to tachyonic third generation squarks in solid grey, a lack of REWSB in solid blue, and the LEP bound on the A mass for $\tan\beta = 2$ in translucent green. We see that there is an allowed region with $\Delta^{-1} \approx \mathcal{O}(10\%)$ where $\nu \approx 10^2$ TeV and $M_3 \gtrsim 2.5$ TeV. We note that in a complete model which addresses the physical Higgs boson mass, there may be additional sources of fine-tuning.

IV. CONCLUSIONS AND FUTURE DIRECTIONS

In this work, we have presented a mechanism which acts as an attractor for the more minimal supersymmetric standard model and radiative electroweak symmetry breaking, while also generating the hierarchical structure of the quark Yukawa matrix. We have presented the simplest realization, which is accomplished by adding a new $SU(3)$ gauge group under which the third generation quarks are charged. The model flows to a strongly interacting fixed point where these quarks acquire order one negative anomalous dimensions, while the Higgs gets a positive anomalous dimension. The mechanism applies to generic supersymmetry breaking scenarios, as long as appropriate symmetries ensure that the combinations of masses in Eq. (10) are small. It also leads to a simple solution of

the μ problem. For concreteness, we analyzed the low energy phenomenology starting from unified gaugino masses, finding a natural supersymmetry spectrum with split families.

It would be interesting to build a fully realistic model based on this mechanism. The main points which need to be addressed are unification [which has been explicitly broken here by the extra matter charged under $SU(3)$] and the generation of a realistic physical Higgs mass. This motivates extending our approach by having two copies of the full SM gauge group, instead of just the $SU(3)$ group. One of the nodes will then become strongly coupled, leading to the properties analyzed here. In this context, the NMSSM can naturally become part of the strong dynamics, and unification is in principle possible.

The mechanism itself can also be improved in different directions. Here, we had to assume that certain approximate symmetries of the supersymmetry breaking sector were forbidding the combinations of soft masses given in Eq. (10). In particular, combinations proportional to $U(1)_Y$ can not be screened. This can be avoided if $U(1)_Y$ is embedded into a larger gauge group for the duration of the conformal regime. One possibility would be to weakly gauge the custodial $SU(2)$. In this case, the only combinations which are not sequestered are $\tilde{m}_{\Sigma}^2 - \tilde{m}_{\tilde{\Sigma}}^2$ and $\tilde{m}_{Q_3}^2 - \tilde{m}_{\bar{Q}_3}^2$, where $\bar{Q} = (\bar{u}, \bar{d})$ —both of these combinations can be suppressed by imposing a discrete symmetry. This can lead to a stronger attractor mechanism.

For models which realize this stronger attractor, there is a novel possibility of decoupling the first/second generation squarks beyond the bound of Ref. [25]. If $\tilde{m}_{1,2} \gg v$, at scales below $\tilde{m}_{1,2}$ there will be a quadratically divergent contribution to the stop masses at 2 loops and the Higgs

mass at 1 loop (which is proportional to small Yukawa couplings). If it is possible to construct a CFT which would be strong enough to suppress these quadratic divergences, the contribution to the mass from these effects will be schematically given by $y^2/(16\pi^2)v^2$ for the Higgs soft mass squared and $g_C^4/(16\pi^2)^2v^2$ for the stop soft mass squared. For $v \simeq 50$ TeV, these contributions are small enough to not destabilize our mechanism. Hence, the flavor problem could be completely decoupled in these models.

It may also be possible to find a microscopic realization where the CFT and the sector which breaks supersymmetry are part of the same dynamics. In this setup, the exit scale v would be related to the scale of supersymmetry breaking. This may be done at the level of the superpotential, or by destabilizing some of the flat directions of the CFT. Exploring a concrete supersymmetry breaking sector which minimizes the mass differences in Eq. (10) would also be an interesting avenue for future work.

If nature cares about naturalness, it is plausible that the dynamics between the weak scale and Planck scale could be highly nontrivial. We have demonstrated that coupling the supersymmetric standard model to a new strongly coupled conformal sector can give rise to the flavor hierarchies and the more minimal spectrum.

ACKNOWLEDGMENTS

We thank J. Wacker for collaboration at early stages of this work. We also thank I. Bah, S. El Hedri, T. Gherghetta, D. Green, H. Haber, S. Kachru, J. March-Russell, and M. Peskin for useful conversations and N. Craig for detailed comments on the manuscript. This work was supported by the US Department of Energy under Contract No. DE-AC02-76SF00515.

-
- [1] N. Arkani-Hamed, M. A. Luty, and J. Terning, *Phys. Rev. D* **58**, 015004 (1998); M. A. Luty and J. Terning, *Phys. Rev. D* **62**, 075006 (2000); S. Franco and S. Kachru, *Phys. Rev. D* **81**, 095020 (2010); N. Craig, R. Essig, S. Franco, S. Kachru, and G. Torroba, *Phys. Rev. D* **81**, 075015 (2010); S. R. Behbahani, N. Craig, and G. Torroba, *Phys. Rev. D* **83**, 015004 (2011).
 - [2] A. E. Nelson and M. J. Strassler, *J. High Energy Phys.* **09** (2000) 030.
 - [3] D. Poland and D. Simmons-Duffin, *J. High Energy Phys.* **05** (2010) 079; O. Aharony, L. Berdichevsky, M. Berkooz, Y. Hochberg, and D. Robles-Llana, *Phys. Rev. D* **81**, 085006 (2010); N. Craig, *arXiv:1004.4218*.
 - [4] T. Gherghetta and A. Pomarol, *Nucl. Phys. B* **586**, 141 (2000); M. Gabella, T. Gherghetta, and J. Giedt, *Phys. Rev. D* **76**, 055001 (2007); T. Gherghetta, B. von Harling, and N. Setzer, *J. High Energy Phys.* **07** (2011) 011; G. Larsen, Y. Nomura, and H. L. L. Roberts, *J. High Energy Phys.* **06** (2012) 032.
 - [5] N. Craig, D. Green, and A. Katz, *J. High Energy Phys.* **07** (2011) 045; R. Auzzi, A. Gaiotto, and S. B. Gudnason, *J. High Energy Phys.* **02** (2012) 069; N. Craig, S. Dimopoulos, and T. Gherghetta, *J. High Energy Phys.* **04** (2012) 116.
 - [6] D. E. Kaplan, F. Lepeintre, A. Masiero, A. E. Nelson, and A. Riotto, *Phys. Rev. D* **60**, 055003 (1999); D. E. Kaplan and G. D. Kribs, *Phys. Rev. D* **61**, 075011 (2000); M. Dine, R. G. Leigh, and A. Kagan, *Phys. Rev. D* **48**, 4269 (1993); P. Pouliot and N. Seiberg, *Phys. Lett. B* **318**, 169 (1993); A. Pomarol and D. Tommasini, *Nucl. Phys. B* **466**, 3 (1996); R. Barbieri, E. Bertuzzo, M. Farina, P. Lodone, and D. Pappadopulo, *J. High Energy Phys.* **08** (2010) 024; N. Craig, M. McCullough, and J. Thaler, *J. High Energy Phys.* **06** (2012) 046.

- [7] S. Dimopoulos and G. F. Giudice, *Phys. Lett. B* **357**, 573 (1995); A. G. Cohen, D. B. Kaplan, and A. E. Nelson, *Phys. Lett. B* **388**, 588 (1996).
- [8] K. Agashe and M. Graesser, *Phys. Rev. D* **59**, 015007 (1998); R. Barbieri, E. Bertuzzo, M. Farina, P. Lodone, and D. Zhuridov, *J. High Energy Phys.* **12** (2010) 070; **02** (2011) 044(E); H. Baer, V. Barger, P. Huang, and X. Tata, *J. High Energy Phys.* **05** (2012) 109; H. M. Lee, V. Sanz, and M. Trott, *J. High Energy Phys.* **05** (2010) 079.
- [9] G. Aad *et al.* (ATLAS Collaboration), *Phys. Lett. B* **710**, 67 (2012); S. Chatrchyan *et al.* (CMS Collaboration), *Phys. Rev. Lett.* **107**, 221804 (2011).
- [10] R. Essig, E. Izaguirre, J. Kaplan, and J. G. Wacker, *J. High Energy Phys.* **01** (2012) 074; C. Brust, A. Katz, S. Lawrence, and R. Sundrum, *J. High Energy Phys.* **03** (2012) 103; Y. Kats, P. Meade, M. Reece, and D. Shih, *J. High Energy Phys.* **02** (2012) 115.
- [11] M. Papucci, J. T. Ruderman, and A. Weiler, *J. High Energy Phys.* **09** (2012) 035.
- [12] D. E. Kaplan, G. D. Kribs, and M. Schmaltz, *Phys. Rev. D* **62**, 035010 (2000); Z. Chacko, M. A. Luty, A. E. Nelson, and E. Ponton, *J. High Energy Phys.* **01** (2000) 003.
- [13] N. Seiberg, *Nucl. Phys. B* **435**, 129 (1995); K. A. Intriligator and N. Seiberg, *Nucl. Phys. B, Proc. Suppl.* **45**, 1 (1996).
- [14] H. Fukushima, R. Kitano, and M. Yamaguchi, *J. High Energy Phys.* **01** (2011) 111.
- [15] C. T. Hill and E. H. Simmons, *Phys. Rep.* **381**, 235 (2003); **390**, 553(E) (2004).
- [16] M. Dine, J. L. Feng, and E. Silverstein, *Phys. Rev. D* **74**, 095012 (2006).
- [17] A. E. Nelson and M. J. Strassler, *J. High Energy Phys.* **07** (2002) 021.
- [18] T. Kobayashi and H. Terao, *J. High Energy Phys.* **07** (2004) 026.
- [19] M. Schmaltz and R. Sundrum, *J. High Energy Phys.* **11** (2006) 011.
- [20] K. A. Intriligator and B. Wecht, *Nucl. Phys. B* **667**, 183 (2003); E. Barnes, K. A. Intriligator, B. Wecht, and J. Wright, *Nucl. Phys. B* **702**, 131 (2004).
- [21] D. Green, Z. Komargodski, N. Seiberg, Y. Tachikawa, and B. Wecht, *J. High Energy Phys.* **06** (2010) 106.
- [22] G. F. Giudice and R. Rattazzi, *Phys. Rep.* **322**, 419 (1999).
- [23] R. Barbieri, L. J. Hall, S. Raby, and A. Romanino, *Nucl. Phys. B* **493**, 3 (1997).
- [24] N. Arkani-Hamed, G. F. Giudice, M. A. Luty, and R. Rattazzi, *Phys. Rev. D* **58**, 115005 (1998); T. Cohen, A. Hook, and B. Wecht, *Phys. Rev. D* **85**, 115004 (2012).
- [25] N. Arkani-Hamed and H. Murayama, *Phys. Rev. D* **56**, R6733 (1997).
- [26] G. Aad *et al.* (ATLAS Collaboration), *Phys. Lett. B* **710**, 49 (2012); S. Chatrchyan *et al.* (CMS Collaboration), *Phys. Lett. B* **710**, 26 (2012).
- [27] P. Batra, A. Delgado, D. E. Kaplan, and T. M. P. Tait, *J. High Energy Phys.* **02** (2004) 043; A. Maloney, A. Pierce, and J. G. Wacker, *J. High Energy Phys.* **06** (2006) 034.
- [28] S. Schael *et al.* (ALEPH and DELPHI and L3 and OPAL and LEP Working Group for Higgs Boson Searches Collaborations), *Eur. Phys. J. C* **47**, 547 (2006).
- [29] S. Chatrchyan *et al.* (CMS Collaboration), *Phys. Lett. B* **713**, 68 (2012).
- [30] A. Djouadi, *Phys. Rep.* **459**, 1 (2008).
- [31] M. Artuso, D. M. Asner, P. Ball, E. Baracchini, G. Bell, M. Beneke, J. Berryhill, A. Bevan *et al.*, *Eur. Phys. J. C* **57**, 309 (2008).
- [32] R. Kitano and Y. Nomura, *Phys. Lett. B* **631**, 58 (2005).