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# Suppressing supersymmetric flavor violations through quenched gaugino-flavor interactions

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## Abstract

Realizing that couplings related by supersymmetry (SUSY) can be disentangled when SUSY is broken, it is suggested that unwanted flavor and CP violating SUSY couplings may be suppressed via quenched gaugino-flavor interactions, which may be accomplished by power-law running of sfermion anomalous dimensions. A simple theoretical framework to accomplish this is exemplified, where a strongly coupled CFT is achieved after SUSY is softly broken. The defeated constraints are tallied. One key implication of the scenario is the expectation of enhanced top, bottom and tau production at the LHC, accompanied by large missing energy. Also, direct detection signals of dark matter may be more challenging to find than in conventional SUSY scenarios.

## I. INTRODUCTION

Low energy Supersymmetry (SUSY) is motivated as a solution to the weak-scale hierarchy problem. However, one of the challenges this theory presents is the potential introduction of new large flavor and CP violating contributions to observables that cannot accommodate significant new additions from new physics. There are many excellent ideas to solve this problem. In our work we explore yet another approach to solving the problem – the quenching of gaugino-sfermion-fermion interactions. There are several reasons to investigate this, as will become clearer throughout the discussion. It is a largely unexplored approach to solving the flavor problem. It does not add additional finetuning or naturalness problems compared to other conventional scenarios of supersymmetry that solve flavor only by fiat. And there may be interesting new avenues of supersymmetry breaking that lead to this approach.

Roughly speaking, approaches to supersymmetry have two extremes. One extreme is to think of supersymmetry as non-existent or broken dramatically, such that there is no SUSY in the low-scale spectrum that has much hope of being seen anytime soon by current colliders and experiments. The other extreme is to maintain that supersymmetry is broken very softly, and all superpartners are “nearby” in the spectrum with full coupling strengths, and are just out of reach of experiment but could be discovered very soon by slight increases in energy or luminosity at the LHC or by increased precision on low-energy flavor experiments.

However, nature may choose a middle way, where supersymmetry breaking dynamics is much more rich than the very simple minded effective approaches we have employed in most studies so far. Compactifications from higher dimensions, couplings to conformally symmetric sectors, or other so-far unrecognized dynamics may lead to a supersymmetric spectrum that has apparent unique patterns of couplings or hard-breaking interactions in the low-scale spectrum. Often times such couplings lead to a form of supersymmetry that is harder to see at high-energy colliders, but have the advantage of solving some outstanding problem, such as the flavor and CP violation problems of SUSY. It is this middle way that we propose to study, and our specific target is the elimination of gaugino-sfermion-fermion couplings. Many salient phenomenological features arise by invoking this idea. Furthermore there are reasonable theoretical approaches that may be able to accomplish exactly this needed pattern. That is our goal.

Any approach that preferentially quenches gaugino-flavor couplings necessarily involves the dynamics and/or transmission of supersymmetry breaking to the MSSM sector, since it is only through SUSY breaking that relations between couplings among SM particles can be disentangled from those involving superparticles. The conventional RG running only modifies the differences between couplings logarithmically [1–3]. So one must do something much more. One can imagine many different ways to split the couplings by much more than the standard logarithmic amount, but the approach that we used to illustrate the approach in this study is power-law running. Fast power-law running can be obtained if some superparticles are involved in a strongly coupled conformal field theory (sCFT) and obtain large anomalous dimensions (AD).

In this article, we study the possibility of employing fast power-law running that differentiates between particles and their superparticles as an illustrative means by which to quench gaugino-flavor interactions, thereby eradicating flavor and CP violation problems in supersymmetry. We show that these problems can be solved automatically with a reasonable energy range of power-law running.

In Sec. II, we give a brief review of the Nelson-Strassler mechanism in which power law running through RGE is used to generate a hierarchy among the Yukawa couplings. In Sec. III, we emphasize that the power law running does not require supersymmetry; nevertheless, we present a general framework of split-coupling SUSY where large suppressions on couplings such as squark-gluino-quark are achieved. Although supersymmetry is not necessary, it helps to control the calculability of the theory, especially the estimation of anomalous dimensions as well as mass spectra of superparticles if an approximate SUSY exists in the strongly coupled CFT. Thus, within that section, we consider a scenario where the MSSM is embedded in an  $N = 2$  SUSY theory and the strongly coupled CFT maintains an approximate  $N = 1$  SUSY orthogonal to that associated with the MSSM.

In Sec. IV, we go through a very detailed QFT description on the running of coupling constants and show how our theory is different from the NS mechanism. Especially, we demonstrate the differences between gauge couplings and their SUSY related couplings to explain where exactly the differences in RG running come from.

A particular subtlety one may be worried about in our scenario is whether the squarks can still be naturally around TeV in our setup since we break SUSY at the (higher) scale where sCFT begins. We address this point in Sec. V, where we argue that the soft mass

terms can be dramatically suppressed by RG running in the conformal region and are not necessarily large. This is especially relevant if there is an approximate SUSY existing in the strong CFT, which shows up in the  $N = 2$  embedding. In this case, soft masses of squarks are expected to be lower than the scale where sCFT ends, so they can be naturally small.

Neither constraints nor theory require the third generation gaugino-flavor couplings to be quenched. In the power-law scheme that we consider this means that the third generation squarks and slepton do not need to obtain large anomalous dimensions like the other two generations. When large anomalous dimensions are allowed for first two generations but not for the third, the couplings in different generations become split, but this is not the case for the squark masses themselves as occurs in other split family supersymmetry models [4, 5]. Suppressing the first two generation couplings alone can avoid flavor and CP violation constraints of supersymmetry, and is at least as beneficial from the naturalness point of view as other approaches. In Sec. VI, we discuss naturalness implications. We also compare split family SUSY and our scenario. We show that with the same level of fine tuning, the suppression on flavor/CP problems can be more efficient in our scenario.

In Sec. VII, we present the detailed calculations on how much quenching is needed in order to get around the constraints from flavor and CP violation measurements. In Sec. VIII, we also discuss interesting phenomenological implications including collider physics signatures as well as DM direct detections. At last in Sec. IX, we present a benchmark scenario and explicitly demonstrate the amount of suppression to be expected for various couplings.

## II. NELSON-STRASSLER MECHANISM

Our goal is to quench gaugino-flavor interactions, and one mechanism by which this may be possible is through coupling SUSY to a sCFT to initiate power-law running. Induced power-law running from a sCFT has been used in the past to address the flavor problem in related ways. Nelson-Strassler (NS) [6, 7] was one of the original approaches to fully exploit this feature. They targeted the fermion mass hierarchy problem as well as flavor-violating mixings in the squark/slepton soft SUSY breaking mass matrices and  $A$ -terms. Our illustrative approach targets different parts of the theory, namely the gaugino-sfermion-fermion interactions, but it utilizes similar CFT-coupling techniques as NS. Let us first

review NS.

Consider the SM gauge group  $S$  times another gauge group  $G$ . We assume that  $G$  runs to a sCFT at a high scale. Label SM particles as  $X$ , which are gauge singlets under  $G$ . Label exotic particles as  $Z$ , which are charged under  $G$  and may be charged also under  $S$ .

When  $G$  runs to its fixed point as a sCFT,  $\dim(X) > 1$  since  $X$  are  $G$ -singlet operators. One can model build to give large AD to the first two generations and small ones for the third generation and  $H_u/H_d$ . Thus, near the fixed point, the SM Yukawa coupling  $yX_1X_2X_3$  is irrelevant for the first two generations and almost marginal for the third generation. This induces power law suppression to the Yukawa couplings for the first two generations while keeping the third  $\mathcal{O}(1)$ .

### III. SPLIT COUPLING SUSY

The significant lesson so far is that large AD over a wide energy range, induced by a sCFT, can generate large suppression hierarchies from small initial differences. This is simply due to the anomalous dimension induced by sCFT, and it does not require SUSY. If the coupling of a SUSY theory to the CFT does not preserve SUSY, large SUSY breaking effects can occur as well that might ameliorate the flavor and CP violation problems.<sup>1</sup>

Let us consider the case where superparticles ( $\tilde{q}$ ) get large AD while SM particles ( $q$ ) do not. The fast power law running can introduce a large hierarchy between two couplings that were originally related by SUSY. In order to achieve such a scenario, we couple the MSSM to a sector  $G'$  differently for particles than superparticles. For example,  $q$  and  $\tilde{q}$  can couple differently to particles in  $G'$ . If  $\tilde{q}$  and  $G'$  form a sCFT at a particular energy scale,  $\tilde{q}$  can obtain large AD while  $q$  does not.

To be more specific, the Lagrangian can be formally written as

$$L \supset \kappa O_{\tilde{q}} O_{G'} + a \kappa O_q O_{\tilde{G}'} + \dots \quad (1)$$

Here ... includes the source of SUSY breaking. The couplings are defined schematically such that if  $a = 1$  supersymmetry is preserved; i.e., the quarks and squarks both couple to the  $G'$  sector in a manner demanded by supersymmetry invariance. In that case, SUSY is preserved

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<sup>1</sup> Here we emphasize that such SUSY breaking effects can be induced by soft SUSY breaking terms, as we will discuss in detail at the end of this section. SUSY is restored above a certain energy scale.

in the MSSM and supersymmetry breaking is achieved by some other, perhaps traditional, approach. This  $a = 1$  limit reduces to the NS scenario. In our proposal, we are suggesting a different limit,  $a \ll 1$ . This breaks supersymmetry and gives dramatically different AD to squarks vs. quarks.

To be more illustrative, let us consider the squark-gluino-quark coupling in the Lagrangian. In the MSSM with preserved SUSY, the coupling between quark and gluon is correlated with the coupling between squark-gluino-quark,

$$L \supset -iq^\dagger \bar{\sigma}^\mu (\partial_\mu - ig_3 A_\mu^a T^a) q - \sqrt{2} \lambda_3 (\tilde{q}^* T^a q \tilde{g}^a + c.c.) \quad (2)$$

i.e.,  $\lambda_3 = g_3$ . When SUSY is broken, these couplings become different. In ordinary scenarios, the difference is controlled by a logarithmic running. However if either the squark or gluino gains large AD from a sCFT,  $\tilde{q}^* T^a q \tilde{g}^a$  becomes an irrelevant operator. Thus  $\lambda_3$  enjoys a power law running. Similarly, the couplings between squarks/sleptons and electroweakinos can also be suppressed by power law running.

One subtlety we want to emphasize is that the gauge couplings between superparticles will not be power-law suppressed. This is guaranteed by gauge symmetry. One can confirm this principle by absorbing the gauge coupling into the kinetic term of the gauge boson. The interaction between the gauge boson and a particle is then directly extracted from the particle's kinetic term constructed from the gauge-coupling-less covariant derivative. The covariant derivative maintains itself without alteration upon canonically normalizing the particle's kinetic term.

The above are the generic aspects that can give rise to our scenario. For the rest of this section we give some additional comments about possible specific directions one may wish to pursue to build a specific and complete model. These comments are outside of the mainline of our present work, and the reader may wish to skip to the next section. However, we find the richness of realization possibilities encouraging for this scenario and wish to make a few remarks on them.

First, we imagine SUSY breaking only happens in a soft manner. More explicitly, SUSY is well preserved in the UV and the splitting of SUSY-related couplings is induced after SUSY breaking at a certain energy scale.

To begin with, there are a few examples where a non-SUSY CFT can be constructed. One is the  $\lambda\phi^4$  theory in  $(4 - \epsilon)$  dimension, which has been used as an example in conformal

sequestering models [8]. Another class of non-SUSY CFTs have been conjectured in [9], motivated by the AdS/CFT correspondence. As we argued previously, conformal symmetry is the key to introduce power-law running of particular couplings, while SUSY is not essential. With these non-SUSY CFT concepts in mind, one may even build models in which superparticles can be embedded after soft SUSY breaking and obtain large anomalous dimensions during conformal region.

In contrast to the non-SUSY approaches, one benefit of a SUSY CFT is that the anomalous dimension is associated with the  $R$ -charge of field. This makes the theory well under control computationally. Here we present a model where an approximate SUSY can be applied in order to control the theory while still generating the desired feature for Split Coupling SUSY.

Let us embed the MSSM into  $N = 2$  SUSY content. For example, a quark comes with a hypermultiplet as  $(q, \tilde{q}, q', \tilde{q}')$ . We label this sector as  $MSSM_2$ . There are two ways to split the hypermultiplet in terms of  $N = 1$  SUSY. One is  $(\{q, \tilde{q}\}, \{q', \tilde{q}'\})$  which is compatible with the SUSY generators of the MSSM. The other way is to group the hypermultiplet as  $(\{q, \tilde{q}'\}, \{q', \tilde{q}\})$ , which respects the SUSY generators in  $N = 2$  also, but whose  $N = 1$  embedding is orthogonal to that of the MSSM. The interactions in  $N = 2$  are highly constrained. We explicitly break  $N = 2$  SUSY in  $MSSM_2$  to  $N = 1$  SUSY by Yukawa couplings in the superpotential. Nevertheless, the  $N = 2$  SUSY particle contents still remain.

Now we introduce another sector labeled  $G$ . Assume  $G$  has also  $N = 2$  particle content and the interactions within this section preserve  $N = 2$  SUSY. Here we emphasize that the  $N = 2$  SUSY in sector  $G$  is not exact because at least gravity can mediate the  $N = 2$  SUSY breaking effects from sector  $MSSM_2$  to sector  $G$ . However, one generically expects such effects to be small since the running of dimensionless couplings is logarithmic. And the  $N = 2$  SUSY is still approximately preserved.<sup>2</sup>

Let us assume that there is a relevant operator introducing additional couplings between  $MSSM_2$  and  $G$  sectors. However, this relevant operator only preserves the  $N = 1$  SUSY

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<sup>2</sup> Our model is essentially a dimension deconstructed version of extra dimension model with  $N = 2$  SUSY being explicitly broken at fixed points of orbifolds.

which is orthogonal to that of the  $MSSM$ .

$$O_{rel} = M^n O_G O_{MSSM_2} \quad (3)$$

where  $n$  is a positive number which is determined by the dimensions of  $O_G$  and  $O_{MSSM_2}$ . For example,  $O_{MSSM_2}$  could be identified as a chiral supermultiplet consisting of  $\{q', \tilde{q}\}$ , which preserves the SUSY generator in  $N = 2$  SUSY orthogonal to the SUSY generators of the  $MSSM$ . Here we emphasize that SUSY is still a good symmetry at energy scales higher than  $M$  because the operator in Eq. (3) is a relevant operator. From the viewpoint of sector  $G$ , when the energy is below  $M$ , the approximate  $N = 2$  SUSY is dramatically broken by this relevant operator, but there is still an approximate  $N = 1$  SUSY which is orthogonal to the  $N = 1$  SUSY of the  $MSSM$ . If  $(G + O_{rel})$  runs to a strongly coupled CFT at a scale not far below  $M$ , the anomalous dimension of  $\tilde{q}$  can be calculated by its  $R$ -charge under the approximate  $N = 1$  SUSY in sector  $G$ .

One may worry that although  $\tilde{q}$  obtains a large anomalous dimension and its coupling to quark and gluino is power law suppressed,  $\tilde{q}'$  can remain untouched. Especially, if  $N = 2$  SUSY is preserved in the  $MSSM_2$  sector,  $\tilde{q}' - \psi_A - q$  has to appear, where  $\psi_A$  is a fermion paired with a gaugino in the  $N = 2$  gauge supermultiplet. However, as discussed above,  $N = 2$  SUSY is not preserved by interaction terms in  $MSSM_2$ . Such operators do not necessarily appear or do not have couplings comparable to the strong gauge coupling. Further,  $\tilde{q}'$  does not directly couple to our Higgs bosons. Therefore its mass can be large and does not generate large soft mass terms to Higgs doublets at 1 loop.

#### IV. QFT DESCRIPTION

Let us describe in more detail the underlying QFT picture for the suppression and splitting of couplings. We start with a simplified Lagrangian, writing explicitly only the terms that we are interested in:

$$\begin{aligned} L \supset & -\frac{1}{4}F^2 - |(\partial_\mu - igA_\mu)\tilde{q}_{L/R}|^2 - iq_{L/R}^\dagger \not{D} q_{L/R} \\ & - \sqrt{2}\lambda \tilde{q}_{L/R}^* q_{L/R} \tilde{\gamma} - \frac{y}{2}\phi q_L q_R + \dots \end{aligned} \quad (4)$$

Here we have kept the squark gauge coupling, the Yukawa coupling from the superpotential, and the interaction between squark-gaugino-quark. For simplicity, we consider a  $U(1)$  gauge

coupling with charge one. Generalizing to non-abelian gauge symmetry is straightforward.  $\lambda$  is equal to the gauge coupling  $g$  if SUSY is exact.

Now let us integrate out an energy shell and renormalize the Lagrangian,

$$\begin{aligned}
L \supset & -\frac{1}{4}Z_A F^2 - Z_{\tilde{q}_{L/R}} |(\partial_\mu - iZ_g g A_\mu) \tilde{q}_{L/R}|^2 \\
& -iZ_{q_{L/R}} q_{L/R}^\dagger \not{D} q_{L/R} - \sqrt{2}\lambda Z_\lambda \tilde{q}_{L/R}^* q_{L/R} \tilde{\gamma} \\
& -\frac{y}{2} Z_y \phi q_L q_R + \dots
\end{aligned} \tag{5}$$

Here  $Z_A$ ,  $Z_{\tilde{q}_{L/R}}$  and  $Z_{q_{L/R}}$  come from wavefunction renormalization.  $Z_g$ ,  $Z_\lambda$  and  $Z_y$  are from 1-PI renormalization of the interaction vertices. If a field obtains a large AD, its wavefunction renormalization  $Z_i$  is power-law enhanced,  $Z_i \sim (\frac{1}{\epsilon_i})^2 \sim (M_{c,b}/M_{c,e})^{2\gamma_i}$ . Here  $M_{c,b}(M_{c,e})$  is the beginning (ending) energy scales of the conformal regime.  $\gamma_i$  is the anomalous conformal dimension.

The 1-PI vertex renormalization factors depend on the details of how strongly the  $G'$  sector couples; it is assumed to be  $O(1)$  or larger. After canonically normalization,

$$\begin{aligned}
L \supset & -\frac{1}{4}F^2 - |D_\mu \tilde{q}_{L/R}|^2 - i q_{L/R}^\dagger \not{D} q_{L/R} \\
& -\sqrt{2}\lambda (Z_\lambda Z_{\tilde{q}_{L/R}}^{-1/2} Z_{q_{L/R}}^{-1/2} Z_{\tilde{\gamma}}^{-1/2}) \tilde{q}_{L/R}^* q_{L/R} \tilde{\gamma} \\
& -\frac{y}{2} Z_y Z_{q_L}^{-1/2} Z_{q_R}^{-1/2} Z_\phi^{-1/2} \phi q_L q_R + \dots
\end{aligned} \tag{6}$$

For gauge coupling terms, the factors of wavefunction renormalization from matter fields cancel precisely the 1-PI interaction vertex correction in the covariant derivative terms. This is required by gauge symmetry. If SUSY is preserved, the cancelation also happens in the squark-quark-gaugino vertex, i.e.  $Z_\lambda = Z_g \sim (\frac{1}{\epsilon_i})^2$ , which guarantees that the coupling constant  $\lambda$  is the same as the gauge coupling  $g$  at all scales.

In contrast, within the NS scenario, cancelation does not occur for the Yukawa vertex. Appealing to the non-renormalization theorem,  $Z_y = 1$  is fixed. Nevertheless, wavefunction renormalization introduces a power-law suppression on the Yukawa couplings as  $y_q \sim \epsilon_{q_L} \epsilon_{q_R}$ .

Now let us turn to our scenario where the squark and quark couple differently to the sector  $G'$ . If squark+ $G'$  runs to a sCFT at a particular scale, the squark may obtain a large AD, i.e.  $Z_{\tilde{q}} \sim (1/\epsilon_{\tilde{q}})^2 \gg 1$ . Since SUSY is broken,  $Z_\lambda$  no longer has a rigid supersymmetric relation with  $Z_g$ . This disconnect is maximal if  $G'$  does not directly couple to quarks and gauginos, but does couple to squarks. This results in  $Z_\lambda$ ,  $Z_q$  and  $Z_{\tilde{\gamma}}$  expected to be  $\mathcal{O}(1)$ , whereas

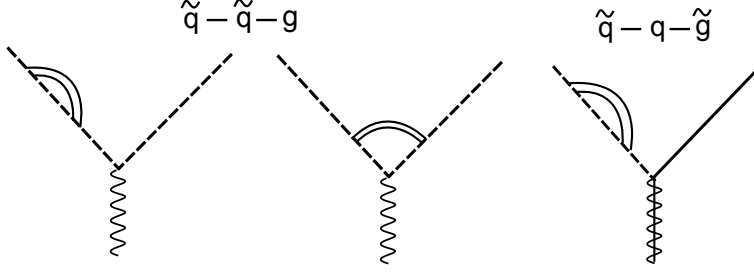


FIG. 1: Here we show a few examples of Feynman diagrams where a strongly coupled CFT may contribute at 1-loop. Double lines indicate large renormalization from the strongly coupled CFT. If sCFT only couples to squarks, both the squark propagator and its gauge coupling vertex receive large corrections, as shown in the first two diagrams. However, there is no large 1-PI vertex correction to  $\tilde{q} - q - \tilde{g}$  vertex as shown in the third diagram.

$Z_{\tilde{q}_{L/R}}^{-1/2} \sim \epsilon_{\tilde{q}_{L/R}}$ . Thus, the only significant suppression factors in our theory arise from wave function renormalization factors on squarks and sleptons, and possibly higgsinos also, as will be discussed later. This leads to quenching of the gaugino-flavor interactions and possibly also the higgsino-flavor interactions. The mismatch between couplings of quarks/leptons vs. squarks/sleptons is schematically illustrated in Fig. 1.

The considerations outlined in this section enable us to define various options of spectra for the masses of the minimal supersymmetric particles. Various benchmark possibilities will be discussed in detail later in sec. IX.

## V. SOFT SUSY BREAKING TERMS

It is also important to understand how the soft mass terms run while in the conformal region. Naively, one would expect that the soft mass terms are large in our scenario because the SUSY relations among couplings are broken at the beginning of the conformal region, i.e.  $E_{c,b}$ . Especially, one would generically expect that the squark soft mass square is quadratically sensitive to  $E_{c,b}$ . However, the running of soft terms is quite subtle in the conformal region and highly depends on how the  $MSSM$  couples to the sCFT. Such running can be estimated explicitly if the CFT has an approximate supersymmetry. The details are studied in [7].

Let us take our example where the  $MSSM$  is embedded into an  $N = 2$  SUSY framework,

as discussed in Sec. III. In such a model, from the viewpoint of sector  $G$ , an  $N' = 1$  SUSY, orthogonal to that of MSSM, is still an approximate symmetry due to strong couplings, and the running of soft terms can still be estimated in the context of SUSY. In [7], the authors show that some particular linear combinations of soft masses can be suppressed during RG flow in the conformal region, while others are not, depending on the residual global symmetries when coupling the two sectors.

The argument behind this claim is the following. The soft SUSY breaking terms can be combined with coupling constants into a supermultiplet. For example, a general form of superpotential can be written as

$$W = \sum_s y_s \prod_i (\phi_i)^{p_s^i}. \quad (7)$$

One can form a vector supermultiplet whose lowest component is a function of  $y_s$ , whose  $\theta^2$  component is related to the A-term and whose  $\theta^4$  component is a linear combination of soft mass terms of  $\phi_i$ 's,

$$m_s^2 \equiv \sum_i p_s^i m_i^2. \quad (8)$$

Due to supersymmetry, the beta functions of these three terms are related to each other. It can be shown that [7], by requiring the conformal fixed point to be stable, the linear combination of the soft mass terms will be suppressed dramatically through RG running. A similar argument can be applied to gauge couplings as well in order to prove that certain linear combinations, related to the gauge charges of the fields, of soft mass terms can be naturally small after RG running,

$$m_\lambda^2 \equiv \frac{T_G |M|^2 - \text{Tr } m^2}{-2\pi + \alpha T_G}. \quad (9)$$

Here  $M$  is the gaugino mass which has no quadratic sensitivity to  $E_{c,b}$ . On the other hand, if there is an exact global symmetry preserved in the Lagrangian when coupling to the sCFT, there is a particular linear combination of soft mass squares unaffected by RG flow. A more detailed discussion on the running of soft terms can be found in Appendix B of [7].

The point of the discussion above is to show that the soft mass terms of squarks in our scenarios are not necessarily large, i.e. not quadratically sensitive to  $E_{c,b}$ , thanks to the approximate  $N' = 1$  SUSY preserved in the sCFT. Meanwhile, flowing into the IR after the conformal region, the soft mass terms run as normal. Since the SUSY-related couplings do

not match anymore, one expects the soft mass squares to be at least quadratically sensitive to the end of conformal region, i.e.  $E_{c,e}$ . Thus in order to have squark at  $O(\text{TeV})$ , we expect  $E_{c,e}^2$  to be at most 1 loop factor larger than squark soft mass squared, i.e.  $E_{c,e} \sim O(10) \text{ TeV}$ .

To make our discussion more clear, as well as to illustrate that a large anomalous dimension with suppressed soft mass term can actually be achieved, let us present a simple toy model.<sup>3</sup> We label the  $N' = 1$  supermultiplet which contains MSSM squark as  $\Phi_1^a$ , where  $a$  is the  $SU(3)_c$  index. We further introduce two  $N' = 1$  chiral supermultiplets  $\Phi_2^a$  and  $\Phi_3$ . Both fields transform as fundamental representation of a strongly coupled  $SU(2)_s$  gauge group. Further,  $\Phi_2$  transforms as anti-fundamental representation of  $SU(3)_c$ . At last, we introduce a gauge singlet  $\Phi_4$ . A superpotential is introduced to couple  $\Phi_i$  as

$$W = \lambda_1 \Phi_1^a \Phi_2^a \Phi_3 + \lambda_2 \Phi_3^2 \Phi_4 + \lambda_3 \Phi_4^3. \quad (10)$$

Note that the gauge coupling of  $SU(3)_c$  is small compared to the strong  $SU(2)_s$ , thus  $SU(3)_c$  index should be treated as a flavor symmetry index. From superpotential, the RG running of  $\lambda_i$  implies three relations among soft mass terms according to Eq. (8),

$$\begin{aligned} m_1^{a2} + m_2^{a2} + m_3^2 &\rightarrow 0 \\ 2m_3^2 + m_4^2 &\rightarrow 0 \\ 3m_4^2 &\rightarrow 0 \end{aligned} \quad (11)$$

at  $E_{c,e}$ , where the superscript  $a$  is the  $SU(3)_c$  index (not part of the exponent). At the meanwhile, from the running of  $SU(2)_s$  gauge coupling, according to Eq. (9), we have

$$3m_2^{a2} + m_3^2 \rightarrow 0 \quad (12)$$

Now we see that the soft mass term of the squark in  $\Phi_1$  is forced smaller during the RG running, which indicates that the squark mass is not quadratically sensitive to  $E_{c,b}$ .

Now let us consider the strongly coupled  $SU(2)_s$ . We have  $N_c = 2$  and  $N_f = 4$  which indicates that it confines.<sup>4</sup> One can further use Seiberg duality to describe the dual theory.

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<sup>3</sup> We refer readers who are interested in a more detailed model building to [6]. The authors in that paper presented multiple ways to achieve a realistic model with detailed MSSM field embedding in an analogous context that can be utilized for our new application here.

<sup>4</sup> Although this theory never reaches exact conformal fixed point, it is well-enough approximately conformal in the IR, aka below  $E_{c,b}$ , for the anomalous dimension to be large to drive our wanted dynamics. One can construct another (more complicated) theory that is strictly conformal in the IR, if desired, by adding another two states in fundamental representation of  $SU(2)_s$ .

Especially, the binary  $\Phi_2\Phi_3$  becomes a mesonic field in the magnetic theory. The R-charge of the mesonic field is  $2(N_f - N_c)/N_f = 1$  which implies its conformal dimension as  $3/2$ . Thus  $\Phi_1$  obtains an anomalous dimension as large as  $1/2$  during the conformal region.

In summary, we show that it is possible to realize a scenario where squarks are as light as TeV while their couplings dramatically deviate from their SUSY related ones. Thus the SUSY flavor/CP problems are resolved through the coupling splitting but not the mass splitting, and our scenario is very different from the split SUSY scenario.

## VI. THE FIRST TWO GENERATIONS VERSUS THE THIRD

If squarks get large anomalous dimensions during the conformal region, all couplings involving squarks, such as squark-gluino-quark or squark-higgs quatic couplings, dramatically deviate from its SUSY required values. Thus the cancelations among contributions to higgs soft mass terms fail, and fine tuning becomes a concern. In particular, the mis-match on coupling constants starts at the scale  $E_{c,b}$ . The third generation quarks, especially top quark, have relatively large Yukawa couplings to the higgs boson. If stop gets large anomalous dimension, the corrections to higgs soft mass terms come in at 1-loop level. Given the current status on stop and gluino searches at the LHC,  $O(1) \sim O(10)\%$  fine tuning is already necessary in the MSSM [10]. If we take that as a benchmark point,  $M_{c,b}$  can only be as large as  $O(1) \sim O(10)$  TeV. This is phenomenologically unacceptable due to unseen exotic colored resonances or unseen large deviations of precision electroweak measurements.

Fortunately, SUSY flavor/CP problems are mainly induced by the mixing among the first two generation squarks. Thus one attractive scenario is to only give large anomalous dimensions to the first two generation squarks, while keep the third generation weakly coupled to the strong sector <sup>5</sup>. The leading contributions to higgs soft mass terms appear at 2-loop level through electroweak gauge couplings. Again, by requiring less than  $O(10)\%$  fine tuning,  $M_{c,b}$  is required to be at most  $O(100)$  TeV. This result is comparable to that from the naturalness concerns in split family SUSY because the scale where coupling constants are different by  $O(1)$  is analogous to the scale where such super-particles are removed in split family SUSY models.

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<sup>5</sup> We envision giving the same anomalous dimension to all superparticles within each generation in order to avoid 1-loop quadratic sensitivities from  $D$ -terms.

It is interesting to compare the efficiencies to resolve the flavor/CP problems in our mechanism with split family SUSY where the first two generations of squarks are raised to  $O(10) \sim O(100)$  TeV. In the language of effective operators, squark masses are identified as the suppression scale of the higher dimension operators responsible for flavor changing or CP violating processes. For example, a dimension 6 operator is proportional to  $\frac{1}{m_{\tilde{q}}^2}$ . In split family SUSY, that is identified to be  $O(10) \sim O(100)$  TeV. On the other hand, in our scheme, the suppression scale is much lower, e.g.  $m_{\tilde{q}} \sim O(1)$  TeV. However, as we will show in the next section, the diagrams which generate such higher dimension operators may have a high power dependence on coupling constants. For example, a box diagram generating  $\Delta F$  by two units is proportional to the fourth power of the coupling constants. Thus the efficiency on suppressing the flavor changing processes can be much higher in our scheme than simply raising superparticles' masses in split family SUSY.

## VII. REMEDYING THE SUSY FLAVOR/CP PROBLEM

The primary motivation of our work is to suppress the flavor and CP problems of supersymmetry via quenched gaugino-flavor couplings which can be accomplished, for example, by power-law suppressions of squark/slepton couplings. We first review the constraints and then apply those constraints to our scenario and show that reasonable parameters lead to safe phenomenology.

Flavor and CP measurements put stringent constraints on the parameter space in supersymmetry [11]. All these constraints can be characterized by flavor off-diagonal soft SUSY breaking mass terms.

Let us first focus on hadronic systems. There are two classes of processes, characterized by the change of flavor number, i.e.  $\Delta F = 1$  or  $2$ .

$\Delta F = 2$  processes can be described by dimension 6 operators, e.g.

$$O = \frac{1}{\Lambda^2} (\bar{d}_L \gamma^\mu s_L) (\bar{d}_L \gamma_\mu s_L). \quad (13)$$

These are induced by a box diagram with four squark-gluino-quark vertices. Since the flavor number is changed by 2, we need at least two insertions of flavor off-diagonal soft mass elements. Thus the suppression scale of such operators scales as

$$\frac{1}{\Lambda^2} \sim \lambda^4 \frac{(\delta_{ij}^a)^2_{AB}}{m_{SUSY}^2} \quad (14)$$

$(\delta_{ij}^a)_{AB} = \frac{(m_{ij}^a)^2_{AB}}{m_{SUSY}^2}$ , assuming all flavor-diagonal squark masses and gluino mass equal to  $m_{SUSY}$  for simplicity.  $\{A, B\}$  label left/right squarks,  $\{a\}$  indicates up/down type,  $\{i, j\}$  are generation indices. We also explicitly write the dependence on squark-gluino-quark couplings ( $\lambda$ ), without specifying flavors – flavor violations are all embedded in the  $\delta_{ij}$  factors. If  $\lambda$  is much smaller than the gauge coupling  $g_3$ , which power-law running can accomplish ( $\lambda \sim \epsilon_{\bar{q}} g_3$ ), the constraints from flavor/CP can be weakened.

Assuming flavor universality, the strongest constraints on  $\Delta F = 2$  processes are from  $K - \bar{K}$  mixing. This system also manifests a CP-violating phenomenon characterized by the parameter  $\varepsilon$ . We suppress all indices of  $\delta$  assuming all  $(\delta_{ij}^a)_{AB}$  are comparable and complex. From  $K - \bar{K}$  system measurements, constraints become

$$\begin{aligned} \left(\frac{\lambda}{g_3}\right)^2 |\text{Re}(\delta)| &< 10^{-3} \left(\frac{m_{SUSY}}{500 \text{ GeV}}\right) \\ \left(\frac{\lambda}{g_3}\right)^2 |\text{Im}(\delta)| &< 10^{-4} \left(\frac{m_{SUSY}}{500 \text{ GeV}}\right). \end{aligned} \quad (15)$$

If  $\frac{\lambda}{g_3}$  is smaller than  $10^{-2}$ , which is easily accomplished by our power-law running scenario, one can have  $\mathcal{O}(1)$  flavor mixing with TeV scale squarks.

For  $\Delta F = 1$  processes, both box diagrams and penguin diagrams contribute. For illustration, we show one operator for each kind:

$$O_{box} = \frac{(\bar{d}_L \gamma^\mu s_L)(\bar{q}_L \gamma_\mu q_L)}{\Lambda_{box}^2}; O_{pen} = \frac{\bar{d}_L \sigma^{\mu\nu} s_R F_{\mu\nu}}{\Lambda_{pen}}. \quad (16)$$

Since  $\Delta F = 1$ , one only needs one flavor-changing mass insertion in the loop. Thus we have

$$\frac{1}{\Lambda_{box}^2} \sim \lambda^4 \frac{(\delta_{ij}^a)_{AB}}{m_{SUSY}^2}. \quad (17)$$

Penguin diagrams have subtlety on chirality. Depending on whether the chirality is changed in the soft mass insertion, the operator effectively is either dimension 5 or 6:

$$\frac{1}{\Lambda_{pen}} \sim c_1 \lambda^2 \frac{(\delta_{ij}^a)_{AA} m_q}{m_{SUSY}^2} + c_2 \lambda^2 \frac{(\delta_{ij}^a)_{LR}}{m_{SUSY}}. \quad (18)$$

The strongest constraints on  $\Delta F = 1$  processes come from a CP-violating measurement in Kaon system,  $\varepsilon'/\varepsilon$ . If the chirality is not changed by a soft mass insertion, i.e. by  $(\delta_{ij}^a)_{AA}$ , the box and penguin contributions tend to cancel with each other when  $\lambda = g_3$ . However the box diagram has a different  $\lambda$  dependence from that of the penguin diagram, and when  $\lambda \ll g_3$  the penguin diagrams always dominate.

Assume all  $\delta$ 's are comparable complex numbers, we have

$$\begin{aligned}\left(\frac{\lambda}{g_3}\right)^2 |\text{Im}(\delta)_{AA}| &< 10^{-1} \left(\frac{m_{SUSY}}{500 \text{ GeV}}\right)^2 \\ \left(\frac{\lambda}{g_3}\right)^2 |\text{Im}(\delta)_{LR}| &< 10^{-5} \left(\frac{m_{SUSY}}{500 \text{ GeV}}\right).\end{aligned}\quad (19)$$

When  $m_{SUSY}$  is  $\mathcal{O}(\text{TeV})$ , if  $\frac{\lambda}{g_3} < 3 \times 10^{-3}$ ,  $\delta$  can be  $\mathcal{O}(1)$ .

Flavor-changing lepton decays,  $\ell_i \rightarrow \ell_j + \gamma$ , especially muon decay, impose constraints on the leptonic sector. Assuming  $m_{\tilde{l}} \sim m_{\tilde{\gamma}} \sim m_{SUSY}$ ,

$$\begin{aligned}\left(\frac{\lambda'}{e}\right)^2 |\text{Im}(\delta)_{AA}| &< 10^{-2} \left(\frac{m_{SUSY}}{100 \text{ GeV}}\right)^2 \\ \left(\frac{\lambda'}{e}\right)^2 |\text{Im}(\delta)_{LR}| &< 10^{-6} \left(\frac{m_{SUSY}}{100 \text{ GeV}}\right).\end{aligned}\quad (20)$$

$\lambda'$  is the slepton-photino-lepton coupling. If  $\frac{\lambda'}{e} < 10^{-3}$ , one can get around the muon decay constraint with  $\mathcal{O}(1)$  mixing in the slepton soft mass matrix, even if the slepton and photino are as light as 100 GeV.

Finally, the flavor diagonal soft mass matrix, i.e.  $(\delta_{ii})_{LR}$ , can be strongly constrained by radiative mass corrections and electric dipole moments.

The radiative mass corrections scale as

$$\begin{aligned}\Delta m_q &\sim \lambda^2 m_{SUSY} \text{Re}(\delta_{ii}^q)_{LR} \\ \Delta m_l &\sim \lambda'^2 m_{SUSY} \text{Re}(\delta_{ii}^l)_{LR}.\end{aligned}\quad (21)$$

$m_{SUSY}$  refers to the gluino and photino masses for the quark and lepton mass corrections respectively. The strongest constraint comes from the first generation. Requiring the radiative masses to not exceed the quark/lepton mass, one gets

$$\begin{aligned}\left(\frac{\lambda}{g_3}\right)^2 \text{Re}(\delta_{11}^q)_{LR} &< 2 \times 10^{-3} \left(\frac{500 \text{ GeV}}{m_{SUSY}}\right) \\ \left(\frac{\lambda'}{e}\right)^2 \text{Re}(\delta_{11}^l)_{LR} &< 8 \times 10^{-3} \left(\frac{100 \text{ GeV}}{m_{SUSY}}\right).\end{aligned}\quad (22)$$

The electric dipole moments of the neutron and electron scale as

$$\frac{d_q}{e} \sim \lambda^2 \frac{\text{Im}(\delta_{11}^q)_{LR}}{m_{SUSY}}, \quad \frac{d_e}{e} \sim \lambda'^2 \frac{\text{Im}(\delta_{11}^l)_{LR}}{m_{SUSY}}. \quad (23)$$

The constraints are,

$$\begin{aligned} \left(\frac{\lambda}{g_3}\right)^2 \text{Im}(\delta_{11}^q)_{LR} &< 3 \times 10^{-6} \left(\frac{m_{SUSY}}{500 \text{ GeV}}\right) \\ \left(\frac{\lambda'}{e}\right)^2 \text{Im}(\delta_{11}^l)_{LR} &< 4 \times 10^{-7} \left(\frac{m_{SUSY}}{100 \text{ GeV}}\right). \end{aligned} \quad (24)$$

If  $\delta$ 's are  $\mathcal{O}(1)$  complex numbers, one needs  $(\frac{\lambda}{g_3})$  or  $(\frac{\lambda'}{e})$  smaller than  $\mathcal{O}(10^{-3})$  for light squarks and sleptons, which are readable accessible with power-law running suppressions.

## VIII. SIGNATURES IN PHENOMENOLOGY

### A. Collider signatures

One interesting phenomenological consequence is its unique collider signatures. First, the production of superparticles are not dramatically suppressed because all squarks and the gluino have unsuppressed gauge couplings. The production channels of squark/gluino are mainly through gluon fusion and  $q\bar{q} \rightarrow g^*$ , whereas quark-gluino associated production is suppressed.

We assume the LSP is a neutralino which has suppressed couplings to the first two generations of squarks. We also assume all squarks have comparable masses for simplicity. When the gluino is heavier, the squark directly decays to neutralino. Interestingly the dominant decay channel would be through mixing to the third generation squarks,  $\tilde{u} \rightarrow \tilde{t} \rightarrow t\tilde{\chi}^0$ , where  $\tilde{\chi}^0$  is a neutralino. Thus the signature would be mainly two third-generation quarks plus missing energy (MET). If two gluinos are produced, the signature would be four bottom/stop quarks plus MET.

If squarks are heavier, gluino production is dominant. They will decay through off-shell squarks that will in turn produce four third-generation quarks plus MET. If squarks are produced, there can be six third-generation quarks plus MET. However, it is possible that the stop/sbottom-gluino-top/bottom vertices receive a stronger suppression than those with neutralino, squarks then would decay to neutralino without passing through a gluino. The detailed signatures depend on the UV model. Nevertheless, the main theme is clear: the preponderance of third-generation quarks and leptons accompanied by missing energy.

## B. Dark matter and direct detection

In the MSSM, the lightest neutralino is a good candidate for DM. At tree level, they can interact with nucleons through  $t$ -channel higgs exchange,  $t$ -channel  $Z$  exchange and  $s$ -channel squark exchange, which potentially allows for their discovery in suitable laboratory detectors.

The mixture among the neutralino interaction eigenstates to form mass eigenstates originates from their couplings to higgsinos, i.e.  $h^\dagger \tilde{h} \tilde{b}$  and  $h^\dagger \tilde{h} \tilde{w}$ . If the mixing is small, a neutralino barely couples to the Higgs and  $Z$  bosons. If either higgsinos or the gauginos (bino/wino) obtain large AD, the mixing can be highly suppressed. Thus it is quite generic to expect the lightest neutralino to be almost a pure state. If the LSP is a higgsino and all neutralinos have mass  $\mathcal{O}(\text{TeV})$ , the suppression factor needs to be no stronger than  $\mathcal{O}(10^{-2})$ , in order for the higgsino to not be a Dirac fermion from the DM detection point of view. However, a pure state like this creates experimental detection challenges. In particular, the mixing is small even when the bino, wino and higgsino have comparable masses, and so DM detection via Higgs or  $Z$  exchange is not effective method to problem the theory.

$s$ -channel squark exchange may also induce scattering between DM and nucleons. However, a strongly coupled CFT can suppress such couplings efficiently. The conclusion remains that the strong constraints on SUSY from lack of direct detection of DM do not apply here.

## IX. COUPLING HIERARCHIES AND BENCHMARK SCENARIOS

Let us briefly summarize the physics effects introduced by the split coupling scenario. For simplicity, we assume the strongly coupled sector does not introduce large 1-PI vertex renormalization except for gauge vertices. After canonical normalization, power-law suppression factors to non-gauge vertices arise from the external leg corrections, which include squarks, sleptons and Higgsinos. We identify these factors  $\epsilon_{\tilde{q}}$ ,  $\epsilon_{\tilde{\ell}}$ , and  $\epsilon_{\tilde{h}}$ , all of which are  $\epsilon \ll 1$ . However, we will assume that third generation squarks and sleptons do not receive large power-law suppression factors,  $\epsilon_{\tilde{t}, \tilde{b}, \tilde{\tau}} \sim 1$ . This choice is not necessary for the viability and unique attractiveness of the theory; however, it is allowed experimentally due to the relatively weak flavor bounds on third generation processes and does not exacerbate the naturalness problem.

Here we write a few examples of how the vertices are altered:

$$h^2\tilde{\ell}^2 \rightarrow (\epsilon_{\tilde{\ell}}^2)h^2\tilde{\ell}^2; \quad \tilde{q}\tilde{g}q \rightarrow (\epsilon_{\tilde{q}})\tilde{q}\tilde{g}q; \quad h\tilde{h}\tilde{w} \rightarrow (\epsilon_{\tilde{h}})h\tilde{h}\tilde{w}, \quad (25)$$

where we have ignored Lorentz indices, etc. Other interactions are not suppressed in our approach, most notably  $h\bar{f}f$  (where  $f = q, \ell, \nu$ ) and all gauge interactions arising from covariant derivative interactions, such as  $A_\mu\bar{f}\gamma^\mu f$ ,  $A_\mu\tilde{f}^*\partial^\mu\tilde{f}$ , etc. Furthermore, as we have discussed earlier, one is likely to be able to build the soft supersymmetric masses to obtain any value wanted. For this reason, we view them as unconstrained parameters from the theory point of view, which enables one freely to pursue spectra that minimize the finetuning of electroweak symmetry breaking, for example, and more importantly suggests that all mass scales and hierarchies should be considered when constructing benchmark models and searching for experimental signatures of the scenario.

## X. DISCUSSION

In this article, we have discussed the possibility of quenching gaugino-flavor and higgsino-flavor interactions for the purposes of remedying the flavor problem of SUSY. There are potentially many different ways to achieve this phenomenological aim. For illustrative purposes we have considered the prospect of coupling the MSSM to a SUSY-breaking sector  $G'$  which flows to a sCFT at a particular energy scale. The SUSY flavor/CP problems are naturally solved by quenching the gaugino-flavor couplings. Such split coupling models are analogues to split SUSY or split family models, depending on whether the third generation squarks also get large anomalous dimensions. However, instead of having a split mass spectrum, it is the SUSY-related couplings that develop large separations. Neutralino DM is likely to be a pure state, whose direct detection signals are suppressed.

If only the first two generations of squarks get large anomalous dimensions, the collider signatures are similar to those in the natural SUSY approach, where third generation quarks and leptons dominate the final state, but with different signal rates. As for the SM fermion mass hierarchy, it can be solved separately by applying the NS-mechanism. In that case, the superparticle couplings would be further suppressed. Finally, it is also interesting to consider the suppression of  $R$ -Parity violating vertices through a similar mechanism. We leave the details of this possibility for future study.

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