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SUSY, the Third Generation and the LHC

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Abstract: We develop a bottom-up approach to studying SUSY with light stops and sbottoms, but with other squarks and sleptons heavy and beyond reach of the LHC. We discuss the range of squark, gaugino and Higgsino masses for which the electroweak scale is radiatively stable over the "little hierarchy" below 10 TeV. We review and expand on indirect constraints on this scenario, in particular from flavor and CP tests. We emphasize that in this context, R-parity violation is very well motivated. The phenomenological differences between Majorana and Dirac gauginos are also discussed. Finally, we focus on the light subsystem of stops, sbottom and neutralino with R-parity, in order to probe the current collider bounds. We find that 1/fb LHC bounds are mild and large parts of the motivated parameter space remain open, while the 10/fb data can be much more decisive.

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Contents

1 Introduction

Supersymmetry (SUSY) remains a strong contender for the mechanism underlying electroweak stability. If one puts great stock in a particular high-energy SUSY model, the model couplings can be run down and matched onto an LHC-energy effective Lagrangian, $\mathcal{L}_{\rm eff}$, which can then be used to carefully tailor experimental searches. However, as is becoming increasingly clear, SUSY is a broad paradigm with several possible motivated incarnations and a complex parameter space, and it is a challenge for experiments to cover all of the phenomenological bases. One way to proceed is to try to constrain the form of $\mathcal{L}_{\rm eff}$ by more bottom-up criteria, and to use the results to guide experiments, committing perhaps to some broad UV principles but not committing to a specific UV model.

The most important such criterion is that SUSY-breaking in $\mathcal{L}_{\rm eff}$ be compatible with the radiative stability of the electroweak scale within the domain of validity of the effective theory, up to roughly 10 TeV. The significance of the 10 TeV scale is that almost all experiments, up to and including the LHC, only have sensitivity to new physics $\lesssim 10$ TeV, be it through direct searches or virtual effects. (In this regard, flavor physics tests are exceptional in probing vastly higher scales and consequently they require special consideration.) The fact that the non-supersymmetric Standard Model (SM) is already fine-tuned in this regime is known as the "little hierarchy problem", and provides the most immediate motivation for new physics accessible to colliders.

We must further weigh the relevance for the effective theory of other general concerns of the SUSY paradigm, which at least partly relate to very high energies:

- the SUSY Flavor Problem
- Grand Unification
- proton stability and R-parity
- superpartner dark matter candidates
- SUSY-breaking dynamics
- Higgs mass

In this paper, we will focus on the *minimal* effective theories that arise from the above viewpoint. They are "minimal" in terms of the particle content and parameter space of $\mathcal{L}_{\rm eff}$. This does not imply, however, that their UV-completions, above LHC energies, are also minimal in some way. Conversely, the Minimal Supersymmetric Standard Model (MSSM), is a minimal visible sector from the high-energy perspective, but is non-minimal in the sense that matters to the LHC effective theory and phenomenology, as we will review. The central observation, mentioned in [1, 2] and developed in [3], is that radiative stability between the weak scale and ~ 10 TeV, does not require a superpartner in the effective theory for every Standard Model (SM) particle, but just for those particles with order one couplings to the Higgs boson and electroweak breaking. In this way, the minimal superpartner content is given by the gauginos, Higgsinos, stop and sbottom, without sleptons or first and second generation squarks.[1] The omitted superpartners may have masses above LHC reach and may play a crucial role in weak scale stability up to much higher inscales, but all this is outside the scope of the effective theory and outside the grasp of the LHC. Ref. [4] dubbed this kind of structure, "Effective Supersymmetry". Since [3, 4], a number of quite different approaches to far-UV dynamics have converged on such a "more minimal" spectrum at accessible energies [5–8].

Of course, there is no guarantee that at accessible energies new SUSY physics will be turn out to be minimal. Rather, we study minimal LHC-effective theories for three reasons: (1) they represent possible SUSY phenomenology, and there do exist UV SUSY dynamics that match onto them, (2) a great deal of the natural parameter space remains open after one year of LHC data, and yet discoverable within the next year, (3) minimal models in any arena of exploration represent an important departure point for thinking more broadly.

In this paper, we will take a more UV-agnostic approach to the minimal effective theory at LHC energies than has been previously considered. We do not do this blindly, but only after discussion of the general SUSY concerns listed above. We will argue that modern developments in model-building and SUSY field theory have proliferated the range of UV options that relate to these issues, and it is precisely for this reason that we advocate thinking more modularly about them, and with less commitment to any one UV plot. Our goal will be to use electroweak naturalness, flavor constraints, minimality, and earlier searches as a guide to the LHC phenomenology, to discuss qualitative options (such as R-parity versus R-parity violation) and to organize the different possible channels and relevant parameter spaces. We will use this platform to study the LHC phenomenology in more detail, and in future work to broaden and help optimize experimental search strategies. We will adopt the name "Effective SUSY" to refer to this minimalist and UV-agnostic approach to the LHC-effective theory. Our study of effective SUSY coincides with the accumulation of significant LHC data. However, there are earlier collider studies relevant to effective SUSY on which our work expands, such as [9–15].

The paper is organized as follows. In section 2, we derive the minimal effective SUSY Lagrangian subject to electroweak naturalness with a cutoff of $\sim 10\,\mathrm{TeV}$. Here we impose R-parity and make the useful idealization that the third generation does not mix with the first two generations. We also make the standard assumption that the Higgsino mass arises from a supersymmetric μ term. In section 3, we perform the same exercise but with a cutoff of only $\sim 1\,\mathrm{TeV}$, in a sense increasing our agnosticism towards what lies above the early $7\,\mathrm{TeV}$ LHC reach. One possibility, but not the only one, is that this $1\,\mathrm{TeV}$ effective theory derives straightforwardly from the $10\,\mathrm{TeV}$ effective theory of section 2. In section 4, we study the possibility that Higgsinos obtain mass from soft SUSY breaking rather than a μ term, and we write an even more minimal set of effective Lagrangians with $10\,\mathrm{TeV}$ and $1\,\mathrm{TeV}$ cutoffs. In section 5 we put back consideration of third-generation mixing, and review and extend the constraints provided by low-energy flavor and CP tests. We emphasize the considerable safety of the effective SUSY scenario. In section 6, we make the case for R-parity violation as a very plausible option, write the effective SUSY R-parity

[1] We will be more precise later.

violating interactions, and discuss some of the low-energy constraints. In section 7, we discuss the interesting possibility of Dirac gauginos and how this can considerably affect the collider phenomenology and low-energy constraints. Section 8 is devoted to discussing collider phenomenology, in particular the 7 TeV LHC. As a first foray, we focus mostly on the minimal subsystem of stops, sbottoms and neutralino with R-parity. We also make brief remarks about other phenomenological regimes of effective SUSY. Section 9 provides our outlook.

While this paper was being completed, we became aware of three other groups pursuing partially overlapping work [16–18].

2 Effective SUSY $\lesssim$ 10 TeV

Let us start with the MSSM field content and ask which superpartners are minimally needed in order to maintain electroweak naturalness below 10 TeV, roughly the collider reach in the years to come. We will not ask here what physics lies above this scale. Therefore at the technical level, $\Lambda_{UV} \equiv 10\,\text{TeV}$ provides the cutoff for any UV divergences encountered in the effective theory, and this allows us to estimate electroweak fine-tuning and check where in parameter space effective SUSY solves the "little hierarchy problem" of the SM.

SM particles with order one couplings to the Higgs boson must certainly have superpartners in the effective theory because they would otherwise give rise to quadratically divergent Higgs mass-squared contributions at one loop, $\sim \Lambda_{UV}^2/(16\pi^2)$, big enough to require significant fine-tuning. In order to supersymmetrically cancel these divergences, the effective theory must therefore include the left-handed top and bottom squarks, $\tilde{q}_L \equiv (\tilde{t}_L, \tilde{b}_L)$, and the right-handed top squark, $\tilde{t}_R$, as well as the up-type Higgsino, $\tilde{h}_u$, and electroweak gauginos, $\lambda_{1,2}$.

Considerations beyond SUSY itself imply that we need to retain even more superpartners. Electroweak gauge anomaly cancellation implies that $\tilde{h}_u$ must be accompanied by $\tilde{h}_d$ in the effective theory. Indeed, one might have anticipated that down-type Higgs *bosons*, h_d, are required anyway to give masses to the down-type fermions, and that $\tilde{h}_d$ provide the required superpartners.[2] With the h_d bosons present in the effective theory, there is a new quadratic divergence, even in the supersymmetric limit, in the form of a (supersymmetric) hypercharge D-term. It is associated by supersymmetry with the mixed hypercharge-gravity triangle anomaly. The quadratic divergence vanishes only if $\mathrm{Tr}(Y) = 0$, where Y is the hypercharge charge matrix over the scalar fields of the effective theory. With the field content described, including h_d, this condition is not satisfied, and the theory remains unnatural despite superpartners for the main players in the SM. Vanishing $\mathrm{Tr}(Y)$ can be arranged by retaining the right-handed bottom squark, $\tilde{b}_R^c$, within the effective theory.

For the most part, two-loop quadratic divergences $\sim \Lambda_{UV}^2/(16\pi^2)^2$ are not important for Higgs naturalness, with a cutoff as low as 10 TeV. But the QCD coupling is an exception. In particular, the $\tilde{q}_L, \tilde{t}_R^c$ masses must themselves be so light in order to protect Higgs

[2]We proceed with this logic in this section, although there is a loop-hole whereby h_u can provide down-type fermion masses in the effective theory, and h_d bosons are not needed. We discuss this option in section 4.

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naturalness at one loop order, that they suffer from their own naturalness problem due to one-loop mass corrections from QCD. This one loop QCD destabilization of the squarks, hence two-loop destabilization of the Higgs, requires the gluino, λ_3, to be in the effective theory.

In this way, the effective theory has complete supermultiplets,

$$\begin{aligned} Q \equiv \begin{pmatrix} T \\ B \end{pmatrix} &\equiv (\tilde{q}_L, q_L) \equiv \left(\begin{pmatrix} \tilde{t}_L \\ \tilde{b}_L \end{pmatrix}, \begin{pmatrix} t_L \\ b_L \end{pmatrix} \right) \\ \bar{T} &\equiv (\tilde{t}^c_R, t^c_R) \\ \bar{B} &\equiv (\tilde{b}^c_R, b^c_R) \\ H_u &\equiv (h_u, \tilde{h}_u) \\ H_d &\equiv (h_d, \tilde{h}_d) \\ V_1 &\equiv (B_\mu, \lambda_1) \\ V_2 &\equiv (W_\mu, \lambda_2) \\ V_3 &\equiv (G_\mu, \lambda_3) \end{aligned} \tag{2.1}$$

where we use the lower case "h" to distinguish just the scalars of the Higgs chiral supermultiplet, "H".

2.1 Effective Lagrangian, neglecting third-generation mixing

Above, we have introduced squarks belonging to only the "third generation", and yet this notion is slightly ambiguous because generation-numbers are not conserved, even in the SM. However, CKM mixing involving the third generation is at least highly suppressed, so we will begin by considering the "zeroth order" approximation in which third-generation number is exactly conserved. For most purposes in LHC studies of the new physics, this approximation is sufficient. But for complete realism and to check the viability of the theory in the face of very sensitive low-energy flavor constraints, the extra subtlety of third-generation mixing must be taken into account. We defer this discussion until section 5. For now, this mixing is formally "switched off". Further, we will impose R-parity on effective SUSY, and defer the discussion of possible R-parity violating (RPV) couplings to section 6.

With the field content described above, the effective Lagrangian is given by

$$\begin{aligned} \mathcal{L}_{\text{eff}} = & \int d^4\theta K + \left(\int d^2\theta \left(\frac{1}{4}\mathcal{W}_\alpha^2 + y_t \bar{T} H_u Q + y_b \bar{B} H_d Q + \mu H_u H_d \right) + \text{h.c.} \right) \\ & + \mathcal{L}_{kin}^{light} - \left(\bar{u}_R Y_u^{light} h_u \psi_L + \bar{d}_R Y_d^{light} h_d \psi_L + \text{h.c.} \right) + \mathcal{L}_{lepton} \\ & - m^2_{\tilde{q}_L} |\tilde{q}_L|^2 - m^2_{\tilde{t}^c_R} |\tilde{t}^c_R|^2 - m^2_{\tilde{b}^c_R} |\tilde{b}^c_R|^2 - m^2_{h_u} |h_u|^2 - m^2_{h_d} |h_d|^2 \\ & - \left(m_{i=1,2,3} \lambda_i \lambda_i + B\mu h_u h_d + A_t \tilde{t}^c_R h_u \tilde{q}_L + A_b \tilde{b}^c_R h_d \tilde{q}_L + \text{h.c.} \right) \\ & + \mathcal{L}_{hard} + \mathcal{L}_{non-ren.}, \end{aligned} \tag{2.2}$$

where the first line is in superspace/superfield notation, while the remaining lines are in components. Here, K is the standard gauge-invariant Kähler potential for the chiral

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superfields of eq. (2.1), and $\mathcal{L}_{kin}^{light}$ denotes the standard gauge-invariant kinetic terms for the light SM quarks (that is, not the top and bottom), $u_R, d_R, \psi_L \equiv (u_L, d_L)$. $\mathcal{L}_{lepton}$ denotes all terms involving leptons, with Yukawa couplings to h_d (neglecting neutrino mass terms). The super-field strength tensors are implicitly summed over all three gauge groups of the standard model, both here and throughout the paper. Even the second line can be thought of as the result of starting from the supersymmetric MSSM, but then deleting all superpartners for light SM fermions. As mentioned above, we ignore the third generation mixing with the first two generations (until section 5). The third and fourth lines are soft SUSY breaking terms for the superfields of the effective theory.

The absence of superpartners for the light fermions will necessarily induce *hard* SUSY-breaking divergences at one-loop order. To renormalize these, we must include hard SUSY breaking couplings into the effective Lagrangian, and naturalness dictates that the renormalized couplings be at least of one-loop strength, $\gtrsim 1/(16\pi^2)$. These couplings are included in the last line, in $\mathcal{L}_{hard}$. Such couplings can then appear within one-loop Higgs self-energy diagrams, yielding two-loop sized quadratic divergences, $\gtrsim \Lambda_{UV}^2/(16\pi^2)^2$. While this is acceptable from the viewpoint of naturalness, we see that we cannot tolerate order one hard breaking couplings. UV completions of effective SUSY theory can contain mechanisms to naturally yield such non-vanishing, but suppressed, hard breaking terms, for example [5, 7]. Because the hard breaking is necessarily small, it is largely negligible for early LHC phenomenology. On the other hand, at a later stage of exploration, measuring hard SUSY breaking such as a difference between gauge and gaugino couplings may provide a valuable diagnostic.

Effective SUSY is expected to arise from integrating out heavy physics above 10 TeV, some of which is crucial in solving the hierarchy problem to much higher scales. It should therefore be a non-renormalizable effective theory, with higher-dimension interactions suppressed by $\sim$ 10 TeV or more. These are contained in $\mathcal{L}_{non-ren}$ on the last line. Again, these will be largely irrelevant for early LHC phenomenology, but can very important in precision low-energy experiments, such as CP or flavor tests. The most stringent of such tests imply that at least some non-renormalizable interactions have to be suppressed by effective scales much beyond 10 TeV. Again, there are UV completions of effective SUSY which possess natural mechanisms to explain this required structure.

2.2 Higgs mass

The experimental bounds on the lightest *physical* neutral Higgs scalar provide some of the most stringent constraints on weak scale SUSY. The dominant couplings of our effective Lagrangian are just those of the MSSM, so the electroweak symmetry-breaking and Higgs-mass predictions are essentially the same. This is problematic because naturalness dictates stops lighter than a few hundred GeV, while the physical Higgs mass constraints require higher stop masses. One difference with the high-scale MSSM is that in effective SUSY we have hard SUSY breaking couplings, among which can be Higgs quartic couplings which ultimately contribute to the physical Higgs mass. However, these contributions are modest, just a few GeV, since the hard SUSY-breaking couplings must be suppressed for electroweak naturalness. Instead, sizeable upward contributions to physical Higgs mass

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require new particle content beyond the MSSM (see e.g. [19] and references therein). For example, this is readily accomplished by adding a chiral superfield gauge singlet to the effective theory [20–23],

$$\delta\mathcal{L}_{\text{eff}} = \int d^4\theta |S|^2 + \int d^2\theta \left(\kappa S H_u H_d + \frac{1}{2}\sigma S^2\right) + \text{h.c.}$$
$$-m_s^2|s|^2 + \text{ other soft terms} \tag{2.3}$$

which contains a new contribution to the Higgs quartic couplings, $\sim \kappa^2$. The soft scalar mass-squared term m_s^2 can be $O(\text{TeV}^2)$ without destabilizing EWSB. It can also ensure that the singlet does not acquire a vacuum expectation. In principle, in effective SUSY with a 10 TeV cutoff, we must commit to which type of physics, $\delta\mathcal{L}_{\text{eff}}$, accounts for an acceptable physical Higgs mass. But for early LHC superpartner searches, the details of $\delta\mathcal{L}_{\text{eff}}$ need not be relevant, as the new particles can lie above 1 TeV. In such cases, the new physics is just a "black box" which gives viable physical Higgs masses. Indeed, in writing effective SUSY theories with a lower $\sim$ 1 TeV cutoff, we will see that we can formally imagine having integrated out the new physics responsible for new Higgs quartic couplings.

2.3 Naturalness in effective SUSY

Here, we assemble the electroweak naturalness constraints on effective SUSY, thereby giving a rough idea of the motivated regions of its parameter space. For this purpose, we will compute various independent corrections to the h_u mass-squared, and simply ask them to be $\lesssim$ (200 GeV)^2 for naturalness. We will compute these corrections *before* EWSB. Contributions sensitive to EWSB are typically $\sim O((100 \text{ GeV})^2)$, and therefore typically do not compromise naturalness. Given the intrinsically crude nature of naturalness arguments, we see no merit in a more refined analysis.

We begin with a classical "tuning" issue. The μ term gives a supersymmetric $|\mu|^2$ contribution to the Higgs mass-squareds. While the soft terms also contribute to Higgs mass-squareds, naturalness forbids any fine cancellations, so therefore by the criterion stated above,

$$|\mu| \lesssim 200 \text{ GeV}. \tag{2.4}$$

This same parameter then also plays the role of the Higgsino mass parameter, ensuring relatively light charginos and neutralinos in the superpartner spectrum. (Of course, after EWSB, these physical states may also contain admixtures of electroweak gauginos.)

Next, we turn to quantum loops. We assume that $\tilde{q}_L, \tilde{t}_R$ have approximately the same mass, $m_{\tilde{t}}$, for simplicity, and we also neglect the μ and A-terms. We work pre-EWSB since we are concerned with sensitivity to parametrically higher scales. By evaluating the diagrams in figure 1, we find that the $m_{h_u}^2$ parameter receives the following correction:

$$\delta m_{h_u}^2 = -\frac{3y_t^2}{4\pi^2} m_{\tilde{t}}^2 \ln\left(\frac{\Lambda_{UV}}{m_{\tilde{t}}}\right) \tag{2.5}$$

Naturalness therefore requires, very roughly,

$$m_{\tilde{t}} \lesssim 400\text{GeV}. \tag{2.6}$$

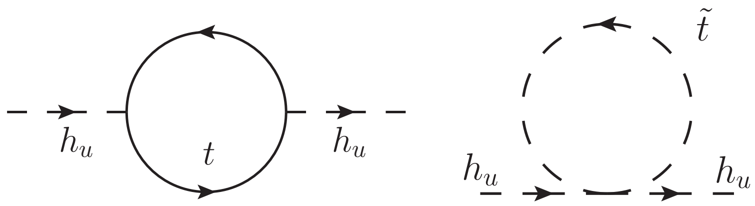

Figure 1. Higgs mass corrections

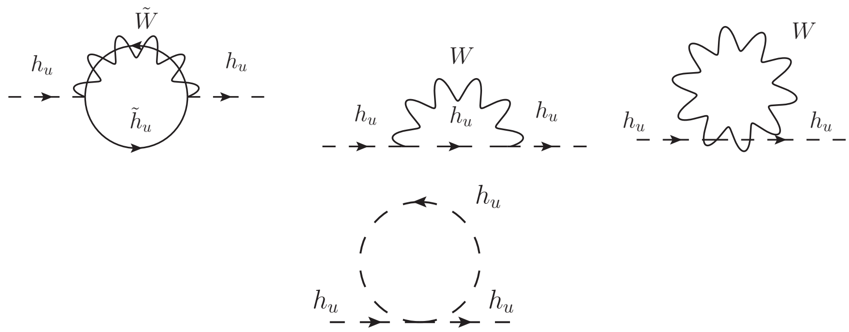

Figure 2. Higgs mass correction

There are also electroweak gauge/gaugino/Higgsino one-loop contributions to Higgs mass-squared. Again, working before electroweak symmetry breaking (gaugino-Higgsino mixing) and just looking at the stronger $SU(2)_L$ coupling, the Higgs self-energy diagrams are in figure 2.

The Higgs mass correction is then given by

$$\delta m_{h_u}^2 = \frac{3g^2}{8\pi^2}\left(m_{\tilde{W}}^2 + m_{\tilde{h}}^2\right)\ln\frac{\Lambda_{UV}}{m_{\tilde{W}}}. \tag{2.7}$$

We identify the Higgsino mass with μ. Because we are already taking $\mu \lesssim 200\,\mathrm{GeV}$, this translates into a roughly natural wino mass range of

$$m_{\tilde{W}} \lesssim \mathrm{TeV}. \tag{2.8}$$

Next, we compute the hypercharge D-term loop contribution to Higgs mass-squared, in figure 3.

This gives rise to a higgs mass correction:

$$\delta m_{h_u}^2 = \sum_{\text{scalars } i} \frac{g'^2 Y_i Y_{h_u}}{16\pi^2}\left(\Lambda_{UV}^2 - m_i^2 \ln\frac{\Lambda_{UV}^2 + m_i^2}{m_i^2}\right). \tag{2.9}$$

Including both the right-handed sbottom and the down-type higgs, as we do in this section, ensures that the quadratic divergence cancels, but there is still a residual correction to the higgs mass. Given that other scalars have already been argued to be relatively light, we can use this correction to estimate the natural range for the mass of $\tilde{b}_R$,

$$m_{\tilde{b}_R} \lesssim 3\mathrm{TeV}. \tag{2.10}$$

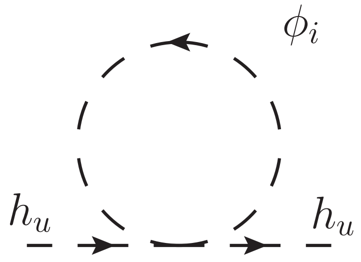

Figure 3. Higgs mass correction

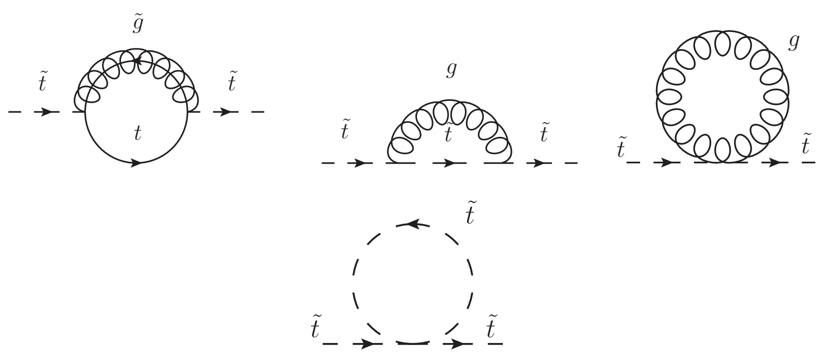

Figure 4. Stop mass correction

Finally, $\tilde{q}_L, \tilde{t}_R$ also being relatively light scalars, suffer from their own naturalness problem, with mass corrections dominated by the diagrams in figure 4.

This gives rise to a stop mass correction:

$$\delta m_{\tilde{t}}^2 = \frac{2g_s^2}{3\pi^2} m_{\tilde{g}}^2 \ln \frac{\Lambda_{UV}}{m_{\tilde{g}}}. \tag{2.11}$$

For squark masses $\sim$ few hundred GeV, naturalness requires

$$m_{\tilde{g}} \lesssim 2m_{\tilde{t}}. \tag{2.12}$$

3 Effective SUSY $\lesssim$ 1 TeV

Although the LHC has a multi-TeV reach in principle, parton distribution functions fall so rapidly at high energies that most parton collisions have sub-TeV momentum transfers. In the early LHC era, statistically signficant effective SUSY signals would be in this regime. For example, in effective SUSY, gluino production would have a cross-section of just a few fb for TeV gluino mass. We can therefore focus our attention on just the early accessible physics by constructing a rough effective SUSY theory with a cutoff $\Lambda_{UV} \sim$ TeV, while not committing strongly to the physics above this scale. With such a low cutoff, only top quark loops in the SM destabilize Higgs naturalness. This is cured by SUSY cancellation upon including the squarks, $\tilde{q}_L, \tilde{t}_R$, to form complete supermultiplets, $Q \equiv (\tilde{q}_L, q_L), \bar{T} \equiv$

$(\tilde{t}^c_R, t^c_R)$, as before. Even hypercharge D-term divergences from the uncanceled $\mathrm{Tr}(Y)$ are not quantitively significant. It therefore appears that we can dispense with Higgsinos, $\tilde{b}_R$, and the gauginos in the effective theory. However, if Higgsino mass arises from a supersymmetric μ term, as discussed in subsection 2.3, then electroweak naturalness also forces the Higgsinos to be light. We will continue with this assumption in this section, and therefore retain complete supermultiplets, $H_{u,d} \equiv (h_{u,d}, \tilde{h}_{u,d})$.

Even though we do not commit here to the structure of the theory above 1 TeV, one possibility is that it is just that of the last section.[3] But in that case, by eq. (2.11), we should include the gluino in the sub-TeV effective theory. However, non-minimal physics in the $1 - 10$ TeV window can change this conclusion, and indeed the gluino might naturally be considerably heavier than 1 TeV. We illustrate such new physics in section 7, with the example of a Dirac gluino. It exemplifies the general theme that non-minimal UV physics can lead to more minimal IR physics, while still being compatible with naturalness. Here, we merely check within the TeV effective theory that naturalness indeed requires stops, but that these stops *do not require* gluinos. The first statement follows from eq. (2.5), where naturalness up to $\Lambda_{UV} \sim 1$ TeV then implies

$$m_{\tilde{t}} \lesssim 700\mathrm{GeV}. \tag{3.1}$$

The second statement follows from eq. (2.11), where we see that with the logarithm of order one and gluino mass ~ 1 TeV, we can naturally have stops as light as 300 GeV. In our phenomenological studies of section 8, we mostly keep in mind lighter stops, $m_{\tilde{t}} \lesssim 400\mathrm{GeV}$, compatible with either 1 or 10 TeV cutoffs as discussed in section 2.

3.1 Effective Lagrangian, neglecting third-generation mixing

Given the light superpartner content described above, the R-parity conserving effective theory below a TeV is given by

$$\begin{aligned}\mathcal{L}_{\rm eff} = &\int d^4\theta K + \left(\int d^2\theta \left(\frac{1}{4}\mathcal{W}_\alpha^2 + y_t \bar{T} H_u Q + y_b \bar{B} H_d Q + \mu H_u H_d\right) + \text{h.c.}\right)\\ &+\mathcal{L}^{light}_{kin} - \left(\bar{u}_R Y_u^{light} h_u \psi_L + \bar{d}_R Y_d^{light} h_d \psi_L + \text{h.c.}\right) + \mathcal{L}_{lepton}\\ &-m^2_{\tilde{q}_L}|\tilde{q}_L|^2 - m^2_{\tilde{t}^c_R}|\tilde{t}^c_R|^2 - m^2_{h_u}|h_u|^2 - m^2_{h_d}|h_d|^2\\ &-\left(B\mu h_u h_d + A_t \tilde{t}^c_R h_u \tilde{q}_L + \text{h.c.}\right)\\ &+\mathcal{L}_{hard} + \mathcal{L}_{non-ren.},\end{aligned} \tag{3.2}$$

[3]While the LHC might be dominated by sub-TeV physics, as explained above, electroweak precision tests at lower energy machines are famously sensitive to multi-TeV scales via virtual processes. In the 1 TeV effective theory, this translates to precision test sensitivity to higher dimensional operators. In the case, where this effective theory merely originates from our 10 TeV effective theory, such higher-dimensional operators are suppressed by the 10 TeV scale and are safe from electroweak precision tests. We take this as an existence proof that multi-TeV physics of the sort we contemplate can easily yield sufficiently suppressed higher-dimensional operators in the TeV effective theory to be safe, and make it an assumption for our consideration of TeV effective theory in general. We do not further specify the structure of such operators, given their lower relevance for LHC processes.

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which is to be interpreted as in eq. (2.2) except that all *terms involving gauginos or $\tilde{b}^c_R$ are to be thrown away* after expanding the superspace expressions in components.

With the cutoff as low as 1 TeV, the hard SUSY-breaking can now include $|h_u|^4$ couplings strong enough to give contributions to the physical Higgs mass of tens of GeV without making EWSB scale unnatural. One can think of these terms as arising from new fields, such as discussed in section 2.2, heavier than 1 TeV, which have therefore been integrated out. One virtue of this sub-TeV theory is that we do not have to commit to just what UV physics contributed to Higgs mass; whatever it might be is parametrized by the effective hard couplings.

3.2 Dark Matter considerations

In our TeV effective theory, we must take the Higgsinos as the lightest superpartners in order to avoid phenomenologically dangerous colored (collider-)stable particles in the form of stops or sbottom. Such Higgsinos will then form charginos and neutralinos at the ends of superpartner decay chains. Higgsino neutralinos would have a thermal relic abundance smaller than needed to fully account for all of dark matter. This is not an issue if dark matter is dominated by other physics not (soon to be) accessible to the LHC. Another possibility is that the wino and bino, $\lambda_{1,2}$, which are not required to be light by naturalness, are nevertheless light and in the effective theory, and a linear combination of gaugino-Higgsino forms a neutralino LSP. It is possible then that such a hybrid LSP has the correct thermal relic abundance to account for dark matter. This computation still remains to be checked in the effective SUSY context however. Even in this case, our minimal effective theory is still useful, in that for the purposes of *early* LHC phenomenology the details of charginos/neutralinos are not as important as their existence and the LSP mass. The Higgsino LSP in our effective theory can therefore serve as a toy model of whatever the real chargino/neutralino degrees of freedom are. More refined modeling can wait until the new physics is discovered.

4 Effective SUSY with heavy Higgsinos

4.1 Effective SUSY with 10 TeV cutoff

As alluded to earlier, given that we necessarily have hard SUSY breaking couplings in effective SUSY, we can reduce the particle content even further by eliminating h_d *bosons* and the right-handed bottom squark $\tilde{b}^c_R$ from the effective theory. See refs. [24–26] for earlier related works. This move maintains the vanishing of $\mathrm{Tr}(Y)$ required for naturalness with 10 TeV cutoff, but forces us to obtain Yukawa-couplings for down-type fermions by coupling them to

$$h_u^* \equiv i\sigma_2 h_u^\dagger, \tag{4.1}$$

where σ_2 is the second weak-isospin Pauli matrix. This is the usual approach to getting down-type fermion masses in the SM with a single Higgs doublet. In the SUSY context, such a coupling cannot arise from a superpotential, which can only depend on H_u, not $H_u^\dagger$. Instead, it represents a hard SUSY breaking effect (though it may arise from soft SUSY breaking from the vantage of a UV completion). It poses no threat to naturalness if the couplings are $\ll 1$. This is certainly the case for all the down-type Yukawa couplings.

4.1.1 Effective Lagrangian, neglecting third-generation mixing

With the particle content described above, the R-parity conserving effective Lagrangian is given by

$$\begin{aligned}
\mathcal{L}_{\text{eff}} = & \int d^4\theta K + \left(\int d^2\theta \left(\frac{1}{4}\mathcal{W}_\alpha^2 + y_t \bar{T} H_u Q \right) + \text{h.c.} \right) \\
& + \mathcal{L}_{kin} - \left(\bar{u} Y_u^{light} h_u q_L + y_b \bar{b} h_u^* q_L + \bar{d} Y_d^{light} h_u^* q_L + \text{h.c.} \right) + \mathcal{L}_{lepton} \\
& - m_{\tilde{q}_L}^2 |\tilde{q}_L|^2 - m_{\tilde{t}_R^c}^2 |\tilde{t}_R^c|^2 - m_{h_u}^2 |h_u|^2 \\
& - \left(m_{i=1,2,3} \lambda_i \lambda_i + A \tilde{t}_R^c h_u \tilde{q}_L + m_{\tilde{h}} \tilde{h}_u \tilde{h}_d + \text{h.c.} \right) \\
& + \mathcal{L}_{hard} + \mathcal{L}_{non-ren.}.
\end{aligned} \tag{4.2}$$

The Kahler potential K consists of the gauge-invariant kinetic terms for the chiral superfields, $\bar{T}, Q, H_u$, while compared with eq. (2.2), the kinetic terms for the (now un-superpartnered) fermions b_R and $\tilde{h}_d$ have now been added to $\mathcal{L}_{kin}$. The second to fourth lines still follow from the MSSM after deleting fields that are absent in our effective theory, except for the small Yukawa couplings of h_u^* to down-type fermions, which we pointed out above are a form of hard SUSY breaking. Other hard breaking as well as non-renormalizable couplings appear on the last line. Our discussion of the physical Higgs mass, and contributions to it, is similar to subsubsection 2.2. However a singlet coupling to $h_u h_d$ is not possible since we have removed h_d, but in an electroweak triplet coupled to $H_u H_u$ is possible and results in a $|h_u|^4$ terms in the potential [27–29].

4.1.2 Higgsino mass

Note that the Higgsino mass now takes the form of a soft SUSY-breaking mass term, $m_{\tilde{h}}$, as opposed to a supersymmetric μ term as in section 2. In this way, it is uncorrelated with any contribution to Higgs boson mass-squared. Therefore, there is only one modification to the bounds obtained in section 2.3; namely, that now $m_{\tilde{h}}$ is only constrained by eq. (2.7), so that

$$m_{\tilde{h}} \lesssim \text{TeV}. \tag{4.3}$$

4.2 Effective SUSY $\lesssim 1$ TeV

In the most minimal of our effective theories, all gauginos and Higgsinos can naturally be heavier than a TeV and thus integrated out of the sub-TeV effective theory. If we identify h_u with the SM Higgs doublet, the only new particles are $\tilde{t}_L, \tilde{b}_L, \tilde{t}_R$.

4.2.1 Effective Lagrangian, neglecting third-generation mixing

The effective Lagrangian with R-parity is then given by

$$\begin{aligned}\mathcal{L}_{\text{eff}} = &\ \mathcal{L}_{SM} + \mathcal{L}_{kin}^{squarks} - V_{D-terms} - y_t^2(|h_u\tilde{q}_L|^2 + |\tilde{t}_R^c\tilde{q}_L|^2 + |\tilde{t}_R^c h_u|^2) \\ &- m_{h_u}^2|h_u|^2 - m_{\tilde{q}_L}^2|\tilde{q}_L|^2 - m_{\tilde{t}_R^c}^2|\tilde{t}_R^c|^2 - \left(A\tilde{t}_R^c h_u\tilde{q}_L + \text{h.c.}\right) \\ &+ \mathcal{L}_{hard} + \mathcal{L}_{non-ren.},\end{aligned} \tag{4.4}$$

$\mathcal{L}_{SM}$ is the SM Lagrangian with h_u playing the role of the SM Higgs doublet, but with no Higgs potential. The Higgs potential is a combination of the soft Higgs mass term in the second line, the D-term potential and possible hard SUSY-breaking couplings $\sim |h_u|^4$. As discussed in subsection 2.1, these hard SUSY breaking couplings can be large enough to easily satisfy the Higgs mass bound without spoiling naturalness.

With exact R-parity, one of the colored superpartners would necessarily be stable and phenomenologically dangerous. However, we can use the above effective Lagrangian as the minimal departure point for adding R-parity violating corrections. We take this up in section 6.

4.2.2 Effective Lagrangian $\lesssim$ TeV, with neutralino LSP

Another possibility is that R-parity is exact but there is a neutralino LSP in the spectrum, even though it is not required by electroweak naturalness. It may or may not be the dominant constituent of dark matter. Since we cannot determine its identity by theoretical considerations alone, we will just add a temporary "place-holder", that allows the squarks to decay promptly while preserving R-parity. We choose this to be the bino, λ_1, even though taken literally, it would predict too large a thermal relic abundance of dark matter. A more refined description of the neutralino would not add much to the early LHC search strategy. In this option, as compared to that of subsection 3.2 and eq. (3.2), we do not have a chargino.

The effective Lagrangian then takes the form

$$\begin{aligned}\mathcal{L}_{\text{eff}} = &\ \mathcal{L}_{SM} + \mathcal{L}_{kin}^{squarks} - V_{D-terms} - y_t^2(|h_u\tilde{q}_L|^2 + |\tilde{t}_R^c\tilde{q}_L|^2 + |\tilde{t}_R^c h_u|^2) \\ &- m_{h_u}^2|h_u|^2 - m_{\tilde{q}_L}^2|\tilde{q}_L|^2 - m_{\tilde{t}_R^c}^2|\tilde{t}_R^c|^2 - \left(A\tilde{t}_R^c h_u\tilde{q}_L + \text{h.c.}\right) \\ &+ i\bar{\lambda}_1\partial.\sigma\lambda_1 - (m_1\lambda_1\lambda_1 + \text{h.c.}) - \sqrt{2}g'\left(\frac{1}{6}\bar{q}_L\bar{\lambda}_1\tilde{q}_L + \frac{1}{6}\bar{\tilde{q}}_L\lambda_1 q_L - \frac{2}{3}\bar{t}_R^c\bar{\lambda}_1\tilde{t}_R^c - \frac{2}{3}\bar{\tilde{t}}_R^c\lambda_1 t_R^c\right) \\ &+ \mathcal{L}_{hard} + \mathcal{L}_{non-ren.},\end{aligned} \tag{4.5}$$

5 Flavor-Changing Neutral Currents and CP Violation

Above, we have worked in the drastic approximation that the mixing between the third generation with the first two generations vanishes, so that the meaning of "third generation" squarks, $\tilde{q}_L, \tilde{t}_R^c, \tilde{b}_R^c$, is completely unambigous. In this limit, there is a conserved third-generation (s)quark number. In the real world, third generation mixing is non-zero but small. In Wolfenstein parametrization, mixing with the second generation is of order ϵ^2 and mixing with the first generation is of order ϵ^3, where $\epsilon \sim 0.22$ corresponds to Cabibbo mixing. Given this fact, it is more natural to have comparable levels of violation of third-generation (s)quark number in the physics we have added beyond the SM.

In practice this means that for every interaction term in which the squarks currently appear, where third-generation number is conserved by the presence of t or b quarks (in electroweak gauge basis), we now allow more general couplings, with the third generation quarks replaced by quarks of the first and second generations. The associated couplings with second generation quarks are taken to be of order ϵ^2, while those with first generation quarks are taken to be of order ϵ^3, all in electroweak gauge basis. All these couplings involving the squarks are technically hard breaking of SUSY, but $\epsilon^{2(3)}$ is so small that, like other hard breaking in the effective theory, they do not spoil Higgs naturalness below 10 TeV. For most, but not all, of the LHC collider phenomenology the small $\epsilon^{2(3)}$ effects are negligible and we can proceed with our earlier effective Lagrangians. (We must of course keep SM third generation mixing effects, so that, for example, the bottom quark decays.) But in the more realistic setting with third-generational mixing, we must confront the SUSY flavor problem. In effective SUSY, this problem has two faces, IR and UV.

The UV face of the problem is contained in the non-renormalizable interactions of eq. (2.2). For example, they can include flavor-violating interactions such as $\bar{s}d\bar{s}d$. If such a non-renormalizable interaction were suppressed only by $(10\,\text{TeV})^2$, it would lead to FCNCs in kaon mixing, orders of magnitude greater than observed. It is therefore vital for the non-renormalizable interactions to have a much more benign flavor structure. Whether this is the case or not is determined by matching to the full theory above 10 TeV, IR effective SUSY considerations alone cannot decide the issue. Refs. [5, 7] are examples of UV theories which reduce to effective SUSY at accessible energies and automatically come with the kind of benign UV flavor structure we require. In this paper, we simply assume that the UV-sensitive non-renormalizable interactions are sufficiently flavor-conserving to avoid conflict with FCNC constraints.

There remain FCNC effects that are UV-insensitive but are assembled in the IR of the effective theory through the small $\epsilon^{2(3)}$ flavor-violating couplings. Many of these have been studied in refs. [30] and are small enough to satisfy current constraints. Indeed this feature is one of the selling points of effective SUSY. Here, we illustrate one such FCNC effective interaction for (CP-violating) $K - \bar{K}$ mixing arising as a SUSY "box" diagram. Similar processes were studied in [31–34], with minor adaptations needed in our case. While the effect is suppressed by $\mathcal{O}(\epsilon^{10})$ in effective SUSY, it is more stringently constraining than $B_d - \bar{B}_d$ mixing or $B_s - \bar{B}_s$ mixing, even though these are suppressed by just $\mathcal{O}(\epsilon^6)$ and $\mathcal{O}(\epsilon^4)$ respectively. We show that with our rough flavor-changing power-counting the $\tilde{b}_R$ squark is constrained to lie above several TeV in the absence of flavor-parameter tuning.

In a low-energy effective Lagrangian to be run down to the hadronic scale, we match onto effective operators of the form

$$\mathcal{L}_{\text{eff}} \supset \kappa(\bar{s}_L d_R)(\bar{s}_R d_L). \tag{5.1}$$

Strictly speaking there are two different operators depending on color contraction. As shown in [31] an operator $\mathcal{O}_5 \propto \bar{d}^i_R s^j_L \bar{d}^j_L s^i_R$ (where i, j are color indices) is not enhanced by QCD running and has $1/N_c$-suppressed QCD matrix element. Therefore we concentrate

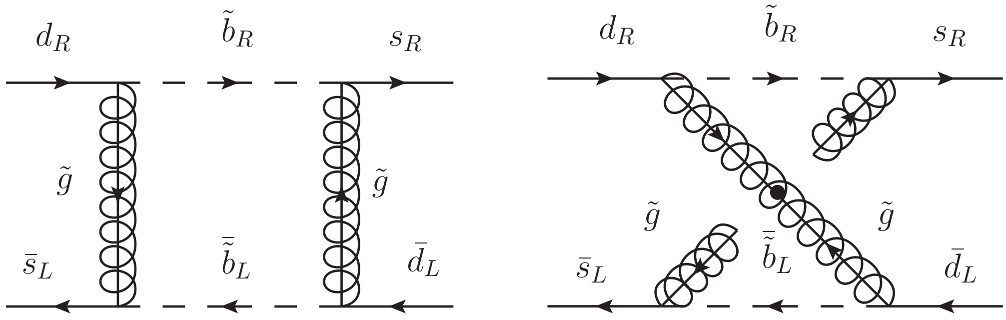

Figure 5. Contributions to $K-\bar{K}$ mixing

on $\mathcal{O}_4 \propto \bar{d}^i_R s^i_L \bar{d}^j_L s^j_R$, which has enhanced QCD running and large hadronic matrix element. Therefore, for the purpose of our simple estimate, in (5.1) we only study the case where each bilinear is a color singlet.

Integrating out the superpartners yields (see figure 5 for corresponding diagrams):

$$
\begin{aligned}
\kappa \sim &-\frac{g_s^4\epsilon^{10}}{4\pi^2}\frac{2}{3}\frac{m_3^2}{(m_3^2-m^2_{\tilde{q}_L})^2(m_3^2-m^2_{\tilde{b}^c_R})^2(m^2_{\tilde{q}_L}-m^2_{\tilde{b}^c_R})} \\
&\qquad\times\Bigg((m^2_{\tilde{q}_L}-m^2_{\tilde{b}^c_R})(m_3^2-m^2_{\tilde{b}^c_R})(m_3^2-m^2_{\tilde{q}_L})+m^2_{\tilde{q}_L}(m_3^4+m^4_{\tilde{b}^c_R})\ln\frac{m^2_{\tilde{q}_L}}{m_3^2} \\
&\qquad+2m_3^2m^2_{\tilde{q}_L}m^2_{\tilde{b}^c_R}\ln\frac{m^2_{\tilde{b}^c_R}}{m^2_{\tilde{q}_L}}+m^2_{\tilde{b}^c_R}(m^4_{\tilde{q}_L}+m_3^4)\ln\frac{m_3^2}{m^2_{\tilde{b}^c_R}}\Bigg) \\
&\quad-\frac{g_s^4\epsilon^{10}}{8\pi^2}\frac{1}{12}\frac{1}{(m_3^2-m^2_{\tilde{q}_L})^2(m_3^2-m^2_{\tilde{b}^c_R})^2(m^2_{\tilde{q}_L}-m^2_{\tilde{b}^c_R})} \\
&\qquad\times\Bigg((2m_3^2m^4_{\tilde{q}_L}m^2_{\tilde{b}^c_R}-m_3^4m^4_{\tilde{q}_L})\ln\frac{m_3^2}{m^2_{\tilde{q}_L}}+(2m_3^2m^2_{\tilde{q}_L}m^4_{\tilde{b}^c_R}-m_3^4m^4_{\tilde{b}^c_R})\ln\frac{m^2_{\tilde{b}^c_R}}{m_3^2} \\
&\qquad+m^4_{\tilde{q}_L}m^4_{\tilde{b}^c_R}\ln\frac{m^2_{\tilde{q}_L}}{m^2_{\tilde{b}^c_R}}+m_3^2(m_3^2-m^2_{\tilde{q}_L})(m_3^2-m^2_{\tilde{b}^c_R})(m^2_{\tilde{q}_L}-m^2_{\tilde{b}^c_R})\Bigg)\,,
\end{aligned}
\tag{5.2}
$$

where, as discussed above, the squark couplings to second generation quarks are assigned strength $\sim g_s\epsilon^2$, while squark couplings to the first generation are $\sim g_s\epsilon^3$. We neglect $\tilde{b}_L$-$\tilde{b}_R$ mixing (after EWSB). Note that our result contains large logarithms of the form $\ln m^2_{\tilde{b}_R}/m^2_{\rm squark}$, which in principle should be resummed (for example, see [34]). However, we do not do this since, again, we only seek an estimate for κ.

Current constraints on ϵ_K require that [35]

$$
(\mathrm{Im}(\kappa)) \lesssim \left(\frac{1}{3\times 10^5\ \mathrm{TeV}}\right)^2. \tag{5.3}
$$

For $m_3 \sim$ TeV and $m_{\tilde{q}_L} \sim 350\,$GeV, this translates into a bound on $\tilde{b}_R$ mass of roughly $m_{\tilde{b}^c_R} \gtrsim 17$ TeV.

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Of course, this bound is extremely sensitive to our estimates for the flavor-changing vertices. For example, if each flavor-changing vertex were only half as strong as our above estimates, the bound would be relaxed to $m_{\tilde{b}^c_R} \gtrsim 4$ TeV, roughly consistent with the requirements of naturalness in subsection 2.3. Alternatively, there may be small phases present in the vertices that further suppress κ. In the even more minimal effective supersymmetry structure of section 4, $\tilde{b}^c_R$ is completely absent and there is no robust infrared contribution to κ at one-loop order to worry about.

There are also CP-violating effects unrelated to flavor-changing, in particular electric dipole moment (EDM) constraints. From [36, 37] (see also references therein), we see that again effective SUSY has a relatively safe IR structure, with large regions of viable parameter space.[4] For example, see case II of table III in ref. [37] and the surrounding discussion. We show that these constraints are even more relaxed in the case of Dirac gauginos, in section 7.

6 R-Parity versus R-Parity Violation

R-parity plays a central role in theory and phenomenology within the weak scale SUSY paradigm. We will review some of the reasons for this, and argue that in light of several modern theoretical developments, the case for R-parity conservation in effective SUSY is less compelling. We are therefore more strongly motivated to take seriously an R-parity violating phenomenology. Quite apart from these theoretical considerations, we believe that this RPV phenomenology of effective SUSY is quite distinctive, and has so far received little attention. We will take up its study in future work.

6.1 Proton decay

The standard motivation for R-parity is that it leads to conserved baryon number. But it does not follow in complete generality. In the MSSM, baryon-number conservation only follows from R-parity after restricting to renormalizable interactions. For example, R-parity conserving but non-renormalizable superpotential interactions of the general form $W \propto \bar{U}\bar{U}\bar{D}\bar{E}$ give rise to proton decay. If the MSSM is taken as valid up to an extremely high scale, such a non-renormalizable term, and the resulting proton decay rate, would be suppressed by that high scale. However, if the MSSM is an effective theory emerging only below some lower threshold, then the non-renormalizable operator can be suppressed by just this lower threshold scale, leading to excessive proton decay. This is precisely the issue in many SUSY GUT theories, where such an effective interaction arises in the effective MSSM after integrating out a color-triplet GUT-partner of the Higgs. The moral only gains strength in effective SUSY, with a 10 TeV cutoff. For example, a dimension-6 R-parity conserving operator such as $u_L d_L u_R e_R$ can be viewed as a remnant of a supersymmetric non-renormalizable Kahler potential term. It gives rise to extremely rapid proton decay if suppressed by just $(10 \text{ TeV})^2$. Such an operator might well arise upon integrating out new thresholds above 10 TeV.

[4] Here, we are discussing the supersymmetric CP problem as opposed to the Strong CP problem. We assume for concreteness that the Peccei-Quinn mechanism with an axion resolves the Strong CP problem.

We conclude that R-parity is not by itself enough to protect against proton decay in effective SUSY, in general we need some other symmetry, such as baryon-number or lepton-number symmetry.[5] Clearly then, the proton-stability motivation for R-parity is gone.

6.2 Unification

Traditionally, the reason for arguing against new physics thresholds between the GUT and weak scales is because such new physics generally spoil the success of gauge coupling unification. But this is evaded if the new physics comes in complete GUT multiplets. For example, this is what is typically assumed for the messenger threshold of gauge-mediated SUSY-breaking models. In the model-building of recent years, we have seen that even quite radically new intermediate structure can maintain the success of gauge coupling unification by following this basic rule of GUT-degenerate thresholds [38]. There also exist new unification mechanisms that improve on the imperfect unification of SM via strong coupling effects over intermediate scales [39]. Therefore, we cannot have confidence that there is a Weak-GUT desert, as is often assumed. There may well be important new physics (not far) above 10 TeV, and in this context R-parity does not save us from excessive proton decay, as discussed above.

Another GUT-related reason in favor of R-parity is that in the context of traditional GUT models, imposing baryon- or lepton-number symmetry conflicts with the unification of quarks and leptons, whereas imposing R-parity does not. However, such traditional GUT models also suffer from other difficulties such as the notorious doublet-triplet splitting problem. In more recent years, it has been understood that some of the successes attributed to SUSY GUTs can arise more generally, in particular in the context of Orbifold GUT models (see [40, 41] and references therein). Such models employ "split multiplets", in which quarks and leptons can naturally arise as incomplete parts of *separate* GUT multiplets, and the Higgs doublet and triplet are also neatly split in the same manner. In this orbifold unification context, one can straightforwardly impose baryon- or lepton-number symmetry, safeguarding proton stability without requiring R-parity.

In this way, the unification considerations that originally favored R-parity over baryon- or lepton-number symmetry are less compelling.

6.3 Dark Matter

There is a second traditional motivation for R-parity, namely that the lightest R-odd superpartner is stable, and therefore may account for the dark matter of the Universe, enjoying the rough quantitative success known as the "WIMP-miracle". RPV interactions spoil this stability and seem to rob us of such a dark matter candidate. However, it is entirely possible that dark matter does consist of weak scale WIMPs, but these WIMPs are stabilized by carrying a different symmetry than R-parity, under which the SM is inert. This natural possibility leads us to separate the question of modeling dark matter from the questions of electroweak and Higgs naturalness, at least for the immediate purpose

[5] While baryon number (lepton number) is broken by anomalies, just as in the SM, this need only imply baryon number violation via non-perturbatively small interactions, which can easily be well below any experimental bounds.

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of pursuing collider phenomenology. In the traditional view, every superpartner produced cascade decays down to the dark matter particle. But more generally, we can have R-parity violation and dark matter may or may not be at the end of superpartner decay chains.

6.4 RPV and FCNCs

A final reason for favoring R-parity is that in standard weak scale SUSY, large parts of RPV parameter space lead to excessive FCNCs, only exacerbating the usual SUSY Flavor Problem. However, this point is mitigated, though not completely evaded, in effective SUSY, because of the greatly reduced squark content, as discussed below. Again, this makes RPV a more motivated possibility in the effective SUSY context.

In the end, we think that both R-parity and RPV alternatives are plausible in the effective SUSY context, and make for very different phenomenological features and search strategies. Below we discuss RPV with proton decay protected by lepton number symmetry, and alternatively by baryon number conservation.

6.5 RPV with Lepton number conservation

The standard renormalizable RPV SUSY couplings preserving lepton number are of the superpotential form $W \propto \bar{U}_I \bar{D}_J \bar{D}_K$, with generational indices I, J, K. Such couplings give rise to a variety of RPV Yukawa couplings and (after SUSY breaking) RPV A-terms which can decisively affect superpartner decays and flavor physics. Here, we specialize to the most minimal particle content of effective SUSY, as discussed in section 4, with beyond-SM field content given by $\tilde{q}_L, \tilde{t}^c_R, \tilde{h}_{u,d}, \lambda_i$. While there is the up-type scalar singlet $\tilde{t}^c_R$, there is no down-type scalar singlet, and therefore no RPV A-terms are possible in the effective SUSY theory. The only RPV Yukawa couplings that come from truncating the above type of superpotential to effective SUSY are of the form

$$\mathcal{L}_{RPV} = \kappa_{IJ} \tilde{t}^c_R d^{cI}_R d^{cJ}_R . \tag{6.1}$$

We will consider this to be added to the minimal 10 TeV effective Lagrangian of eq. (4.2), or the 1 TeV effective Lagrangian of eq. (4.4).

Flavor constraints on these couplings, reviewed in ref. [42], easily allow RPV coupling strengths that lead to prompt squark decays into quarks at colliders. But while lepton-number conservation is sufficient to protect against proton decay (assuming the gravitino or other non-minimal fermions are heavier than the proton), it does not forbid neutron-antineutron oscillations. This is because (accidental) U(1) baryon-number symmetry is incompatible with the combination of RPV couplings, gaugino-squark-quark coupling, and Majorana gaugino masses. The bounds on neutron-antineutron oscillations are stringent (see [43] for review), even in effective SUSY where CKM suppressions are incurred in mediating such effects via the third generation squarks and gauginos. Again, RPV couplings can straightforwardly be strong enough to lead to prompt squark decays to quarks at colliders. And yet, they cannot be order one in strength. Theoretically, having RPV couplings $\ll 1$ is plausible enough, related perhaps to the smallness of ordinary Yukawa couplings. Experimentally, small RPV couplings imply that squarks cannot be *singly* produced at colliders.

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Remarkably, there is a way of recovering U(1) baryon number symmetry consistent with *order one* RPV couplings of the form of eq. (6.1), but it requires realizing gauginos as components of Dirac fermions. Observing single squark production can then be an interesting diagnostic of supersymmetry breaking, even those parts out of direct reach of the 7 TeV LHC. We will show how this works in section 7.

6.6 R-parity violation with Baryon number conservation

The standard renormalizable RPV SUSY couplings preserving baryon number are superpotential terms of the form, $W \sim LL\bar{E}, QL\bar{D}, LH_u$. Let us again consider truncating to the minimal beyond-SM field content described in section 4, $\tilde{q}_L, \tilde{t}^c_R, \tilde{h}_{u,d}, \lambda_i$. Again, there are no A-terms of the forms of these superpotentials possible, and the $LL\bar{E}$ completely vanishes. The bilinear superpotential turns into a mixing mass term $\ell\tilde{h}_u$. Since $\tilde{h}_d$ and the left-handed leptons, ℓ, share the same gauge quantum numbers, we can choose a new basis for them such that there are no $\ell\tilde{h}_u$ terms. The only surviving RPV Yukawa couplings are then of the form,

$$\mathcal{L}_{RPV} = \kappa'_{IJ} d^{cI}_R \ell^J_L \tilde{q}_L. \tag{6.2}$$

We defer the study of the flavor constraints and the LHC implications of this type of baryon-number conserving RPV interactions within effective SUSY to future work. Ref. [42] reviews such interactions in the more general SUSY context.

7 Dirac Gauginos

We have argued in the context of our 10 TeV effective SUSY theories that naturalness requires sub-TeV gluinos, which provides a very significant and visible SUSY production channel at the LHC. Yet, if we remain uncommitted to the structure of physics above 1 TeV, we have argued that the gluino need not be present in the sub-TeV effective theory. At first sight, these two statements might seems in conflict, but in fact they merely exemplify a general theme in SUSY models: a very minimal field content in the far IR often *requires* a less minimal field content at higher energies. This is the case with regard to gauginos, and gluinos in particular due to their stronger couplings. The idea of Dirac gauginos [44–46] is to have extra field content in the form of a chiral superfield, Φ_i, in the adjoint representation of each SM gauge group, with soft SUSY breaking such that the Φ_i fermion, χ_i, and the gaugino, λ_i, get a Dirac mass with each other, $m_{\lambda_i}\lambda_i\chi_i$. With such non-minimal field content below 10 TeV we will see that it is natural to have the Dirac gauginos heavier than 1 TeV.

The 10 TeV effective theory with Dirac gauginos, analogous to the construction of eq. (2.2), is given by

$$\begin{aligned}\mathcal{L}_{\text{eff}} = &\int d^4\theta K + \left(\int d^2\theta \left(\frac{1}{4}\mathcal{W}^2_\alpha + y_t\bar{T}H_uQ + y_b\bar{B}H_dQ + \mu H_uH_d + (\sqrt{2}m_i\theta^\alpha)\mathcal{W}_{i\alpha}\Phi_i\right) + \text{h.c.}\right)\\ &+\mathcal{L}^{light}_{kin} - \left(\bar{u}Y^{light}_u h_u q_L + \bar{d}Y^{light}_d h_d q_L + \text{h.c.}\right) + \mathcal{L}_{lepton}\\ &-m^2_{\tilde{q}_L}|\tilde{q}_L|^2 - m^2_{\tilde{t}^c_R}|\tilde{t}^c_R|^2 - m^2_{\tilde{b}^c_R}|\tilde{b}^c_R|^2 - m^2_{h_u}|h_u|^2 - m^2_{h_d}|h_d|^2 - m^2_{\phi_i}|\phi_i|^2\\ &-B\mu h_u h_d - A_t\tilde{t}^c_R h_u\tilde{q}_L - A_b\tilde{b}^c_R h_d\tilde{q}_L + \text{h.c.}\\ &+\mathcal{L}_{hard} + \mathcal{L}_{non-ren.},\end{aligned} \tag{7.1}$$

where the explicit Grassmann θ^α dependence parametrizes the soft SUSY breaking Dirac gaugino mass term in superspace notation, and m_ϕ^2 in the third line gives soft mass-squared to the scalars in the adjoint superfield Φ. The remaining terms are as discussed below eq. (2.2).

Similarly, the 10 TeV effective theory with Dirac gauginos, analogous to the construction of eq. (4.2), is given by

$$\begin{aligned}\mathcal{L}_{\text{eff}} = & \int d^4\theta K + \left(\int d^2\theta \left(\frac{1}{4}\mathcal{W}_\alpha^2 + y_t \bar{T} H_u Q + (\sqrt{2} m_i \theta^\alpha)\mathcal{W}_{i\alpha}\Phi_i\right) + \text{h.c.}\right) \\ & + \mathcal{L}_{kin} - \left(\bar{u} Y_u^{light} h_u q_L + y_b \bar{b} h_u^* q_L + \bar{d} Y_d^{light} h_u^* q_L + \text{h.c.}\right) + \mathcal{L}_{lepton} \\ & - m_{\tilde{q}_L}^2 |\tilde{q}_L|^2 - m_{\tilde{t}_R^c}^2 |\tilde{t}_R^c|^2 - m_{h_u}^2 |h_u|^2 - m_{\phi_i}^2 |\phi_i|^2 \\ & - \left(A \tilde{t}_R^c h_u \tilde{q}_L + m_{\tilde{h}} \tilde{h}_u \tilde{h}_d + \text{h.c.}\right) \\ & + \mathcal{L}_{hard} + \mathcal{L}_{non-ren.}.\end{aligned} \tag{7.2}$$

This scenario was first emphasized and studied in detail in the context of full supersymmetry in [26].

7.1 Naturalness

Expanding the soft gaugino mass term from superspace into components yields couplings,

$$\mathcal{L} \supset \sqrt{2} m_{\lambda_i} D^i (\phi_i + \bar{\phi}_i) - m_{\lambda_i} (\chi^i \lambda_i + \bar{\lambda}^i \bar{\chi}_i) \tag{7.3}$$

The D-term contributes mass to the *real* part of ϕ_i so that the total mass-squared is $m_{R_i}^2 = 2(m_{\lambda_i}^2 + m_{\phi_i}^2)$, while the imaginary part has mass-squared of just m_ϕ^2. In addition, the D-term generates a coupling of the real part of ϕ to the other scalars charged under the related gauge group. For the case of Dirac gluinos, we obtain the coupling $\mathcal{L} \supset -\sqrt{2} m_{\lambda_3} g_s (\phi_3^i + \bar{\phi}^a)(\bar{\tilde{q}} T^a \tilde{q})$, where T^a are the Gell-Mann color matrices. This provides a new correction to the stop mass-squared at one loop which cancels the logarithmic divergence found in eq. (2.11) [46]. Eq. (2.11) is then replaced by a UV-finite total correction,

$$\delta m_{\tilde{t}}^2 = \frac{2 g_s^2 m_{\tilde{g}}^2}{3\pi^2} \ln \frac{m_{R_3}}{m_{\tilde{g}}}. \tag{7.4}$$

Taking the stop much lighter than the gluino and the scalar gluon ("sgluon") to be comparable to the gluino mass (the above logarithm ~ 1), and requiring naturalness of the stop mass, yields

$$m_{\tilde{g}} \lesssim 4 m_{\tilde{t}}. \tag{7.5}$$

This implies it is natural to have gluinos above a TeV for stops as light as $\sim 300\,\text{GeV}$. In such cases, it is sensible to remove the gluino and sgluons from the sub-TeV effective theory, and from early LHC phenomenology.

Boson	q	Fermion	q
h_u	0	$\tilde{h}_u$	-1
		$\tilde{h}_d$	1
$\tilde{q}_L$	$\frac{4}{3}$	q_L	$\frac{1}{3}$
		(u_L, d_L), (c_L, s_L)	$\frac{1}{3}$
$\tilde{t}^c_R$	$\frac{2}{3}$	t^c_R	$-\frac{1}{3}$
		u^c_R, d^c_R, s^c_R, c^c_R, b^c_R	$-\frac{1}{3}$
		leptons	0
A_μ	0	λ	1
ϕ	0	χ	-1

Table 1. R-charges of particles in theory with eq. (6.1) and Dirac gaugino masses.

7.2 R-parity violation

As advertized in subsection 6.5, Dirac gauginos are also important for the case of lepton-number conserving RPV because they completely relax the stringent constraints from neutron-antineutron oscillations by allowing one to have a U(1) baryon number symmetry. The trick is that this symmetry is realized as an R-symmetry in the sense that different fields in a supermultiplet carry different charges. The charges of the fields are given in table 1. One can then check that eq. (7.2) and the RPV couplings of eq. (6.1) respect such a baryon number R-symmetry in the absence of the A term.

With baryon R-symmetry, neutron-antineutron oscillations are forbidden, even when RPV couplings are sizeable, which raises the possibility that stops can be singly produced at colliders.[6] But we first have to ask if this is plausible in light of flavor physics and CP constraints. A useful way to think of the new flavor structure of RPV couplings of $\tilde{t}^c_R$ in effective SUSY is that they effectively make this antisquark a “diquark”, even up to its baryon number. In this way, the general discussion and constraints of flavor structure for scalars with $d_R d_R$ diquark couplings given in [48] applies to the effective SUSY setting here. In particular, ref. [48] discusses the different plausible hierarchical structures for such couplings and the mechanisms underlying their safety from FCNC and CP-violating constraints. As is shown there, it is indeed plausible for the $\tilde{t}^c_R$ to have order one couplings to light quarks, and therefore be singly produced.[7]

Baryon-number R-symmetry, by forbidding the A-term, also makes for an interesting signature for pair-production of $\tilde{q}_L$ since they can no longer mix with $\tilde{t}^c_R$ after electroweak symmetry breaking. These squarks do not directly couple to quark pairs, unlike $\tilde{t}^c_R$, which means that each $\tilde{q}_L$ will decay into two third generation quarks *plus* a quark pair.

[6]Ref. [47] discusses a model in which it is $\tilde{b}_R$ that is singly produced (at the Tevatron), and in which neutron-antineutron oscillation placed important constraints. Dirac gauginos would also loosen these constraints in this context. (Our flavor estimates suggest that $\tilde{b}_R$ lighter than TeV is disfavored, but perhaps this is possible with a more special flavor structure.)

[7]A similar analysis is possible for (non-R-symmetry) baryon-number preserving RPV and loosening the constraints from lepton-number violation tests such as neutrinoless double-β decay.

7.3 Electric dipole moments

With the baryon R-symmetry as described above, it is straightforward to check that all the soft SUSY breaking parameters can be made real by appropriate rephasing of fields in eq. (7.2). Therefore there are no new CP-violating contributions to electric dipole moments from this Lagrangian. However, as discussed in section 5, we should more realistically add third-generation flavor-changing corrections to any such Lagrangian, which can contain new CP-violating phases. However, as discussed there these new terms will be suppressed by $\mathcal{O}(\epsilon^2)$. In this way, we expect non-vanishing but highly suppressed new contributions to EDMs. These observations for effective SUSY are closely related to the observations made in refs. [46, 49, 50].

8 Collider Phenomenology

In this section we will demonstrate three things:

1. After ~ 1/fb LHC running, there are analyses that put non-trivial constraints on the motivated parameter space of effective SUSY.
2. Nevertheless, very large parts of the parameter space, fully consistent with electroweak naturalness, are still alive.
3. The most constraining searches for effective SUSY, so far, are not always those optimized for more standard SUSY scenarios.

While effective SUSY has many interesting experimental regimes, we will not attempt a complete study in this paper. Rather, we will focus on the simplest natural setting, and do enough of the related phenomenology to make the points $(1-3)$ above. We will pursue the R-parity conserving scenario for a few related reasons. This naturally provides the phenomenological handle of sizeable missing energy, which can stand out in even the early LHC data. Secondly, there is greater familiarity in the community of pursuing these event topologies. Our results for effective SUSY can then be compared with the phenomenology of more standard SUSY scenarios. In the paper, we have however emphasized that R-parity violation is a particularly well-motivated option within effective SUSY, and it does display distinctive phenomenological features. We will pursue a more detailed study of this kind of the RPV phenomenology of effective SUSY in future work.

The central consideration for effective SUSY phenomenology is the great reduction in new colored particles, squarks, compared with standard SUSY scenarios. In effective SUSY we keep just the minimal set of superpartners below TeV needed to stabilize the electroweak hierarchy. This has the effect of lowering the new physics cross-sections substantially. Furthermore, in standard SUSY settings one typically entertains higher superpartner masses than is technically natural, partly a result of renormalization group running of superspectra from very high scales, and partly in order to radiatively raise the physical Higgs boson mass above the experimental bound. In our bottom-up effective SUSY, with less UV prejudice, we have only tried to constrain the spectrum from the viewpoint of naturalness and the little hierarchy problem. As we have seen, other mechanisms for raising the physical Higgs mass work well within effective SUSY. Therefore, we favor the regime where stops are lighter than 500 GeV, while gluinos may be so heavy as to be irrelevant in the early LHC. The decay products of lighter stops in effective SUSY can easily fail to pass the harsher cuts on missing energy and jet energies used in searches optimized for heavier superpartners.

In the following subsection, we will study in detail collider constraints which one can put on the most minimal scenario, namely light stops and sbottom (predominantly left-handed) with a neutralino at the bottom of the spectrum. We will briefly review the Tevatron constraints on this scenario and further analyze the constraints arising from LHC data at $\mathcal{L} \sim 1\,\mathrm{fb}^{-1}$. In the subsequent subsection we will survey other variations, but will not go into details. We leave this to future work.

8.1 Neutralino and Squarks

In this particular subsystem we will simplify considerations even further to the effective theory of eq. (4.5), where we have just a bino LSP lighter than the squarks. The neutralino might more generally be an admixture of several neutral gauge eigenstates, but phenomenologically this is not very relevant; the neutralino is simply a way of invisibly carrying off odd R-parity from colored superpartner decays. The bino is a good proxy for such a general neutralino. In the remainder of this section, we focus on the collider phenomenology of eq. (4.5).

One further simplification we make is to take the stops and sbottom to be roughly degenerate. If there is no substantial left-right mixing, this is a very good approximation in the left-handed (LH) sector. The mass difference between the LH stop and sbottom is given by

$$\Delta m \approx \frac{m_W^2 \sin^2\beta}{2m_{\tilde{q}_L}}. \tag{8.1}$$

Since this splitting comes from $\mathrm{SU}(2)\times\mathrm{U}(1)$ D-terms, it is proportional to the mass of the W. Usually if the splitting is dominated by D-terms, one gets that $m_{\tilde{b}} > m_{\tilde{t}}$. This might suggest that one should also consider a decay mode $\tilde{b} \to W^{(*)}\tilde{t}$. However this would imply a three-body decay, which is therefore highly suppressed. More important, stop decay modes $\tilde{t} \to W^{(*)}\tilde{b}$ can become competetive to other stop decay modes, if it is forced to proceed through an off-shell top. However this can happen only if the left-right mixing between the stops is large, and we will neglect this possibility further.

Before considering the LHC, we should note several D0 searches which directly address this scenario. The first relevant search looks for b-jets $+\not{E}_T$ [51]. This search constrains the sbottom mass to be higher than 247 GeV if the neutralino is massless. The constraints become weaker if the neutralino is heaver, but unless there is an accidental degeneracy, the lower bounds on the sbottom are still around 200 GeV. Another search of D0 looks for stops, which are pair-produced and further decay into $b\ l + \not{E}_T$ (where this decay mode is assumed to have 100 % branching fraction). The most updated search used events with opposite flavor pairs [52]. This search also bounds the stop mass at 240 GeV if the neutralino is massless and for massive neutralino (without any accidental degeneracy with the stop) the bound is of order 200 GeV, depending on the neutralino mass.

CDF has a more elaborate search, where it looks for $t\bar{t}+\not{E}_T$. This search was performed in monoleptonic [53] and hadronic [54] channels. The bounds one can put on production cross sections from these two measurements are comparable to each other, but too weak to constrain effective SUSY with its small squark cross section.[8]

Now let us turn our attention to the LHC searches. As we will see, the bounds from the LHC are not very stringent (partly due to an insufficient number of dedicated searches). This is in part because, with the exception of an Atlas top-group search for $t\bar{t}+\not{E}_T$ (which we will discuss later), there are no dedicated searches for this scenario. However there are several general searches, which can be sensitive to the stop/sbottom/neutralino subsystem we are studying here. We explictly considered the following list of searches:

1. jets + $\not{E}_T$ (including simple $\not{H}_T$ search and an α_T search) [56, 57]
2. jets + $\not{E}_T$ with b-tag [58, 59]
3. lepton + jets + $\not{E}_T$ [60]
4. OS dileptons + jets + $\not{E}_T$ [61]
5. lepton + jets with b-tag + $\not{E}_T$ [62]

In order to estimate the bounds on our scenario, we simulated events and checked the acceptances within the channels listed above.[9] The events were generated and decayed with `MadGraph 5` [66] and further showered and hadronized with `Pythia 6` [67]. The events were reconstructed with `FastJet-2.4.4` [68]. We calculated all the NLO cross-sections with `Prospino 2` [69] and reweighted all the events appropriately. We ran each spectrum assuming that the mass difference between the stops and sbottom are negligible. Given the mass difference, eq. (8.1), this is not a bad approximation. (One can of course play with the mass difference between $\tilde{t}_L$ and $\tilde{t}_R$, still keeping the spectrum natural, but we did not perform this study.)

We find that all the searches listed above, except searches for jets + $\not{E}_T$, do not put any interesting bounds on the subsystem that we are discussing here. The searches in leptonic modes put extremely harsh cuts on the H_T of the entire event, and therefore easily miss the stops in the range between 200 and 400 GeV, while the cross sections in the higher mass range are far too small. Unfortunately, the Atlas search for jets $+l+b-$tag$+\not{E}_{\mathrm{T}}$ [62] also does not add interesting constraints, mostly because it is tuned to detect (or exclude) gluinos above 400 GeV which further cascade-decay to bottom, top and neutralino.[10] The jets + $\not{E}_T$ searches indeed put interesting constraints on our stop/sbottom/neutralino

[8]Hereafter we do not consider a mass range of stop below 200 GeV, where the stop mostly decays off-shell. This intriguing possibility is not yet excluded, and the reader is refered to [16, 55].

[9]Whenever both Atlas and CMS have performed closely overlapping searches, we have considered just the CMS representative. The relevant Atlas searches are [63, 64]. We also did not explicitly simulate an additional CMS jets $+\not{E}_T$ search which takes advantage of the m_{t2} variable [65], since it is not expected to have a good acceptance in our case.

[10]This search claims that it looks for events with 4 b-jets with lepton and $\not{E}_T$, however demands only a single b-tag in the event selection. One can probably put more interesting bounds by demanding more than one b-tag.

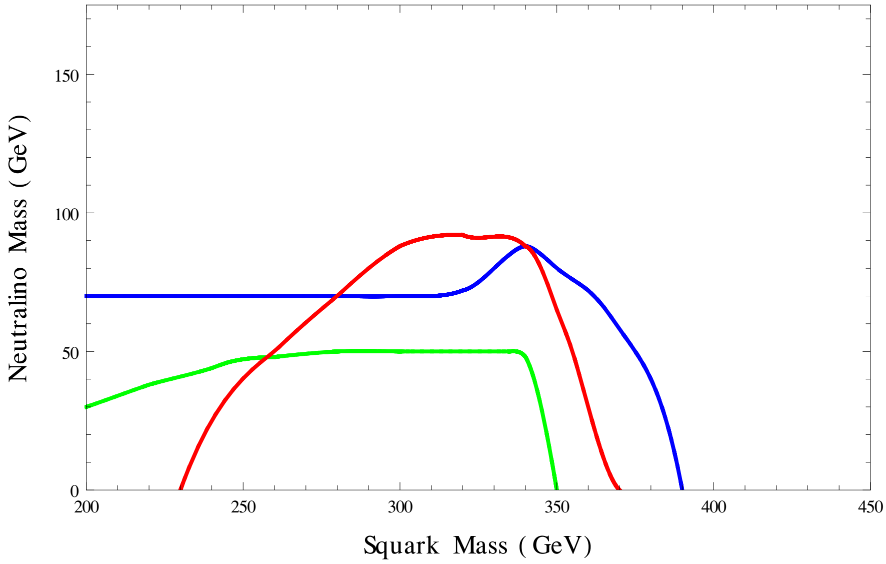

Figure 6. Exclusion curves for our minimal model, eq. (4.5), from three relevant searches as a function of masses for squarks and neutralino. We assume roughly equal masses for all three squark species, two stops and a sbottom. The green line represents exclusion by α_T search, the blue line is an exclusion by $\not{H}_T$ search and the red one is exclusion by $t\bar{t} + \not{E}_T$ search.

subsystem and we show our bounds in figure 6. We found that more than half of all the relevant events which contribute to the exclusion come from sbottom production and decays. In fact, even a single sbottom without any stops would be excluded all the way to 300 GeV with the same searches for massless neutralino. For more general neutralino mass the single-sbottom exclusion plot appears in figure 7. By comparison, the same searches put no bounds on a single stop (or even both stop species), due to extremely bad acceptance in this range of masses.

This, however, does not conclude the full list of searches. There is an additional search by Atlas, which looks precisely for $t\bar{t} + \not{E}_T$ in a monoleptonic channel [70]. This particular search puts almost no bound for production of a single species of stop, but the picture is different when we have both stops roughly degenerate (with double the production cross sections). We show the final exclusion plots on figure 6, where the exclusion due to $t\bar{t} + \not{E}_T$ search is given by the red curve. On figure 8 we show the ranges excluded by this search if we split the masses of the stops (neutralino mass is assumed to be zero). Note that this exclusion is comparable to the exclusion one gets with the jets $+\not{H}_T$ search.

8.2 Overview of some other possibilities

8.2.1 Gluinos

Because of their large color charge and the high multiplicity of their decay products, the biggest phenomenological consideration for the 7 TeV LHC is the presence or absence of gluinos below a TeV. Production cross-section grows significantly as gluinos are taken below 1 TeV in mass, and gluinos decay exclusively into the third generation squarks. This scenario has been studied both in cases when the gluino decays into a sbottom (see abovementioned searches for jets plus $\not{E}_T$ with a b-tag) or into a stop [62]. However,

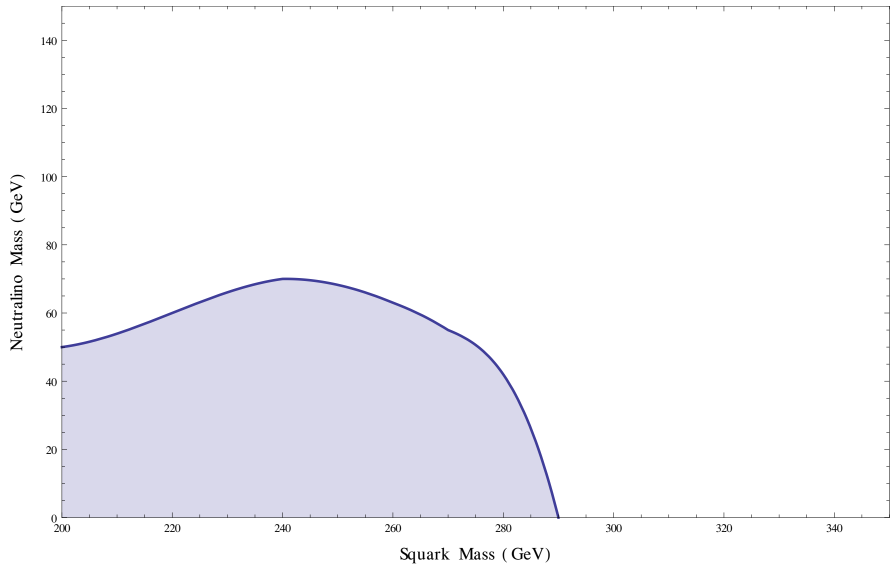

Figure 7. Exclusion of a single sbottom due to jets + $\not{H}_T$ search as a function of a sbottom and neutralino masses.

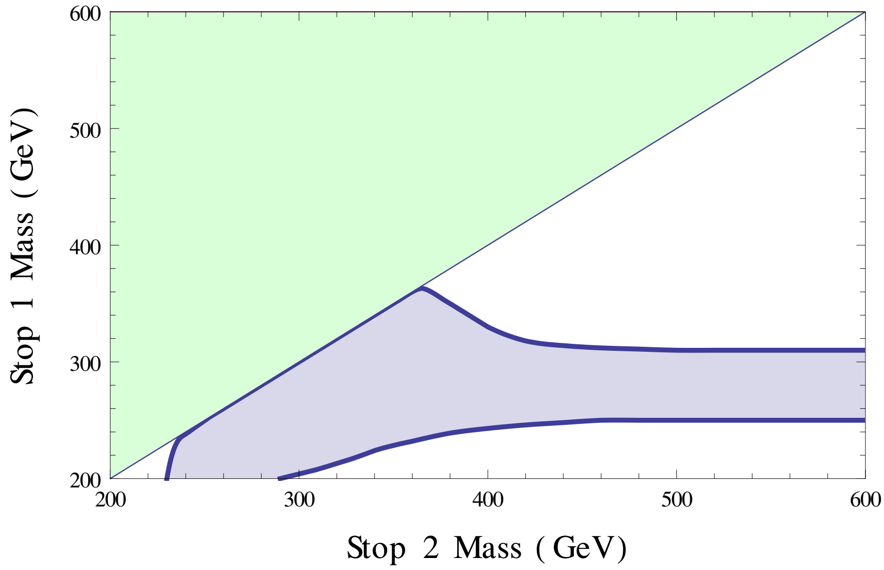

Figure 8. Exclusion curves for two stops with different masses from the Atlas search for $t\bar{t} + \not{E}_T$ in monoleptonic channel [70] . The neutralino mass is assumed to be zero. Note a narrow band between 250 and 290 GeV for the first stop which is excluded even when the second stop is very heavy. This is the region where the sensitivity of the search is maximized.

there are reasons to believe that a monoleptonic channel with one b-tag, which was used in the Atlas search is not optimal. The model of gluinos decaying exclusively to stops was carefully studied in [12] and it was found that with luminosity of $1\,\mathrm{fb}^{-1}$ gluinos up to 650 GeV can be discovered, if one takes advantage of a few competitive channels, like same-sign dileptons, multileptons with or without b-tags (and sometimes multiple b-tags).

8.2.2 Collider-stable squarks

One can also consider the very simple scenario with $\tilde{t}_L$, $\tilde{t}_R$ and $\tilde{b}_L$ at the bottom of the superpartner spectrum. With R-parity, the lightest scalar (either stop or sbottom) is stable. We should of course assume that it decays at some point (for example it can decay into a gravitino, or through some tiny R-parity violating coupling) in order to avoid constraints from searches for ultra-heavy hydrogen atoms [71], but this still allows squarks with cosmological lifetimes [72]. If this is the case, $\tilde{t}$ or $\tilde{b}$ should show up as R-hadrons at the LHC. Recent bounds from CMS impose severe constraints on this scenario if the lightest superpartner is a stop [73].[11] Results of these searches imply that a stable stop in the mass range between 100 and 800 GeV is excluded if its production cross section is of order 10^{-2} pb. Comparing these results to theoretically expected production cross sections [74], we find that these cross-sections are expected for a single stop with mass up to 600 GeV. However in our case, we should at least multiply the cross sections by a factor of three (we have two stops and at least one sbottom), rendering the bound to somewhat higher than 600 GeV. Therefore, if one takes the little hierarchy problem seriously up to ~ 10 TeV, this scenario is disfavored.[12]

8.2.3 Neutralino and Chargino LSPs

A safer option is to consider the effective theory of eq. (4.5), where we see the Higgsinos providing natural neutralino/chargino candidates. If the neutralino is the LSP, bounds on stable charged or colored particles are evaded. Of course, the neutralinos and charginos may more generally be an admixture of several electroweak gauge eigenstates.

In detail, the presence of a chargino as an NLSP makes a phenomenological difference, but we believe that it is less decisive in the present context. The difference from the scenario described in subsection 8.1 is that on top of the decay modes $\tilde{t} \to t\tilde{\chi}^0$ and $\tilde{b} \to b\tilde{\chi}^0$ we have already considered, we will have competing modes $\tilde{b} \to t\tilde{\chi}^\pm$ and $\tilde{t} \to b\tilde{\chi}^\pm$. Since we are mostly interested in the region of mass parameters where the top-quark mass is far from negligible, we conclude that the decay mode $\tilde{b} \to t\tilde{\chi}^\pm$ will be mostly suppressed due to the phase space. Therefore, introducing the chargino at the bottom of the spectrum will usually have a mild effect on sbottom decay modes and the constraints which come from these decays (mostly jets plus $\not{E}_T$). However the stops decay modes will be altered compared to our discussion in subsection 8.1, since the decay mode $\tilde{t} \to b\tilde{\chi}^\pm$ is now phase space unsuppressed. The chargino will consequently decay to the neutralino and W (maybe off-shell). Therefore, this will look roughly similar to the decay modes of a regular stop, even though the kinematics might be different. If the chargino and neutralino are quasi-degenerate, then the decay modes of stops very much resemble those of sbottoms, thereby effectively increasing the production cross sections for jets plus $\not{E}_T$ and making the constraints somewhat more stringent then what we find in subsection 8.1.

[11] Even though the authors of this paper do not interpret there results in terms of stable $\tilde{b}$, there is no reason to believe that this bound would be dramatically different.

[12] However, as noted in subsection 4.2, the effective theory of eq. (4.4) is a useful departure point for adding in RPV phenomenology.

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While the above are reasonable deductions, explicit simulation is still required when charginos are light. We again leave this to future work.

9 Outlook

In this paper, we developed a bottom-up formulation of effective supersymmetry and analysed some of its phenomenological aspects. As we have shown, the constraints on effective SUSY, even with the most conservative approach, R-parity with neutralino at the bottom of the spectrum, are very mild. With these assumptions, the data still allow a spectrum fully consistent with electroweak naturalness.

This conclusion strongly suggests the future research program in this direction. Evidently, current LHC searches are not optimized for this scenario. It would be interesting to see how one can increase the sensitivity of the current searches and vary the cuts so as to allow better acceptance for effective SUSY. We expect that there is a strong opportunity for searches optimized to effective SUSY to make great inroads into discovery or exclusions within ~ 10/fb of LHC running, in the coming year.

Another promising avenue one can take has to do with R-parity violation. As we emphasized in section 6, RPV is highly motivated if effective SUSY indeed describes the physics immediately beyond the SM. Even the signals of RPV SUSY with lepton-number violation can be quite challenging if squark decays into leptons involve τ. The signals of RPV SUSY with baryon-number violation are even more challenging, because the decays of the squarks will mostly results in jets. However, as pointed out for the case with baryon R-symmetry, squarks can have more spectacular decays into several jets, including two with heavy flavor. Current exotica searches [75] put very mild bounds on these RPV scenarios and it is very interesting if one can improve these search strategies to get better sensitivity to the new physics.

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