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Supersymmetry with a Sister Higgs*

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Abstract

Within the context of supersymmetric theories, explaining a 125 GeV Higgs motivates a consideration of a broader range of models. We consider a simple addition to the MSSM of a "Sister Higgs" (Σ_d), a Higgs field that participates in electroweak symmetry breaking but does not give any direct masses to Standard Model matter fields. While a relatively minor addition, the phenomenological implications can be important. Such a field can be naturally charged under an additional symmetry group G_s. If gauged, the Higgs mass is naturally much larger than in the MSSM through an NMSSM-type interaction, but with Σ_d playing the role of H_d. The addition of the sister Higgs allows new R-parity violating operators $\Sigma_d H_d E$, which are less constrained than conventional leptonic R-parity violation. Considerations of unification motivates the presence of colored G_s-charged fields. Production of these G-quarks can lead to new b-rich final states and modifications to decays of gluinos, as well as new opportunities for R-parity violation. Unlike a conventional fourth generation, G-quarks dominantly decay into a light jet and a scalar (potentially the Higgs), which then generally decays to b-jets. The presence of additional sister charges allows the possibilities that lightest sister-charged particle (LSiP) could be stable. We consider the possibility of an LSiP dark matter candidate and find it is generally very constrained.

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I. INTRODUCTION

As data continue to arrive from the LHC, the simplest models of supersymmetry have become increasingly constrained. In particular, the recent discovery of a Higgs-like feature at 125 GeV [1, 2] is tantalizingly close to the tree level prediction of the MSSM, but is large enough to require large radiative corrections, and pushes the MSSM into narrow corners of parameter space [3]. At the same time, no clear evidence for any of the other superpartners has appeared. Both of these results can be accommodated by a higher SUSY scale, but at the cost of increasing the tuning of electroweak symmetry breaking scale.

Alternatives for raising the Higgs mass have been varied: the simplest possibility maybe the NMSSM, in which a singlet S is added to the MSSM together with the superpotential term $\lambda S H_u H_d$. The resulting quartic coupling $|h_u h_d|^2$ raises the Higgs mass at small $\tan\beta$. However, aside from the aesthetic issues with including a complete singlet, achieving a large enough quartic requires a superpotential coupling which must be near the limit of perturbativity at the GUT scale. At the same time, if a much larger value for the quartic were possible, then improved possibilities for naturalness arise [3].

Large quartics can be natural, for instance if the singlet S and/or Higgs is a composite [4, 5] (alternatively the new states can be integrated out yielding quartics from Kähler potential terms [6]). New matter at high energies can boost the quartic somewhat by raising the standard model gauge couplings in the UV [7, 8].

While the D-term quartics are fixed by supersymmetry, they can be enhanced by non-decoupling D-terms from an additional gauge group [9, 10]. While appealing in principle, most models require significant additional matter content to maintain unification.

Extending the MSSM to include new fields is dangerous however, as they can introduce new sources of FCNCs or significant corrections to precision electroweak observables. There is a simple addition, however, that can evade these concerns: namely, the inclusion of additional Higgs multiplets, which we denote Σ_u and Σ_d. These fields carry the same SM quantum numbers as H_u and H_d, respectively. To avoid FCNCs these fields should have no tree-level, renormalizeable couplings to SM fermions. At the same time, there is no reason that they cannot participate in EWSB by acquiring significant vevs. Such fields we refer to as "sister" Higgs fields.

The inclusion of additional Higgs fields is certainly not new. What we are describing is

largely the generalization of the Type I 2HDM to a supersymmetric context. Similar non-coupling Higgses were employed to construct R-symmetric SUSY theories [11]. As we shall explore in this paper, however, there are important phenomenological consequences when allowing fields such as these in the low energy theory.

First, an important point is that for NMSSM-like enhancements of the Higgs mass, the coupling $SH_u\Sigma_d$ is just as effective as SH_uH_d, i.e., the object that acquires a large vev needn't be the object that gives the down-type fermions their masses. This allows a great deal of freedom in reconsidering the dynamics of the light fields in the theory (i.e., the charged Higgs) and whether they carry additional quantum numbers under some new group G_s. In particular, the charges of Σ_d are only tied to the charges of S, not to the charges of down-type fermions.

Secondly, the inclusion of the sister Higgs allows the presence of a new R-parity violating operator, namely $H_d\Sigma_dE$, and with it comes a great deal of new phenomenology. In the presence of R-parity conservation, new LSPs are possible, including LSPs that carry a conserved charge and thus can potentially provide a dark matter (DM) candidate.

Finally, the presence of these fields motivates the presence of additional down-type quarks, also charged under G_s. These G-quarks can have important consequences for signals such as gluino decays, and other opportunities for R-parity violation. Unlike usual 4th generations, these G-quarks decay via scalar (often Higgs) emission, leading to b-rich final states.

Any one of these points motivates a consideration of these theories with extended Higgs sectors. We describe the construction of a sister MSSM in more detail below. In the rest of the paper we discuss the extension to a gauged sister model, before addressing each of the above mentioned points in more detail.

A. A Sister Higgs

A sister Higgs is a Higgs field that participates in EWSB but has no tree level couplings to matter. This is a natural outcome when the sister Higgs transforms as a non-trivial representation of some new symmetry G_s under which the SM fields are trivial. This may be a gauged or global symmetry. As we shall see shortly, although in its simplest manifestation G_s is not gauged, it is likely most interesting in a case when G_s is a gauge symmetry.

To make this more concrete: we extend the MSSM by introducing two sister Higgs fields,

	$SU(3)\times SU(2)\times U(1)_Y$	G_s $[SU(2)_s]$
Σ_u	$(1,2,1/2)$	$\mathbf{r}$ [2]
D_g	$(\mathbf{3},1,-1/3)$	$\mathbf{r}$ [2]
Σ_d	$(1,2,-1/2)$	$\bar{\mathbf{r}}$ [2]
D_g^c	$(\bar{\mathbf{3}},1,1/3)$	$\bar{\mathbf{r}}$ [2]
Φ	$(1,1,0)$	$\mathbf{r}$ [2]
$\bar{\Phi}$	$(1,1,0)$	$\bar{\mathbf{r}}$ [2]
H_u	$(1,2,1/2)$	1 [1]
H_d	$(1,2,-1/2)$	1 [1]

TABLE I: Charges of fields under the SM and the Sister group, G_s, for the generic case [the case of $G_s = SU(2)$]. Together Σ_u and D_g (Σ_d and D_g^c) make a $\mathbf{5}$ ($\bar{\mathbf{5}}$) under $SU(5)$. Note the usual MSSM fields Q, U^c, D^c, L, and E^c are singlets under G_s.

Σ_u and Σ_d, which carry the same charge under the SM as H_u and H_d respectively, while also being vector-like under the new sister group, G_s. At this stage we do not specify whether G_s is a global or gauged symmetry. In order to allow couplings between the sister fields and the MSSM we also introduce fields that only carry charge under G_s. Furthermore, if these fields acquire a non-zero vacuum expectation value (vev) then there will be mass mixing between the sister Higgses and the Higgses of the MSSM. We denote these G_s breaking fields as $\Phi(\bar{\Phi})$, which are in the same representation of G_s as $\Sigma_u(\Sigma_d)$. In addition to these fields, as we will see, it is natural to expect vector-like fields that transform as right-handed quarks and are also charged under G_s, which we dub G-quarks, D_g, D_g^c. Later, in Section II, we will investigate a particularly interesting case where the sister group is a gauged $SU(2)$. The charges of all fields in both the general and this special case are shown in Table I. With this field content and assuming $\mathbf{r} \neq \bar{\mathbf{r}}$ the simplest extension of the MSSM superpotential is

$$\begin{aligned} W \;=\; & \mu_\phi \Phi\bar{\Phi} + \mu_\Sigma \Sigma_u \Sigma_d + \mu_h H_u H_d + \mu_g \bar{D}_g D_g \\ & + \lambda_u \Phi H_u \Sigma_d + \lambda_d \bar{\Phi} \Sigma_u H_d \\ & + Y_u U^c H_u Q + Y_d D^c H_d Q + Y_\ell E^c H_d L \ . \end{aligned} \tag{1}$$

To make this completely general, we can add a term $\bar{\Phi} D_g D^c$ that we shall discuss shortly. Notice that if the representation $\mathbf{r}$ is different from the conjugate $\bar{\mathbf{r}}$, this superpotential

contains an additional $U(1)$ symmetry, which needn't be accidental, which we denote $U(1)_\Sigma$. Under this symmetry Σ_u, Φ have charges $1/2$ and $\Sigma_d, \bar{\Phi}$ have charges $-1/2$. In the interesting case where $G_s = SU(2)$, $\mathbf{r}$ is pseudo-real, and operators $\Phi\Sigma_u H_d + \bar{\Phi}\Sigma_d H_u$ are also allowed, for instance, which break the $U(1)_\Sigma$. Although the concept of a Sister Higgs is independent of G_s, if this group is too large (if the representations of Σ are larger than 4 dimensional) then the SM gauge couplings will reach a Landau pole below the GUT scale.

The most natural origin for SM fermion masses is through the usual Yukawa couplings to H_u and H_d as it is in the MSSM, and we have illustrated this in (1). However, with the inclusion of new fields acquiring vevs, it is possible that down-type fermion masses can instead come from the higher dimension operator $\Phi\Sigma_d Q D^c$. This could be generated by integrating out a massive $\mathbf{5} + \bar{\mathbf{5}}$ pair, which are uncharged under G_s, and have couplings $\bar{\mathbf{5}} Q D^c + \mathbf{5}\Phi\Sigma_d$. We will not focus on this case, but it may be natural to consider extending this model by the presence of additional $\mathbf{5} + \bar{\mathbf{5}}$'s.

II. GAUGED SISTERS AND THE HIGGS MASS

As we have discussed, there has been great activity attempting to accommodate a Higgs at 125 GeV in supersymmetric theories. Such a large Higgs mass in the MSSM requires squark masses in the multi-TeV range and large A-terms, well beyond our expectations from naturalness [12].

Various extensions of the MSSM attempt to alleviate this problem, as we have discussed. While the simplest possibility may simply be of a tuning of the MSSM, the simplest potentially natural model is likely the NMSSM. As is well known [8, 13], the RG flow of the SH_uH_d coupling in the NMSSM tends to make it run quite small in the IR. Unlike the top quark, which has $SU(3)$ color contributing to its anomalous dimension, there is no strong interaction to drive this coupling large. Consequently, a 125 GeV Higgs lies near the edge of perturbativity without appealing to sizable radiative corrections from top loops. This leads us naturally to the possibility of gauging the sister group. If the sister group runs strong then the λ_u and λ_d couplings, which play the same role as SH_uH_d in the NMSSM, can remain large and give a sizeable correction to the Higgs self coupling and therefore the Higgs mass. However, such an approach has a "why-now" problem: the coincidence of the strong coupling scale of the new gauge group and the weak scale. If instead the running

of the gauge coupling of the sister group, g_S, were approximately conformal this why-now problem would be removed. An example of such a setup – where the beta function is zero at one loop – is $G_s = SU(2)_s$ with[1] six vector-like families of matter.

With $G_s = SU(2)_s$ and Φ, $\bar{\Phi}$, Σ_u, Σ_d all in fundamental representations of $SU(2)_s$ a one-loop conformal beta function occurs if there are 3 more $(\square, \overline{\square})$ pairs, the G-quarks. This is exactly the field content shown in Table I and is the same as what would be expected if the sister Higgses were embedded in a $(\mathbf{5}, \overline{\mathbf{5}})$ of $SU(5)$. Remarkably, simply by insisting on conformality as a solution to the "why-now" problem, we are naturally pushed into a unified (i.e. GUT) setup with the G-quarks. While the gauge coupling for $SU(2)_s$ is conformal at one-loop, the SM values are modified by the addition of two additional flavors, i.e., $\beta_3 = -1, \beta_2 = 3, \beta_1 = 43/5$.

The low energy phenomenology of the sister Higgs scenario can be much more complicated than the usual MSSM (or NMSSM). In general, there are as many as six "Higgs" fields (i.e., $SU(2)_W \otimes U(1)_Y$ doublets) as well as four SM singlet scalars, in addition to a Z' and the associated fermionic fields. However, for the most part, we can understand the phenomenology by taking a simplified limit of the theory. As we will see this limit reduces the theory to a Type I 2HDM [14]. [2]

We can take this limit as follows: let us begin by considering the superpotential in (1). This superpotential retains a global $U(1)_s = U(1)_\Sigma + T_s^3$ symmetry after $SU(2)_s$ breaking and an associated "sister charge". Within the $SU(2)_s$ doublet Σ_d, for instance, only one component mixes with the Higgs, while the other component, carrying the sister charge, will not. When writing the scalar potential, we will designate Σ_d^s with the ^{s} to indicate when the field carries sister charge, and similarly with the Φ, $\bar{\Phi}$ and Σ_u. In discussing the Higgs sector, we will for the moment assume the presence of this symmetry as it simplifies the physics by reducing by a factor of two the number of fields that mix with the Higgs boson.

Since Σ_d will take the role of the H_d for the purposes of raising the Higgs mass, λ_d is not important for the Higgs mass phenomenology, so for simplicity we set it to zero. We

[1] This group is also small enough that the additional matter charged under the SM does not lead to Landau poles in the SM gauge couplings.

[2] In a Type I 2HDM, there is a standard model like Higgs H_1 which couples to fermions and a second Higgs H_2 which does not, although it can acquire a vev. A mixing between them induces couplings for both mass eigenstates.

take $SU(2)_s$ to be broken by a $\langle\bar{\Phi}\rangle = v_\phi \cos\beta_s$ and $\langle\Phi\rangle = v_\phi \sin\beta_s$. In general, we will be taking $\tan\beta_s \ll 1$, so the $\bar{\phi}$ is acquiring the larger vev. For the moment, we then set $\beta_s = 0$. We will assume that the ϕ fields have large enough SUSY breaking masses that they can be integrated out without disturbing the phenomenology (i.e., the mass spectrum and the $\lambda_u^2|H_u\Sigma_d|^2$ quartic) and set them to their vevs. Σ_u is for the most part a spectator, so we integrate it out with a SUSY breaking mass.

The resulting theory, with these assumptions has a vev-acquiring sector composed of Σ_d, H_d and H_u, with the scalar potential,

$$\begin{aligned} V \;=\; & m_{H_d}^2|H_d|^2 + m_{H_u}^2|H_u|^2 + (m_{\Sigma_d}^2 + \frac{m_{Z_s}^2}{2})|\Sigma_d|^2 + (m_{\Sigma_d}^2 - \frac{m_{Z_s}^2}{2})|\Sigma_d^s|^2 \\ & + |\lambda_u|^2|H_u\Sigma_d|^2 + (B\mu H_u H_d + \lambda_u\mu_\phi^*\langle\bar{\Phi}\rangle^* H_u\Sigma_d + \text{h.c.}) + \frac{1}{2}D_{MSSM}^2 + \frac{1}{2}D_s^2 \ . \end{aligned} \tag{2}$$

Here, we have taken the fermions to couple to the scalar sector through the usual Yukawa matrices $Y_u Q U^c H_u$ and $Y_d Q D^c H_d$ and D_{MSSM} are the usual D-term contributions from $SU(2)_W$ and $U(1)_Y$, and D_s are the residual $SU(2)_s$ D-term quartics after integrating out the $\phi, \bar{\phi}$.

We can further integrate out H_d in the large m_{Hd}^2 limit by $H_d = B\mu H_u^*/m_{Hd}^2$, which yields the final simplified setup.

$$\begin{aligned} V \;=\; & m_{H_u}^2|H_u|^2 + (m_{\Sigma_d}^2 + \frac{m_{Z_s}^2}{2})|\Sigma_d|^2 + (m_{\Sigma_d}^2 - \frac{m_{Z_s}^2}{2})|\Sigma_d^s|^2 \\ & + |\lambda_u|^2|H_u\Sigma_d|^2 + (\lambda_u\mu_\phi^*\langle\bar{\Phi}\rangle^* H_u\Sigma_d + \text{h.c.}) + \frac{1}{2}D_{MSSM}^2 + \frac{1}{2}D_s^2 \ , \end{aligned} \tag{3}$$

where the fermion masses now arise in the effective theory via the terms $Y_u H_u Q U^c$ and $Y_d \frac{B\mu}{m_{Hd}^2} H_u^* Q D^c$. Put simply, the low energy effective theory is a Type-I 2HDM with an additional spectator Σ_d^s. Although Σ_d^s does not mix with these fields nor acquires a vev, we have left it in because of the relationship between its mass and that of the $U(1)_s$ neutral Σ_d (which can lead to tensions in some regions of parameter space). From here, the scalar phenomenology is fairly simple. H_u acquires a vev $\langle H_u\rangle = v\sin\beta$, while $\langle\Sigma_d\rangle = v\cos\beta$, and the Higgs quartic has an NMSSM like contribution proportional to λ_u^2.

The NMSSM-like boost to the Higgs mass will depend on how large λ_u can be. The β function for λ_u at one loop is

$$\beta_{\lambda_u} = \lambda_u\left(3y_t^2 + 5\lambda_u^2 - \frac{3}{5}g_1^2 - 3g_2^2 - 3g_s^2\right), \tag{4}$$

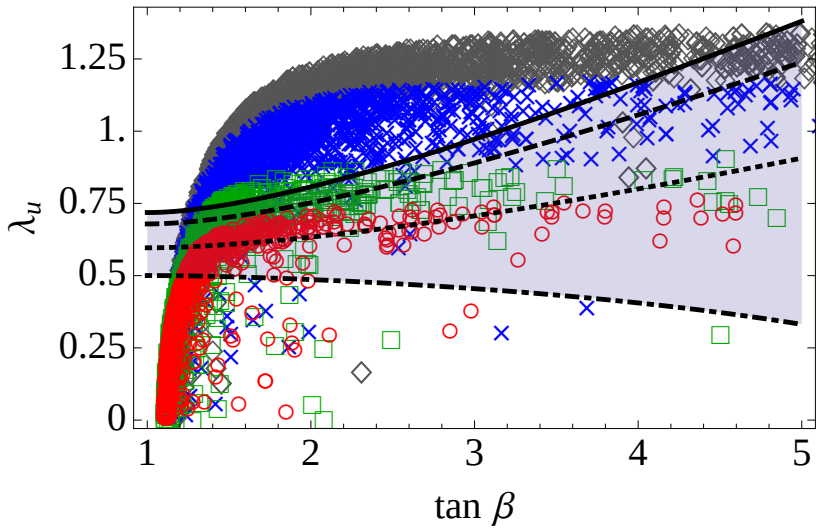

FIG. 1: IR values of $\tan\beta$ and λ_u when scanning over the range [0,4] for the UV inputs g_S, λ_u, and y_t. Different ranges for $g_S(1\,\text{TeV})$ are denoted by different symbols, ○: $g_S < 0.5$, □: $0.5 \leq g_S < 1$, ×: $1 \leq g_S < 1.5$, and ◇: $g_S \geq 1.5$. The lines show the parameters for which a higgs of 125 GeV is achieved at tree level (solid), with stops of 250 GeV (dashed), 500 GeV (dotted) and 1 TeV (dot-dashed).

where $g_{1,2,3}$ are the standard model couplings, and g_s is the coupling of the new sister group. These expressions make clear the simple point: the presence of g_s acts to counter the tendency of λ_u (and similarly for λ_d) to run small in the IR. A simple question, therefore, is how large can the Higgs mass be in this theory, consistent with perturbativity up to the GUT scale? Because some couplings are sizeable, it is best to calculate the RG flow to low energies at two-loop. To do this, we use the publicly available `SARAH` package [15, 16], to evolve the couplings within our theory from $M_{GUT} = 2 \times 10^{16}$ GeV to 1 TeV. We generate a large set of random high scale values of g_s, y_t, λ_u, all lie in the range $[0, 4]$, and run down to low energies. We define $\sin\beta = \frac{m_t(1\,\text{TeV})}{v\ y_t(1\,\text{TeV})}$, where $v \equiv 174$ GeV and we take $m_t(1\,\text{TeV}) = 150.7\,\text{GeV}$ [17]. Points with $\sin\beta > 1$ are disregarded as in conflict with observation.

The results of our scan are shown in Figure 1. While one must be careful with this plot (as there is no sense of a measure on this scatter plot) we can get a sense of correlations between parameters. First, we see that it is quite easy to have very large values of λ_u at low energies. Taking $\alpha_s(M_{GUT}) \gtrsim 0.18$, i.e., $g_s(M_{GUT}) > 1.5$, we can have $\lambda_u \simeq 1.25$ at low energies. This is sufficiently large that no additional correction from stop loops is

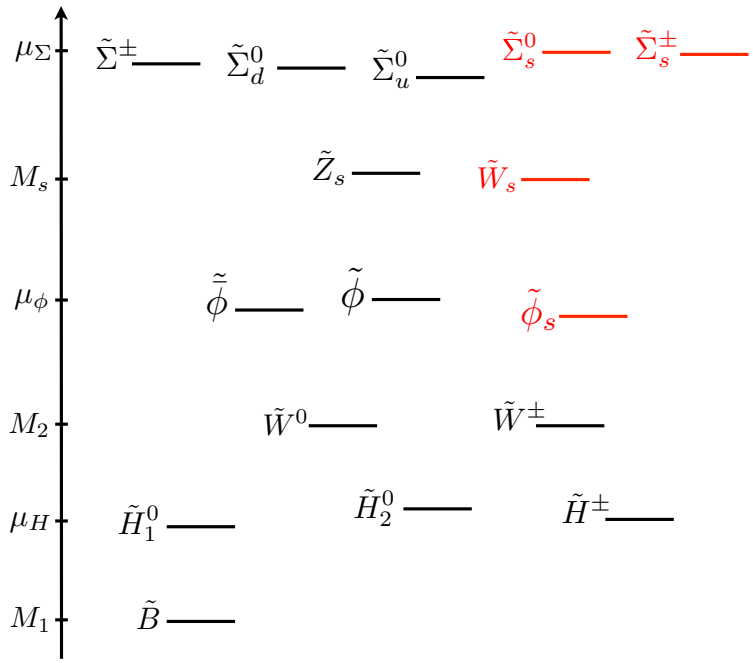

FIG. 2: Example spectrum of fermions, in our simplified limit. States in red are charged under the global $U(1)_s$. No special ordering of mass scales is implied.

necessary. However, including even moderate stop contributions is easy. With $\alpha_s = .02$ (i.e., comparable to the $SU(2)_W$ coupling) λ_u is easily large enough to get $m_H = 125\,\text{GeV}$ with only $\sim 500\,\text{GeV}$ stops.

A final point on which we make a small comment, is that Φ plays the role of the usual S in the NMSSM and, as recently pointed out by [3], for large λ_u the theory can become more natural, in that the soft Higgs mass can be larger, and then mixing the SM singlet scalar (in this case ϕ) with the Higgs, its physical mass can be lowered at the expense of just a moderate tuning. Although we do not explore this quantitatively here, it is clear that the ingredients (a large λ_u and a neutral scalar mixing with the Higgs) are present.

A. New Phenomenology in the Simplified Limit

Even in this limit where we have decoupled $\phi, \bar{\phi}$ and Σ_u, there are very important differences from the MSSM. First, because the second Higgs field in the low energy theory, Σ_d, has no direct couplings to SM fermions, the mass constraints that apply to $H^\pm$ in Type-II 2HDM do not directly apply. In particular, the charged Higgs can be $\lesssim 150$GeV without

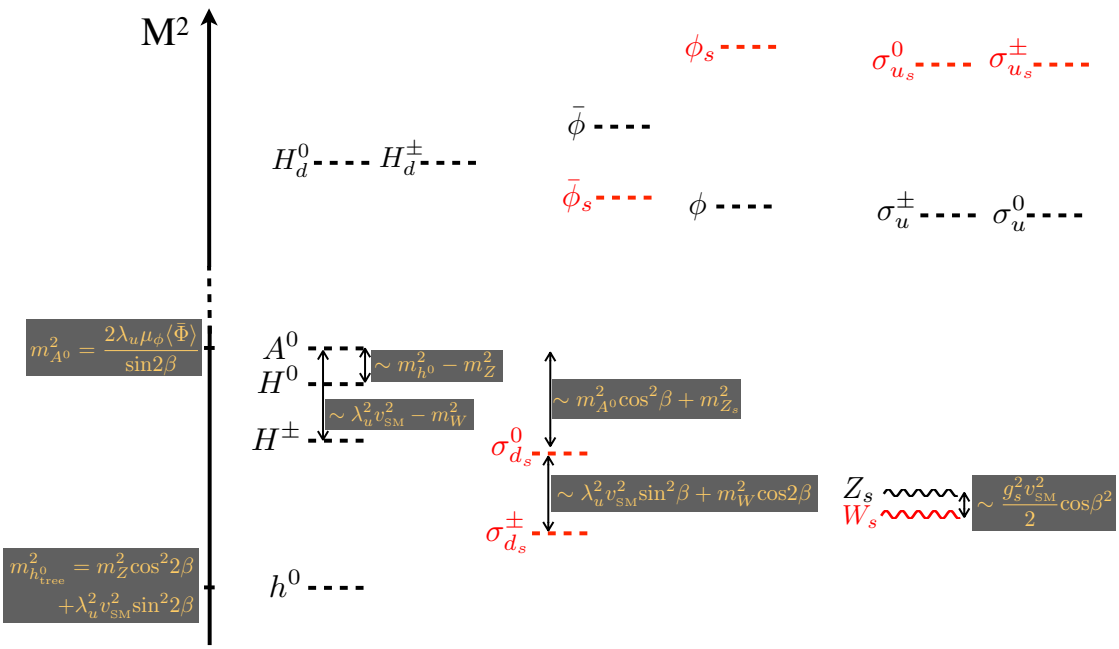

FIG. 3: Example spectrum of scalars and vector bosons, in our simplified limit. States in red are charged under the global $U(1)_s$.

necessarily contributing too largely to the decay of the top, nor excessively to $b \to s\gamma$ [18].

However, unlike a pure Type-I 2HDM, with $G_s = SU(2)$, the sister Higgs is tied to a second field, namely its sister-charged partner Σ_d^s. In particular, the mass of the charged Higgs $H^\pm$ (which is dominantly $\Sigma_d^\pm$) is related to the mass of the sister-charged $\Sigma_d^{s\pm}$ by $m^2_{H^\pm} - m^2_{\Sigma_d^{s\pm}} \sim m^2_{Z_s}$ (assuming $\langle\Phi\rangle \ll \langle\bar{\Phi}\rangle$). With larger $\langle\Phi\rangle$, this becomes $m^2_{H^\pm} - m^2_{\Sigma_d^{s\pm}} \sim (\cos 2\beta_s + 2\lambda_u^2 \sin^2\beta_s/g_s^2)\, m^2_{Z_s}$ (assuming $\tan\beta \gg 1$). Either way, for heavy enough Z_s and light enough $H^\pm$, this will induce a charge-breaking negative squared mass to $\Sigma_d^{s\pm}$.

Probably the simplest way out is to break the global sister charge $U(1)_s$ by including an operator such as $\lambda_u^{\not s}\bar{\Phi}H_u\Sigma_d$. The F-term for H_u then gives an additional contribution to the mass of Σ_d^s, avoiding the unwanted charge-breaking negative $m^2_{\Sigma_d^{s\pm}}$. This operator, combined with $\lambda_u\Phi H_u\Sigma_d$, will also contribute to the squared-mass of H_u through the F-term for Σ_d. In order to avoid fine-tuning, $\langle\Phi\rangle$ and $\langle\bar{\Phi}\rangle$ should not be too large, implying that Z_s should not be too heavy relative to the EW scale. Moreover, including this sister-charge breaking term will affect the running of λ_u, but given that large values of λ_u are already generally quite natural, this is unlikely to be problematic.

Another key element of this setup is the presence of new gauge forces, Z_s^0 and W_s^0. Because Σ_d is charged under both $SU(2)_s$ and $SU(2)_{\rm W}$, its vev will induce mixing between Z_s^0 and the

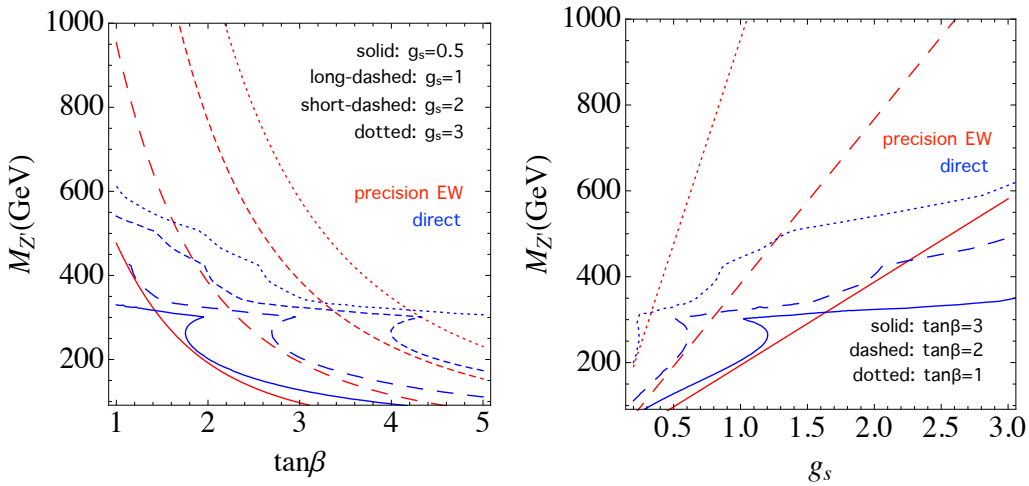

FIG. 4: Bounds on $Z_s - Z$ mixing expressed in terms of m_{Z_s}, $\tan\beta$ and g_s, from precision electroweak constraints of the ρ-parameter (red) and direct collider searches for dilepton resonances [19] (blue). The later are derived assuming that Z_s has the same branching ratios to leptons as the SM Z^0. The region below each line is excluded.

SM Z^0. Such mixing is constrained by precision electroweak (PEW) measurements of the ρ-parameter (although, as has been pointed out previously, e.g., by [20], this can actually help agreement with the standard model) as well as collider searches for Z' resonances. ATLAS performed a search for high-mass dilepton resonances in [19]. We can conservatively re-cast the results of that search in terms of M_{Z_s}, g_s and $\tan\beta$, assuming that Z_s decays as a sequential Z' (*i.e.* assuming that its branching ratios to SM fermions is identical to that of the SM Z^0). These bounds, as well as the PEW constraints, are shown in Fig.4. Note that even for large g_s and moderate $\tan\beta$, Z_s can be relatively light. Finally, we point out that the direct bounds from dilepton resonances are significantly weakened if the dominant decay mode of Z_s is $Z_s \to Z^0 h^0, W^\pm H^\mp$ or even $\phi^*\phi$, which is naturally the case if such on-shell final states are kinematically allowed and $\tan\beta$ is not too large. Despite this, the PEW constraints are stronger and hold regardless of how Z_s decays.

An interesting consequence of $Z_s \to Z^0 h^0$ being the dominant decay mode of Z_s is that it provides a sizable contribution to $Z^0 h^0$ production at the Tevatron. Saturating the precision electroweak constraint on $Z - Z_s$ mixing, the Z_s production can be comparable to that of SM $Z^0 h^0$ production at the Tevatron for relatively light Z_s, as shown in Fig.5. At the LHC such effect is much smaller, however.

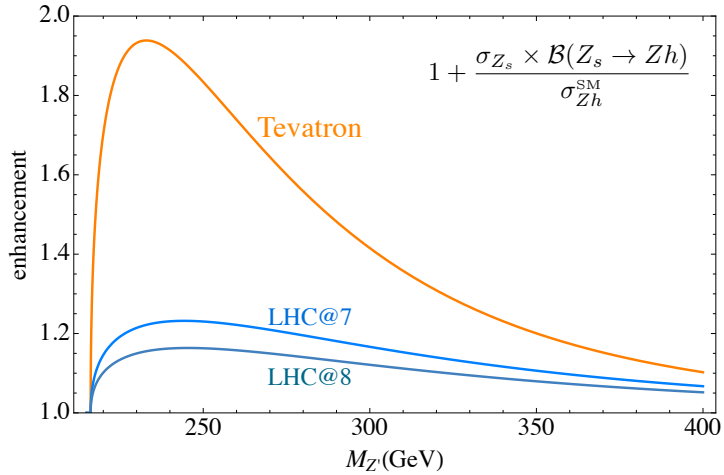

FIG. 5: Saturating the bounds for the ρ-parameter, there can be a sizeable contribution to Z^0h^0 production at the Tevatron through resonant Z_s production. At the LHC, it is a small correction to the SM process.

Finally, we comment on the implications of the Sister Higgs framework on the phenomenology of dark matter. As was pointed out in [21], in the phenomenologically viable region of MSSM parameter space, achieving the correct relic abundance for LSP dark matter requires considerably tuning between parameters. There are several viable regions where the so-called "well tempered" neutralino lives, with just the right admixture of bino/wino or bino/higgsino to allow the LSP to be light ($m_\chi \sim \mathcal{O}(100)$ GeV) and have the correct relic abundance. In the case of the Sister MSSM there are many new neutralino states that can mix with those of the MSSM; in general the neutralino mass matrix is now 13×13. Many of these new states carry charge under G_s and have new interactions, either through the superpotential or, if G_s is gauged, Z_s/W_s exchange. The new gauge bosons can be light and have sizeable gauge couplings (Fig. 4), and so have considerable effect on the process of thermal DM freezeout. The possibility of Dirac dark matter charged under the residual sister $U(1)_s$ will be explored at length shortly, our point here is simply that the neutralino sector is much richer, which allows for a broad range of possibilities for DM.

III. R-PARITY VIOLATION AND A SISTER HIGGS

With the addition of the new fields charged under G_s, the symmetry structure quickly grows complex. As a consequence, some of the new fields may be stable. At the same time,

	$U(1)_B$	$U(1)_L$	$U(1)_\Sigma$	$U(1)_G$
Q, U^c, D^c	$\pm 1/3$	0	0	0
L, E^c	0	± 1	0	0
Σ_u	0	0	$1/2$	0
D_g	0	0	0	1
Σ_d	0	0	$-1/2$	0
D_g^c	0	0	0	-1
Φ	0	0	$1/2$	0
$\bar{\Phi}$	0	0	$-1/2$	0

TABLE II: Charges of fields under the global symmetries. The non-MSSM fields are also charged under a gauged G_s.

we shall see that an intriguing new opportunity for R-parity violation appears in this theory. Let us begin by considering *only* the superpotential of eq.(1). This superpotential has the gauge symmetry group $SM \otimes SU(2)_s$, but has a global symmetry $U(1)_B \otimes U(1)_L \otimes U(1)_\Sigma \otimes U(1)_G \times P_R$. B and L are the usual baryon and lepton numbers, under which only the matter fields of the MSSM are charged, the new fields charged under $SU(2)_s$ are taken to be neutral. G is a conserved quantum number for the G-quarks, P_R is the usual R-parity. With this superpotential, $\Phi, \bar{\Phi}, \Sigma_u$ and Σ_d can be all even or all odd under P_R, and similarly the parity of D_g is unassigned. However, to avoid spontaneous breaking of R-parity, the components of H, Σ, and Φ scalars which acquire a vev must be even under R-parity. Finally, $U(1)_\Sigma$ is a global symmetry under which $\bar{\Phi}, \Sigma_d$ have charges $-1/2$, while Φ, Σ_u have charges $+1/2$.

When $SU(2)_s$ is broken by Φ and/or $\bar{\Phi}$ vevs, the $SU(2)_s \otimes U(1)_\Sigma$ is broken to a global $U(1)_s$, which we refer to as a sister charge. Under this residual symmetry, one component of Φ is charged, while the component getting a vev is neutral, and similarly for $\bar{\Phi}$. Likewise, one $SU(2)_L$ doublet of Σ_d mixes with H_d and is neutral under $U(1)_s$, while another is charged and does not mix, and again with Σ_u and H_u. Thus, in this theory, with no additional operators, the lightest baryon, lepton, R-parity odd field, sister-charged field and G-quark are all stable. The charges of the fields are summarized in Table II.

The presence of additional operators can break the symmetries further. In particular it is possible to break R-parity in several ways. Below we discuss a particularly interesting

case, with novel phenomenology, that breaks R-parity in a new way in the lepton sector, but before doing so we discuss an alternative approach, with less novel phenomenology, of R-parity breaking in the hadron sector.

A. Hadronic RPV

It is simple to break $U(1)_G$, for instance by including a small mass-mixing term[3]

$$\bar{\Phi} D_g D^c. \tag{5}$$

This operator has charge under B, G and Σ, and its inclusion breaks one linear combination of these symmetries. The two unbroken symmetries can be taken to be the combinations $B' = B + G/3$ and $\Sigma' = \Sigma + G/2$. [4] Under this modified baryon symmetry, B', the field D_g carries charge 1/3. This will allow G-quarks to decay, assuming that there is some lighter particle with the residual $U(1)_s$ charge (such as Σ or Φ).

Another operator which affects the $U(1)_G$ is $D_g^c D_g^c U^c$. This operator is charged under $U(1)_G$ and $U(1)_B$, and thus breaks a combination. We can think of the remaining symmetry as a modified baryon number, $B'' = B - G/6$, with this assignment D_g has new baryon number $-1/6$. Note that the G-quarks *cannot* decay with just this operator, since there are no states lighter than them with $B'' = \pm 1/6$. Equivalently, it retains a G-quark parity. Including this operator together with (5) breaks $U(1)_G \otimes U(1)_\Sigma \otimes U(1)_B$ down to a single $U(1)$. We can think of the residual number as a baryon number, under which D^c, D_g, Φ have charges of $-1/3, -1/6$ and $-1/2$.

When $\bar{\Phi}$ acquires a vev it breaks $SU(2)_s \otimes U(1)_\Sigma$ down to $U(1)_s$, which we can now identify as baryon number, $B''' = B'' - (T_s^3 + \Sigma')$. That is, different components of D_g under the $SU(2)_s$ transformation have different baryon numbers. Simply put, since the operator is $D_g^{c1} D_g^{c2} U^c$ (with 1,2 indexing the $SU(2)_s$ component), if we set D_g^{c1} as the field that mass mixes, then we can treat D_g^{c2} as a diquark, and the residual sister symmetry as baryon number. Since, in a similar fashion to the lepton RPV of Section III, we can redefine baryon

[3] We include this rather than $\Phi D_g D^c$ at least initially, as D_g^c is considered to be in a multiplet with Σ_d and thus to have the same G_s representation. Assuming the presence of even an approximate $U(1)_\Sigma$, $\bar{\Phi} D_g D^c$ should be more significant than $\Phi D_g D^c$.

[4] Note that, alternatively, we could have included the operator $Q D_g^c \Sigma_d$, which would have the same effect, although it breaks the three charges down to a different linear combination.

number and therefore R-parity to include the sister charge, this operator does *not* allow the LSP to decay as a consequence.

We can achieve R-parity violation in this fashion if we include the operator $\Phi D_g D^c$, along with $\bar{\Phi} D_g D^c + D_g^c D_g^c U^c$, which explicitly breaks $U(1)_\Sigma$. In essence, this mass mixes both components of the G-quarks, so that $D_g^c D_g^c U^c$, in the mass eigenstate basis, yields the usual UDD RPV operator. Alternatively, if there is an additional $\mathbf{5} + \bar{\mathbf{5}}$ uncharged under $SU(2)_s$, we could integrate it out and leave the operator

$$\frac{1}{M_{5\bar{5}}} \bar{\Phi} D_g^c D^c U^c. \tag{6}$$

In the presence of (5), this acts as the usual baryon number violating operator UDD. Interestingly, this preserves $U(1)_s$ and thus the lightest sister particle (or LSiP) remains a potential dark matter candidate.

Finally, we comment that including *both* terms such as $\Sigma_d \Sigma_d E^c$, discussed in section III B, as well as $\bar{\Phi} D_g D^c + D_g^c D_g^c U^c$ can lead to $B+L$ violating (but $B-L$ conserving and R-parity conserving) operators such as $U D^c D_g^c E H_d \Sigma_d$. Since these operators preserve $B-L$, they will not lead to proton decay, but could lead to neutron decay inside a nucleus. While the constraints on such operators are weaker than with $B-L$ violation, they are clearly constrained, and we leaved detailed studies for this hybrid scenario for future study.

B. Leptonic RPV

Within the MSSM, the operator $H_d H_d E^c$ is never discussed, and for good reason - since the $SU(2)_W$ indices of H_d are contracted antisymmetrically, it vanishes identically. However, as soon as there is an additional flavor of H_d, the operator $H_d^1 H_d^2 E^c$ can be present. This sister-Higgs R-parity violation is qualitatively different from usual RPV in the MSSM, as we shall discuss. Let us first see how this arises in the more complicated scenario where the sister Higgs is an $SU(2)_s$ doublet.

We can naturally include the superpotential operator

$$\kappa_i \Sigma_d \Sigma_d E_i^c \ . \tag{7}$$

This breaks $U(1)_\Sigma \otimes U(1)_L$ down to $U(1)_{L'}$ where $L' = L - \Sigma$ and under which Σ_d carries lepton number 1/2. In addition this operator determines a unique R-parity assignment

within the Σ fields that leaves R-parity preserved. When Φ acquires a vev, we can think of the residual sister charge as identified with lepton number. Hence, this operator does not break R-parity without some additional effect. More explicitly, there is a new, preserved, R-parity we can define, $R' = (-1)^{3(B-Q_s)+2\mathbf{s}}$, where $Q_s = T_s^3 + L'$ and $\mathbf{s}$ is the spin of the particle. The simplest modification that leads to R-parity violation is to include the terms $\bar{\Phi} H_u \Sigma_d$ and/or $\Phi H_d \Sigma_u$ in the superpotential. These explicitly break $U(1)_\Sigma$, which means there is no residual $U(1)_s$, and thus both components of Σ_d will ultimately mix with Higgs bosons. There are no stable states in the theory that may act as dark matter.

Alternatively, if again we consider a $\mathbf{5} + \bar{\mathbf{5}}$ uncharged under $SU(2)_s$, we could integrate it out and leave the operator

$$\frac{\kappa_i'}{M_{5\bar{5}}} \Phi \Sigma_d H_d E_i^c \ . \tag{8}$$

This operator preserves $U(1)_\Sigma$, and hence $U(1)_s$, but it is not possible to define a preserved R-parity in the presence of (1) and (8). This operator breaks R-parity (and lepton number), and thus we are left with an interesting possibility: that R-parity is broken, but that a residual global symmetry is intact, leaving a potential candidate for dark matter. In this case the collider production of DM will either be through electroweak processes, or through the decay of the G-quarks which have QCD production, but may be very heavy and so kinematically suppressed.

In either of these two cases, stable DM candidate or not, there are interesting implications for collider phenomenology. R-parity violation (RPV) will remove the usual $\not{E}_T$ signature of SUSY, but, unlike hadronic RPV, where the $\not{E}_T$ is converted to jets (and can thus be extremely challenging to discover), here, the $\not{E}_T$ is converted to leptons+jets, which is much more tractable. Furthermore, under the assumption that some flavor structure dictates the sizes of couplings, the final state lepton is essentially all τ. Then the final state will contain non-negligible $\not{E}_T$ and the SUSY signatures, while more challenging, are not necessarily impossible. If the dominant decays are to e, μ then the $\not{E}_T$ requirement of many new physics searches may not be met and the constraints will be weaker. We discuss below the possible signatures for decays for various LSPs in the presence of this form of leptonic RPV, and outline which existing searches may be sensitive to these signatures. We leave a detailed study of efficiencies and resulting bounds for future work [22].

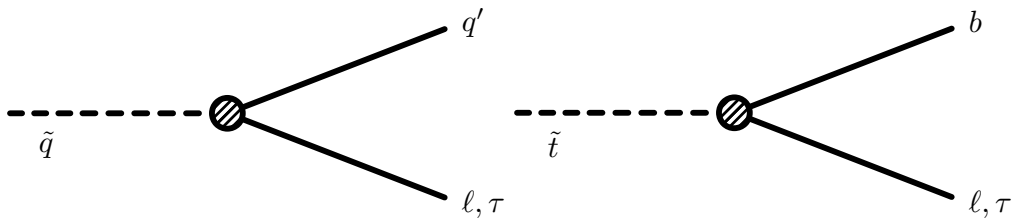

FIG. 6: Squark and stop decay modes in the leptonic sister RPV scenario.

C. Constraints on Sister RPV

Both (7) and (8) will lead to mixing between the RH leptons and the Higgsinos. We concentrate for concreteness on the effects of (7) with one or both of $\Phi H_d \Sigma_u + \bar{\Phi} H_u \Sigma_d$. One linear combination of leptons, the combination $\sum \kappa_i \ell_i$, will have a mass contribution $\sum \kappa_i^2 v^2 \cos^2\beta / M_\Sigma$. Unlike the conventional LH_u and LLE^c RPV operators these contributions do not affect the left-handed leptons, meaning that constraints from neutrino mass are much weaker. Likewise, since there is only a single lepton appearing in this operator, constraints on dilepton resonances are not expected to be relevant. [5]

Rare decays are expected, but because this operator couples to right-handed leptons only, the transition will have a chiral suppression, i.e., we expect a dipole operator for $\ell_i \to \ell_f \gamma$ ($\ell_i = \mu, \tau$, $\ell_f = \mu, e$) of $D_{if} \sim m_\ell \kappa_i \kappa_f e / 16\pi^2 M_\Sigma^2$. Current bounds on the lepton flavor violating branching ratios place constraints on the κ_i of,

$$|\kappa_\mu \kappa_e| \lesssim 3 \times 10^{-6} \left(\frac{M_\Sigma}{100\,\text{GeV}}\right)^2 \quad \text{and} \quad |\kappa_\tau \kappa_\ell| \lesssim 2 \times 10^{-4} \left(\frac{M_\Sigma}{100\,\text{GeV}}\right)^2 . \tag{9}$$

Squark LSP

A squark LSP in the presence of sister RPV is an interesting case to consider. Given that the operator $H_d^1 H_d^2 E_i^c$ mixes the chargino with a charged lepton, the dominant decay mode of a squark LSP may be $\tilde{q} \to q' \ell^\pm$ (Fig. 6), similar to the decay induced by the RPV operator QLD^c in the MSSM. Note, however, that contrary to QLD^c, the sister RPV operator $H_d^1 H_d^2 E_i^c$ does not induce the decay $\tilde{q} \to q \nu_\ell$ at tree-level.

If the squarks belong to the first or second generation and dominantly decay to $\tilde{q} \to q' e^\pm$ or $\tilde{q} \to q' \mu^\pm$, straightforward bounds on $m_{\tilde{q}}$ can be directly extracted from leptoquark

[5] The exception is if the LSP is a sneutrino, which can decay to two leptons.

searches. ATLAS carried out searches for first and second generation leptoquarks with 1 fb^{-1} of luminosity [23, 24], and constrained the masses of leptoquarks to be $m_{LQ} > 685$ GeV for $LQ \to \mu j$ and $m_{LQ} > 660$ GeV for $LQ \to ej$. CMS also searched for second generation leptoquarks with 2 fb^{-1} of luminosity [25], and placed the bound $m_{LQ} > 632$ GeV for $LQ \to \mu j$.

The bounds above on leptoquark masses directly apply to squarks decaying to those same final states, since the production cross section for squarks and leptoquarks are effectively the same in the limit in which the gluino is heavy and has a negligible contribution to squark pair production. However, in generic scenarios in which the first and second generation of right and left handed squarks are degenerate, the squark pair production cross section is a factor of 8 larger than leptoquark pair production. It is straightforward to extract the limit on squark masses in this case: the ATLAS search bounds the production cross section of $\tilde{q} \to q' e^{\pm}$ to be $\sigma(pp \to \tilde{q}\tilde{q}^* \to e^+e^- jj) \leq 5 \times 10^{-3}$ pb, implying that $m_{\tilde{q}} \gtrsim 880$ GeV. Similarly, the strongest constraint on $\tilde{q} \to q' \mu^{\pm}$ comes from CMS. It bounds the production cross section to be $\sigma(pp \to \tilde{q}\tilde{q}^* \to \mu^+\mu^- jj) \leq 3 \times 10^{-3}$ pb, which translates into $m_{\tilde{q}} \gtrsim 900$ GeV.

Those searches cannot, however, be as straightforwardly applied to the case in which the dominant squark decay mode is $\tilde{q} \to q' \tau^{\pm}$. In that case searches specifically targeting new physics signatures with jets and two tau leptons place stronger bounds. ATLAS performed a search for events with jets, at least two tau leptons and missing energy using 2 fb^{-1} of data [26]. We re-interpreted the results of those searches in order to place bounds on first and second generation squarks decaying as $\tilde{q} \to q' \tau^{\pm}$, as well as stop LSPs decaying as $\tilde{t} \to b\tau^+$. We used `MadGraph5` to simulate $pp \to \tilde{q}\tilde{q}^*$ and $pp \to \tilde{t}\tilde{t}^*$ at the 7 TeV LHC. Those events were passed through `Pythia6` for decay, showering and hadronization, and then handed to `PGS4` for a rough simulation of the detector response.

The ATLAS search in [26] selected events with at least two jets, the leading jet with transverse momentum $p_T > 130$ GeV and the sub-leading jet with transverse momentum $p_T > 30$ GeV. In addition, it required two hadronic tau candidates with $p_T > 20$ GeV and missing transverse energy $\not{E}_T > 130$ GeV. It also placed a cut $\Delta\phi(\vec{\not{E}}_T, \vec{p}_{T1,2}) > 0.4$ between the missing transverse momentum and the two leading jets, and $m_{\rm eff} > 700$ GeV. In Figure 7 we display the distribution of events in terms of the sum of the transverse mass of the two tau candidates, for a few benchmark signals. The signal region chosen by ATLAS corresponds to $m_T^{\tau_1} + m_T^{\tau_2} > 80$ GeV, and it bounds the number of signal events in that region to be

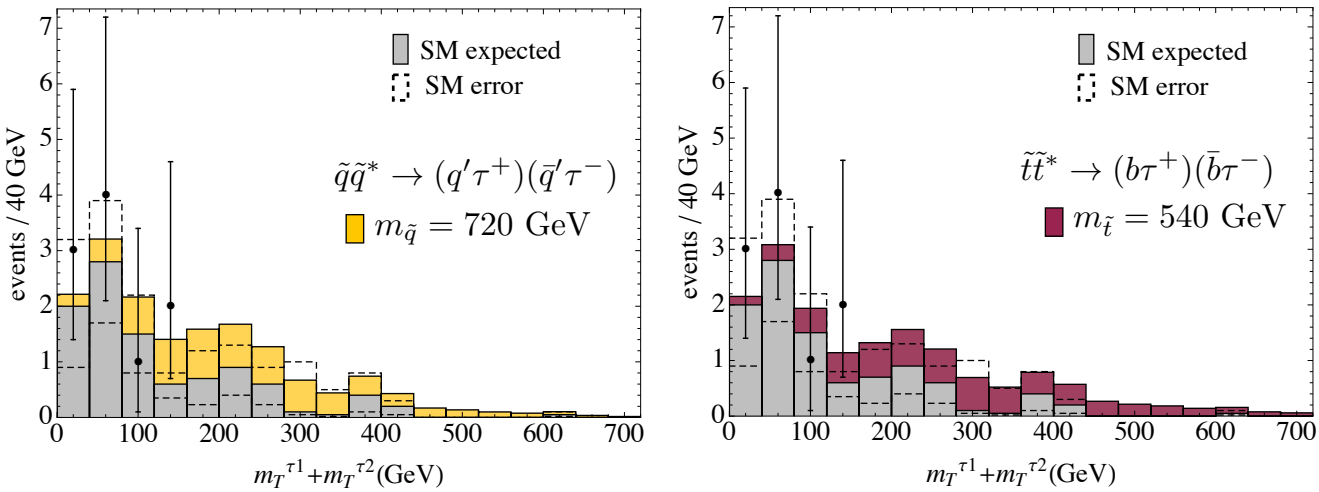

FIG. 7: Distribution of $m_T^{\tau_1} + m_T^{\tau_2}$ for the event selection of the ATLAS search in [26]. The SM background, as estimated in [26], is displayed in gray, and the measured data is shown as black dots. Two benchmark signals with masses at the boundary of the this analysis' reach are displayed as well: (left) degenerate 1st and 2nd generation squark LSP's decaying as $\tilde{q} \to q'\tau^\pm$, and (right) stop LSP decaying as $\tilde{t} \to b\tau^+$.

$\lesssim 6$ events. For the scenario of degenerate first and second generation squarks decaying as $\tilde{q} \to q'\tau^\pm$, that excludes squark masses $m_{\tilde{q}} \lesssim 720$ GeV. For the stop scenario with decay $\tilde{t} \to b\tau^+$ the exclusion is $m_{\tilde{t}} \lesssim 540$ GeV. This bound is surprisingly similar to the one obtained by the CMS dedicated search to this signature [27], which constrains stop masses to be $m_{\tilde{t}} \lesssim 525$ GeV. Searches for charged higgs from stop decays, $t \to bH^+ \to b(\tau^+\nu)$, such as [28], also yield non-trivial bounds. We estimate that the bound from this particular search to be $m_{\tilde{t}} \lesssim 300$ GeV, hence weaker than the bound from [26, 27].

Neutralino LSP

In the presence of the leptonic sister RPV operator $H_d^1 H_d^2 E_i^c$, a neutralino LSP can decay to $\ell^\pm H^\mp$. As we discussed before, the phenomenology of the light Higgs sector is that of a type-I 2HDM, and therefore $H^\pm$ will decay dominantly to $t\bar{b}$ (or $Wb\bar{b}$), $\tau\nu$ and $c\bar{s}$, with branching ratios determined by its mass and whether it is allowed to be on-shell or not. If $H^\pm$ is too heavy, the decay $\chi_1^0 \to \ell^\pm W^\mp$ might dominate. Either way, the expected signatures from neutralinos at the bottom of SUSY cascades should contain a large number of leptons.

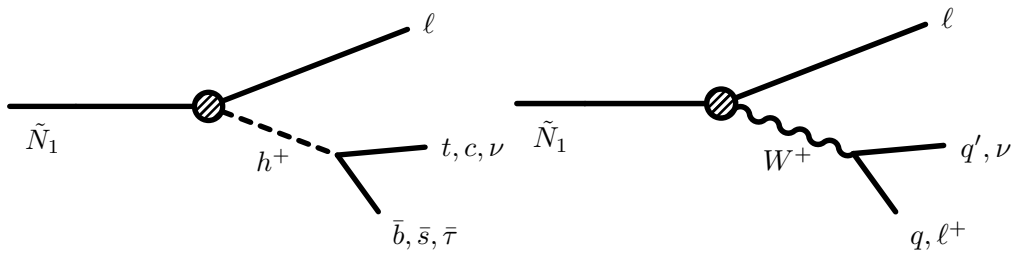

FIG. 8: Neutralino LSP decay modes induced by leptonic sister RPV.

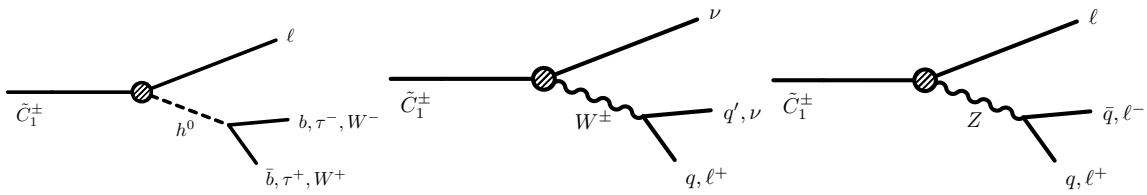

FIG. 9: Chargino LSP decay modes induced by leptonic sister RPV.

If these neutralinos are being produced through strong interactions (such as gluinos or squarks), these scenarios should be significantly constrained by same-sign dileptons [29–33] and multilepton searches [34–39], which have very low backgrounds and hence have very powerful sensitivity even for models with very low branching ratio to multileptons. Other potentially constraining searches are same-sign tops (if the decay $\chi_1^0 \to \ell^\mp H^\pm \to \ell^\mp (tb)^\pm$ is dominant), first and second generation leptoquarks [23–25], or GMSB with LFV [40].

Chargino LSP

It is also interesting to consider the situation where the chargino is the LSP (see [41] for a discussion of scenarios), or when its decay to a neutralino is suppressed due to phase space. In this case, the decay $\chi^\pm \to \ell^\pm h_x^0$ will be quite challenging, see Figure 9. If h_x couples to up quarks, it is possible that we would get a sizeable signal of WW (through its mixing with the Higgs), but if unavailable, then $c\bar{c}$ would dominate. Or, if coupling through down quarks, $b\bar{b}$, which can perhaps aid by adding b-tags along with hard leptons (possibly τs). Alternative decays to $W^\pm\nu$ or $\ell^\pm Z^0$ should be constrained constrained from standard SUSY searches in jets $+\not{E}_T$ and/or multileptons.

Slepton/Sneutrino LSP

Slepton or sneutrino LSPs decay in a similar fashion as in standard leptonic RPV induced by the LLE operator, see Fig.10. Searches for direct slepton production have been recently released by ATLAS [42] and constrain sleptons decaying as $\tilde{\ell}^{\pm} \to \ell^{\pm}\nu$ (where $\ell^{\pm} = e^{\pm}$ or $\mu^{\pm}$) to be $m_{\tilde{\ell}^{\pm}} \gtrsim 185$ GeV. These searches, however, cannot constrain sneutrino pair-production due to the absence of missing energy. However, multi-lepton searches requiring 4 leptons should have sensitivity to $\tilde{\nu}\tilde{\nu} \to 4\ell$.

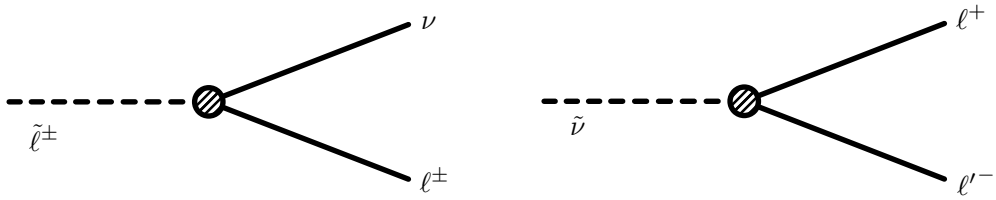

FIG. 10: Slepton and Sneutrino LSP decays induced by leptonic sister RPV

Gluino LSP

Gluino LSPs will decay to dijets plus a charged lepton. If the charged lepton is a $\tau^{\pm}$, standard jets+$\not{E}_T$ plus 0, 1 or 2 charged leptons should have sensitivity to pair-produced gluinos. For gluinos decaying to jets plus $e^{\pm}$ or $\mu^{\pm}$, however, the absence of missing energy makes this signature more challenging. Limits from current searches on leptoquarks, black-holes or same-sign dileptons may be re-cast for $\tilde{g}\tilde{g} \to (qq\ell^{\pm})(qq\ell^{\pm})$; however dedicated searches may have a potentially much higher reach.

IV. DARK MATTER

One appealing element of SUSY is the presence of a dark matter candidate in the form of the LSP. Under the assumption of R-parity conservation, the neutralino often serves as a WIMP candidate. In this theory, because of the sister Higgsinos, the state forming the WIMP is more complicated, with the new forces altering the DM freezeout story. While a thorough study of dark matter is beyond our scope, some discussion is clearly worthwhile.

A. Dark Matter with R-Parity Violation

In general, there is a tradeoff in models between the presence of RPV (which can be an explanation of the so-far elusive nature of supersymmetry) and a dark matter candidate. The standard lore is that since the LSP is the dark matter candidate, violating R-parity (and thus removing the $\not{E}_T$ signal) will remove a cosmologically stable particle from the theory. However, within the sister MSSM, that lore is easily violated by the presence of a residual global symmetry.

In the Sister Higgs scenario, that symmetry can be the global $U(1)_s$, which, if preserved, makes the lightest state charged under this symmetry (the "LSiP") stable and hence a potential dark matter candidate. If the $U(1)_s$ is broken down to a Z_2, then the LSiP is Majorana. As we shall see, the Dirac case is quite constrained by direct detection, but the Majorana case is much less so.

Among the fermions carrying $U(1)_s$ charge, there are three Dirac states comprised of $\tilde{\phi}_s \equiv (\tilde{\Phi}_1, \tilde{\bar{\Phi}}_2)$, $\tilde{\Sigma}_s \equiv (\tilde{\Sigma}_{u1}, \tilde{\Sigma}_{d2})$, and $\tilde{W}_s \equiv (\tilde{W}_{s^+}, \tilde{W}_{s^-})$. The masses and mixings of those fields can be extracted from:

$$\mathcal{L} \supset [\tilde{\Phi}_1 \;\; \tilde{\Sigma}_{u1} \;\; \tilde{W}_{s^+}] \begin{bmatrix} -\mu_\phi & -\lambda_u v s_\beta & g_s\langle\Phi_2\rangle \\ 0 & \mu_\Sigma & 0 \\ g_s\langle\bar{\Phi}_1\rangle & g_s v c_\beta & -M_s \end{bmatrix} \begin{bmatrix} \tilde{\bar{\Phi}}_2 \\ \tilde{\Sigma}_{d2} \\ \tilde{W}_{s^-} \end{bmatrix}. \tag{10}$$

Those fields couple vectorially to the Standard Model Z^0 boson, either directly though their Σ_s-components, or through $Z^0 - Z^0_s$ mixing. If the lightest mass eigenstate is the dark matter, its coupling to the Z^0 is constrained by direct detection experiments. The LSiP-*nucleon* cross section induced by Z^0 t-channel exchange is:

$$\sigma_n \simeq \frac{1}{4\pi}\left(\kappa_{\mathrm{LSiP}-Z}\frac{g_2}{c_w}\frac{1}{m_Z^2}\right)^2 \mu_n^2 \frac{(A-Z)^2}{A^2}, \tag{11}$$

where $\kappa_{\mathrm{LSiP}-Z}$ is the dark matter coupling to Z^0, μ_n is the dark matter-nucleon reduced mass and A (Z) is the nucleus atomic mass (number). This cross section is constrained by XENON100 [43] to be $\sigma_n \lesssim 10^{-44}\mathrm{cm}^2$. This translates into a bound on the dark matter-Z^0 coupling of $\kappa_{\mathrm{LSiP}-Z} \lesssim 3.7 \times 10^{-4}$, which rules out $\tilde{\Sigma}_s$ as a Dirac dark matter candidate since it is charged under $SU(2)_W$ and its coupling to Z^0 ($\kappa_{\tilde{\Sigma}_s - Z^0} = g_2/2c_w$) violates the direct detection bound by three orders of magnitude. This bound also severely constrains $\tilde{\phi}_s$ as a Dirac dark matter candidate, since it mixes significantly with $\tilde{\Sigma}_s$ through $\lambda_u\langle H_u\rangle\Sigma_d\Phi$.

These bounds can be evaded however, by adding small Majorana mass terms, for instance:

$$\mathcal{L} \supset \delta m_1 \tilde{\Phi}_1 \tilde{\Phi}_1 + \delta m_2 \tilde{\bar{\Phi}}_2 \tilde{\bar{\Phi}}_2, \tag{12}$$

or analogous terms for $\tilde{\Sigma}_s$. This term breaks $U(1)_s$ but preserves a Z_2 parity that stabilizes the dark matter. It also splits the mass eigenstates by $\delta m_1 + \delta m_2$. Since those states are Majorana, their Z-mediated elastic scattering is spin-dependent, and moreover suppressed by $\delta m/\mu_\phi$. The spin-independent scattering is inelastic and hence can be suppressed or even shut off for large enough mass splittings [44]. Note that it is non-trivial to achieve this term as superpotential couplings $\Phi\Phi + \bar{\Phi}\bar{\Phi}$ vanish identically by the antisymmetry of the $SU(2)_s$ indices. Thus, generating this term would naturally involve SUSY breaking or higher dimension operators. We shall not pursue this except to note that this is not a simple model addition.

In the absence of such $U(1)_s$-breaking majorana masses, however, the only viable Dirac dark matter is a state that is almost purely sister gaugino, $\tilde{W}_s$. The direct detection bound, in this case, translates into constraints on:

- the mixing of $\tilde{W}_s$ with $\tilde{\Sigma}_s$:

$$\frac{1}{2}\frac{g_2}{c_w}\,\frac{\theta^2_{\tilde{\Sigma}_{u1}\tilde{W}_{s+}} + \theta^2_{\tilde{\Sigma}_{d2}\tilde{W}_{s-}}}{2} < 3.7 \times 10^{-4}, \tag{13}$$

- the mixing of the Standard Model Z^0 with the sister gauge boson Z^0_s:

$$g_s \times \theta_{Z-Z_s} = \frac{g_s^2}{\sqrt{g_2^2 + g_Y^2}}\frac{m_Z^2}{m_{Z_s}^2 - m_Z^2}\cos^2\beta \;\lesssim\; 3.7 \times 10^{-4}. \tag{14}$$

The later constraint is illustrated in the left-hand plot of Fig.11 for $\tan\beta = 4$. In the gray colored region the $Z - Z_s$ mixing violates the bound in (14), predicting a direct detection rate larger than presently allowed. Constraint (13) on the mixing between $\tilde{W}_s$ and $\tilde{\Sigma}_s$ can be satisfied with large μ_ϕ and μ_Σ terms (see eq.(10)). Fig.11 illustrates the allowed regions in the $m_{Z_s} - g_s$ parameter space for $\mu_\phi = \mu_\Sigma = 0.9$ TeV, 1 TeV and 1.2 TeV. One can see that the larger μ_ϕ and μ_Σ are, the more $\tilde{\phi}_s, \tilde{\Sigma}_s$ are decoupled, suppressing their mixing with $\tilde{W}_s$ and allowing a larger parameter space in $m_{Z_s} - g_s$. Thus, insisting on a dark matter candidate of this type pushes us to a very unnatural corner of parameter space, since μ_ϕ and μ_Σ are directly tied to electroweak symmetry breaking, and their TeV scale values imply large cancellations in order to get the correct Z^0 mass.

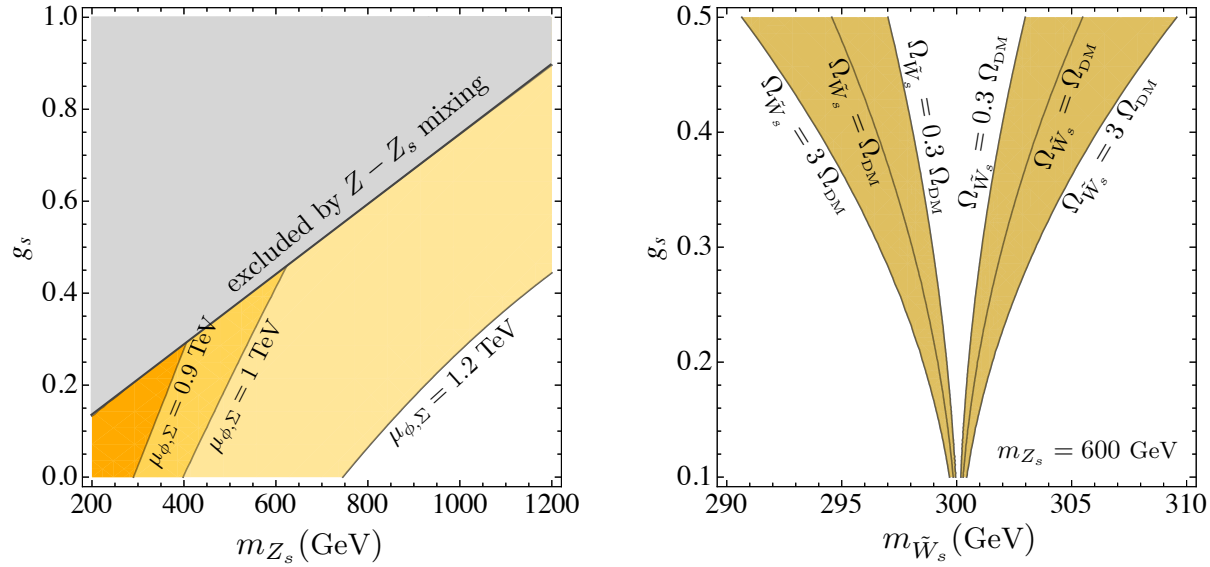

FIG. 11: (Left) Direct detection bounds on g_s vs m_{Z_s} fixing $\tan\beta$=4 and λ_u=0.6, assuming that dark matter is dominantly the gaugino LSiP, $\tilde{W}_s$. The gray region is excluded due to large $Z - Z_s$ mixing, and the colored regions are allowed when the mixing of $\tilde{W}_s$ with $\tilde{\Sigma}_s$ is sufficiently suppressed, which can be achieved with $\mu_\phi, \mu_\Sigma \sim \mathcal{O}(1)$ TeV. (Right) $\tilde{W}_s$ relic abundance, assuming m_{Z_s} = 600 GeV.

In order to complete our dark matter discussion, we also consider the dark matter relic abundance. The dominant $\tilde{W}_s\tilde{W}_s$ annihilation channel will depend on whether there are states with $SU(2)_s$ interactions that are lighter than $\tilde{W}_s$, such as the scalar components of Φ or $\bar{\Phi}$, or G-quarks. If annihilation into those (on-shell) states is kinematically allowed, they will dominate due to large g_s-couplings. If not, then the dominant annihilation channel will be $\tilde{W}_s\tilde{W}_s \to Z_s \to Z^0 h^0$. The final state $W^\pm H^\mp$ is kinematically forbidden since the mass of the charged Higgs, controlled by $\lambda_u \mu_\phi \langle\bar{\Phi}\rangle$, is at the TeV scale. Annihilation through an s-channel Z^0 to Standard Model fermions is subdominant relative to the $Z^0 h^0$ final state when $Z - Z_s$ mixing is suppressed. However, for $m_{Z_s} \lesssim 350$ GeV, the annihilation cross section to SM fermions is larger than 10% of the cross section to $Z^0 h^0$, so we include it as well in what follows.

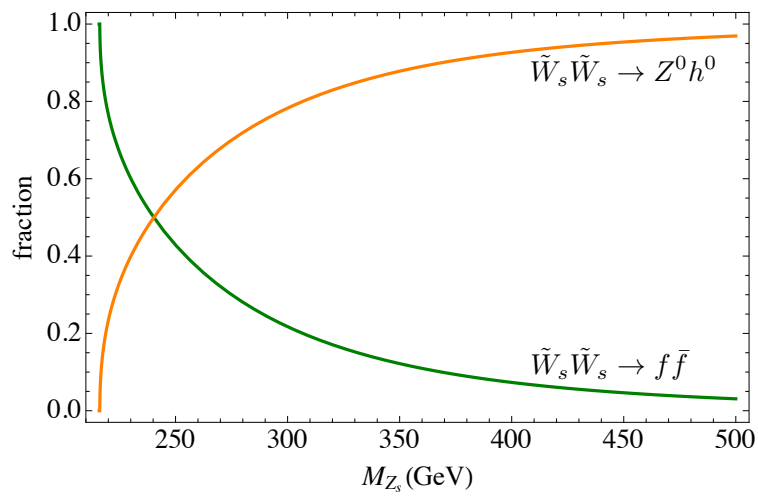

FIG. 12: Relative $\tilde{W}_s\tilde{W}_s$ annihilation cross section to Z^0h^0 vs. light SM fermions as a function of m_{Z_s}. The LSiP mass is set to $m_{\tilde{W}_s} = m_{Z_s}/2$.

The annihilation cross section for $\tilde{W}_s\tilde{W}_s \to Z_s \to Z^0h^0$ is:

$$\langle\sigma v\rangle_{Z^0h^0} = \frac{\pi}{6}\,(\alpha_s\cos\beta\sin\alpha)^2\,\xi\left(\frac{m^2_{Z^0}}{s},\frac{m^2_{h^0}}{s}\right)\,\frac{s+2m^2_{\tilde{W}_s}}{(s-m^2_{Z_s})^2+s^2\Gamma^2_{Z_s}/m^2_{Z_s}} \tag{15}$$

$$\simeq \frac{\pi}{6}\,\alpha_s^2\cos^4\beta\,\xi\left(\frac{m^2_{Z^0}}{4m^2_{\tilde{W}_s}},\frac{m^2_{h^0}}{4m^2_{\tilde{W}_s}}\right)\,\frac{4m^2_{\tilde{W}_s}+2m^2_{\tilde{W}_s}}{(4m^2_{\tilde{W}_s}-m^2_{Z_s})^2}, \tag{16}$$

where $\alpha_s = g_s^2/(4\pi)$, α is the mixing angle that determines the mass eigenstates of the CP-even higgs h^0 and H^0, and we can ignore the Z_s-width since, in the region of parameter space we are considering, it is $< \mathcal{O}(\text{MeV})$. Moreover, in the limit we are considering the CP-odd higgs is very heavy, so that $\sin\alpha \simeq -\cos\beta$. Finally,

$$\xi(x,y) \equiv \left(1+2(5x-y)+(x-y)^2\right)\sqrt{1+x^2+y^2-2(x+y+xy)}. \tag{17}$$

Annihilation into SM fermions through s-channel Z^0 or Z^0_s is given by

$$\langle\sigma v\rangle_{f\bar{f}} \simeq 4.9\pi\,\alpha_s^2\,\cos^4\beta\,\frac{m_Z^4}{(m^2_{Z_s}-m^2_Z)^2}\,(s+2m^2_{\tilde{W}_s})\left|\frac{1}{s-m^2_Z+\frac{is\Gamma_Z}{m_Z}}-\frac{1}{s-m^2_{Z_s}+\frac{is\Gamma_{Z_s}}{m_{Z_s}}}\right|^2 \tag{18}$$

From (15) and (18), we see that the dominant annihilation channel (Z^0h^0 vs $f\bar{f}$) is determined primarily by the mass of Z_s (there is a smaller dependence on $m_{\tilde{W}_s}$ that is only important near threshold). Fig.12 displays the relative fraction between the two annihilation channels (assuming $m_{\tilde{W}_s} \approx m_{Z_s}/2$). As mentioned previously, for $m_{Z_s} \gtrsim 350$GeV, the

branching to $f\bar{f}$ drops below 10%. However, it is the dominant annihilation channel for $m_{Z_s} \lesssim 240$GeV.

In order to get the observed relic abundance the total annihilation cross section must be $\mathcal{O}(1)$ pb i.e. $\langle \sigma v \rangle = \langle \sigma v \rangle_{Z^0 h^0} + \langle \sigma v \rangle_{f\bar{f}} \approx 1$ pb. The right-hand plot of Fig.11 displays the values of g_s vs. $m_{\tilde{W}_s}$ for which the correct relic abundance is obtained, fixing $m_{Z_s} = 600$ GeV and $\tan\beta = 4$.

Alternatively, if we insist on naturalness and still wish to preserve the global $U(1)_s$ symmetry, we can give up on having the LSiP as dark matter and instead interpret the direct detection bounds as a constraint on the LSiP relic abundance. In that case the bound can be re-cast as:

$$\kappa^2_{\text{LSiP}-Z} \frac{\rho}{\rho_0} \lesssim (3.7 \times 10^{-4})^2, \tag{19}$$

where ρ is the local dark matter density and $\rho_0 \sim 0.3$ GeV/cm^3. For an LSiP with electroweak charge such as $\tilde{\Sigma}_s$, that implies:

$$\rho \lesssim 10^{-6}\ \rho_0. \tag{20}$$

Achieving such an under-abundant relic density is challenging. As mentioned before, the annihilation cross section can be significantly increased by allowing the LSiP to go to on-shell states that are charged under $SU(2)_s$, such as ϕ's or G-quarks. Even so, the increase is cross section relative to the $Z^0 h^0$ channel scales as $\cos^{-4}\beta$, a gain of $\approx 3 \times 10^2$ over what is shown in Fig.11. Although an additional boost may be achieved with large values of g_s, naively resonant annihilation through Z_s would still be necessary in order to deplete the LSiP by six orders of magnitude relative to Ω_{DM}. A tension then appears since the Z_s width is also increased to $\Gamma_{Z_s} \sim \mathcal{O}(1-10)$ GeV, intrinsically limiting the efficiency of resonant annihilation. Other channels such as t-channel annihilation into $Z_s Z_s$ cannot easily achieve the required under-adundance either, even for $\alpha_s \approx 1$. Note that need for $\Omega_{\text{LSiP}} \lesssim 10^{-6}\ \Omega_{\text{DM}}$ only holds if the LSiP is dominantly the $SU(2)_{\text{EW}}$ doublet $\tilde{\Sigma}_s$. Hence, stable Dirac LSiPs that are dominantly $\tilde{\Phi}_s$ or $\tilde{W}_s$ may still be viable.

In summary, there seems to be a basic tension to achieving a combination of a) having an intact $U(1)_s$, b) having a stable Dirac particle charged under $U(1)_s$ and $SU(2)_{\text{EW}}$, and c) naturalness. The easiest ways out are to break the $U(1)_s$ completely or to a Z_2 (in both cases the LSP will be Majorana) and making the spin-independent cross section inelastic, or to preserve the $U(1)_s$ but identify it with an existing SM charge such as $U(1)_B$ or $U(1)_L$, in

which case any heavy particles could decay to SM quarks or leptons, so that direct detection constraints will not apply.

B. Bosonic LSiP Dark Matter

So far we have focused on fermionic LSiP dark matter, but it is worth commenting on bosonic LSiP dark matter as well. Many of the same issues will arise here, again due to $Z - Z'$ mixing.

The candidates for a bosonic LSiP WIMP are the $^{(s)}$ charged fields, i.e., $\phi^{(s)}, \bar{\phi}^{(s)}, \sigma_d^{(s)}$ and $\sigma_u^{(s)}$ among the scalars and $W_\mu^{(s)}$ among the vectors. Examining the spectrum of Figure 3 we see an immediate tension in our simplified limit. Although the usual $SU(2)$ D-term pushes $\sigma_0^{(s)}$ down in mass, the NMSSM-like term raises the electrically neutral component in general above the electrically charged component. Thus, even to find a reasonable DM candidate we must explore a broader range of parameter space.

Even so, for a suitable candidate, because it is charged under $SU(2)_s$, the $Z - Z'$ mixing will induce the same large Z-exchange direct detection process, meaning that only a very light WIMP would be allowed, or under a situation where $U(1)_s$ is broken down to Z_2^s.

An interesting possibility in the broader space would be if the $W^{(s)}$ was the LSiP, yielding an interesting possibility of vector DM. It, too, would interact through the Z, and so similar model building constraints apply.

Clearly the model space is broad, and is beyond the scope of what we can accomplish here. We have only begun the discussion of DM in Sister Higgs models.

V. COLLIDER IMPLICATIONS OF AN LSiP LSP

If the LSP carries sister charge, the collider phenomenology of pair produced superpartners is modified in an interesting way. Fig.13 illustrates that for a squark decay. Since the squark does not carry sister-charge, it can only decay to final states with pairs of sister-charged particles. Spin-statistics requires that the final state contains the lightest sister fermion and the lightest sister boson (note that both are stable). If such decay is mediated by an off-shell neutralino, its kinematics will be that of a 3-body decay, where two of the three final states are "invisible".

An analogous situation occurs in the MSSM with an off-shell Bino mediating a 3-body decay of the squark to a quark, a neutrino and a sneutrino. Generically a sneutrino LSP would be severely ruled out by direct detection; however those bounds could be evaded in scenarios with mixed-sneutrino LSP [45], or a sneutrino NLSP which subsequently decays to a neutrino and a gravitino.

An interesting implication of the 3-body kinematics for collider signatures of squarks is a reduction in the visible and missing energy of the events, decreasing the efficiency of these types of signals to pass standard cuts and making its discovery more challenging. Fig.14 contrasts the distributions of missing transverse energy and H_T[6] for the two types of decay, assuming a squark mass $m_{\tilde{q}} = 400$ GeV. For the 2-body topology, the neutralino mass is chosen to be $m_{\tilde{\chi}^0} = 200$ GeV, and for the 3-body topology the two invisible final state particles have masses $m_{\sigma^0} = m_{\tilde{\sigma}^0} = 100$ GeV.

As a brief illustration of the effects of this different kinematics, we perform a Monte Carlo simulation[7] of the two benchmark scenarios introduced in Fig.14, and compare their relative efficiency to pass a standard jets plus $\not{E}_T$ search. Specifically, we consider the baseline selection in the multijets plus missing transverse energy search performed by CMS [46] with 4.98 fb^{-1}. Such selection required events with $H_T > 500$ GeV, $\not{H}_T > 200$ GeV, and at least 3 central jets with $p_T > 50$ GeV, among other requirements. Assuming first and second generation degenerate squarks, this CMS search constrained the production cross section of the 2-body benchmark to be $\sigma_{\tilde{q}\tilde{q}^*} \lesssim 1.2$ pb. This cross section is roughly half of the reference QCD cross section at NLO, and therefore the 2-body benchmark is ruled out. The efficiency

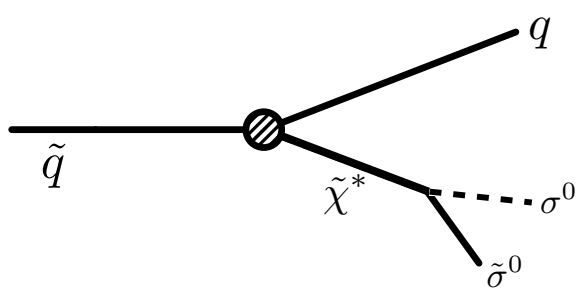

FIG. 13: Topology of a 3-body squark decay to two invisible sister particles.

[6] H_T is defined as the scalar sum of the p_T of all jets in the event.

[7] We use `MadGraph5` for event simulation, `Pythia6` for decays, showering and hadronization and `PGS4` for detector response.

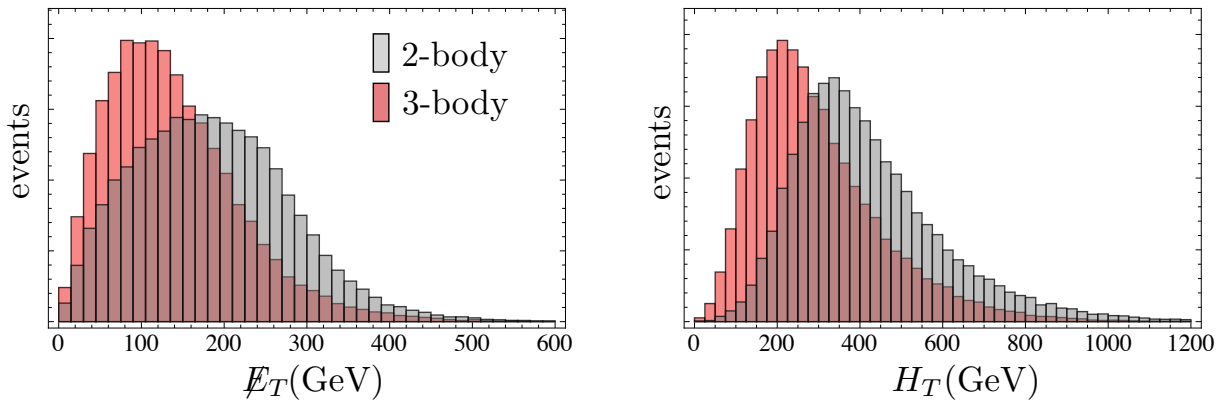

FIG. 14: Contrast between the $\not{E}_T$ and H_T distributions of 2- and 3-body squark decays. We assume $m_{\tilde{q}} = 400$ GeV; for the 2-body decay, the mass of the invisible particle is 200 GeV, and for the 3-body decay the two invisible particles have mass of 100 GeV.

for the 3-body topology is roughly 30% of its 2-body counterpart. Naively applying this lowered efficiency to the cross section limit, we get a bound $\sigma_{\tilde{q}\tilde{q}^*} \lesssim 4$ pb, suggesting that the 3-body topology is slightly beyond this search's reach.

Dedicated searches for more challenging SUSY signatures, such as stop and sbottom pair production are also impacted by the 3-body decay kinematics. We briefly look into one more case, namely, sbottom searches targeting the 2-body topology $\tilde{b}_1 \to b\tilde{\chi}^0$. ATLAS performed a sbottom dedicated search [47] with 2 fb^{-1} of integrated luminosity. The baseline event selection required $\not{E}_T > 130$ GeV and two b-tagged jets with $p_T^{j_1} > 130$ GeV and $p_T^{j_2} > 50$ GeV, among other requirements. Assuming a sbottom mass $m_{\tilde{b}_1} = 300$ GeV, we compared the selection efficiencies for the 2- and 3-body topologies, assuming 50 GeV for the mass of the invisible particle in the 2-body decay, and 25 GeV for both invisible particles in the 3-body decay. The efficiency of the 3-body topology relative to the 2-body topology is $\sim 33\%$. Additional signal regions were defined by placing cuts on the boost-corrected contransverse mass [48], m_{CT}, which have even lower relative efficiencies. The ATLAS search excluded this 2-body decay mass point, bounding the sbottom production cross section to be $\sigma_{\tilde{b}_1\tilde{b}_1^*} \lesssim 0.44$ pb. Conservatively using the relative efficiency of the baseline selection to infer the bound on the 3-body topology, we obtain $\sigma_{\tilde{b}_1\tilde{b}_1^*} \lesssim 1.3$ pb, which is at the borderline of ATLAS' exclusion reach.

While we have only touched on the phenomenology of such decays, we can see that there are clearly interesting results: squarks of 400 GeV and sbottoms of 300 GeV may plausibly

be allowed under current LHC limits. The impact of 3-body topologies in LHC searches is a rich subject that deserves more careful investigation. We defer this to future work.

VI. THE COLLIDER PHYSICS OF G-QUARKS

The phenomenology of the G-quarks is intimately tied into a number of questions: first, are we considering a scenario in which there is a preserved sister charge or not? If there is a sister charge, is it identified with any of the existing global standard model quantum numbers (i.e., baryon or lepton number)? Moreover, if there is a sister charge, and R-parity is conserved, is the lightest sister-charged particle (LSiP) also the lightest supersymmetric particle (LSP)?

Let us begin with the simplest case - there is a single vectorlike pair of G-quarks D_g and D_g^c. These fields can interact with the standard model via $\bar{\Phi} D_g D^c$ which preserves a $U(1)_s$. We could have instead included $\Phi D_g D^c + \Sigma_d D_g^c Q$, which gives opposite $U(1)_s$ assignments, but since D_g falls into a 5 with Σ_d, it is natural to give them the same charge assignments.

If this sister charge is broken, it can be broken spontaneously or explicitly. If broken spontaneously by ϕ and σ vevs, it should be gauged to avoid the presence of a dangerous goldstone boson. In this case the production of G-quarks will naturally yield $D_g^{(c)} \to D^c \phi(\sigma)$. With $\phi(\sigma) \to \bar{b}b$ we can naturally expect six-jet events with three-jet resonances. CMS has performed a search for three-jet resonances in the context of hadronic RPV-gluino decays [49], excluding gluino masses $m_{\tilde{g}} \lesssim 460$ GeV. The limits from this search on G-quarks are much weaker, since the $D_g \bar{D}_g$ production cross section is roughly an order of magnitude smaller than gluino pair-production. A more quantitative study of the limits on such objects can be found in [50].

The G-squarks will mix with SM squarks, and so if they are lighter than the G-quarks, they will produce conventional jet+$\not{E}_T$ signals when pair-produced. However, if heavier than the G-quarks, and assuming a conventional LSP such as a Bino, they will decay $\tilde{d}_g \to D_g + \chi^0 \to j\bar{b}b\tilde{\chi}^0$. Thus, they will produce six jet + $\not{E}_T$ events, with large numbers of b-tags. For large G-quark -G-squark splittings, standard jets+$\not{E}_T$ searches might be sensitive to these signatures, although with smaller reach relative to MSSM squarks due to the reduced G-squark cross section. (The limits on MSSM squarks usually quoted by LHC searches implicitly assume degenerate 1st and 2nd generation squarks, with an effective cross section

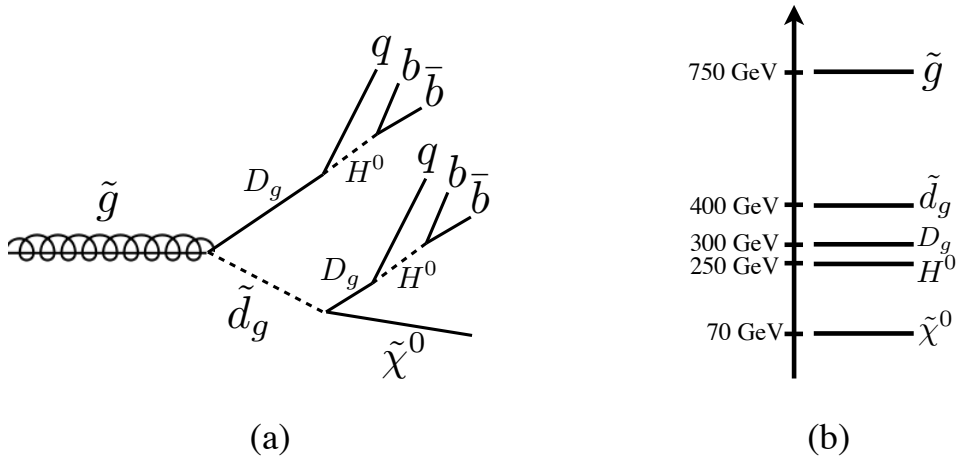

FIG. 15: (a) Gluino decay topology to $\tilde{d}_g D_g$, resulting in final states with $2j + 4b + \tilde{\chi}^0$. (b) Benchmark spectrum with small splitting between $m_{\tilde{d}_g}$ and $m_{D_g} + m_{\tilde{\chi}^0}$.

eight times larger than the cross section of a single squark species). On the other hand, for more compressed spectra the missing energy in the events may be significantly reduced, rendering high-$\not{E}_T$ searches insensitive to such models. However, searches with high-jet multiplicity or/and high b-jet multiplicity may be effective for these topologies. Still, the cross section may be the limiting factor in sensitivity to G-quarks.

An interesting scenario to consider is the case of gluino-boosted G-(s)quark production. If the G-(s)quarks are lighter than the gluinos with all other squarks heavy, the on-shell decay $\tilde{g} \to D_g \tilde{d}_g$ might dominate. This decay topology is illustrated in Fig.15(a) - each gluino decays to $2j + 4b + \tilde{\chi}^0$ final states, for a total of 12 jet + $\not{E}_T$ events (with large numbers of b-tags). Dedicated searches with high b-jet multiplicity could be sensitive to such final states, such as the search performed by ATLAS in missing energy and at least 3 b-jets [51]. This search, however, places a strong requirement on missing transverse energy ($\not{E}_T > 160$ GeV), which narrows its sensitivity to spectra with large mass splittings between G-quarks and G-squarks. In Fig.15(b) we illustrate a typical spectrum which the ATLAS 3-bjets+$\not{E}_T$ search is not sensitive to. In the rest frame of $\tilde{d}_g$, the 3-momentum of its neutralino daughter $\tilde{\chi}^0$ is $|\vec{p}_{\tilde{\chi}^0}| \sim 62$ GeV, and hence this signal has a very low efficiency to pass the $\not{E}_T > 160$ GeV requirement of this search. Despite its low missing energy, this particular spectrum is distinctive enough to be caught by high jet-multiplicity searches, and indeed, through our MC simulations we estimate that this mass spectrum is excluded by the ATLAS search in this channel[52]. This search looks for events with up to 9 hard jets,

and bounds the gluino cross section in the benchmark of Fig.15 to be $\sigma_{\tilde{g}\tilde{g}} \lesssim 0.1$ pb, which is roughly the value of the reference QCD cross section for a 750 GeV gluino. Hence, this point is marginally excluded.

If this sister charge is not broken by ϕ and Σ vevs, there must be a larger gauged sister group, i.e., $SU(2)_s^{\text{gauge}} \times U(1)_\Sigma^{\text{global}}$, which leaves some residual $U(1)_s$. In this case, we must determine the identity of the LSiP. Assuming it is stable, it is most easily neutral, such as the ϕ, in which case the decay $D_g \to D + \text{LSiP}$ would resemble a standard SUSY squark search. Likewise $\tilde{d}_g \to D_g \tilde{\chi}^0 \to D + \tilde{\chi}^0 + \text{LSiP}$ could be discoverable as well with enough luminosity.

The sister charge can be broken explicitly if the sister group is $SU(2)_s^{\text{gauge}} \times U(1)_\Sigma^{\text{global}}$. I.e., we can write $U(1)_\Sigma$ breaking terms $\bar{\phi} D_g D^c + \phi D_g D^c + \Sigma_d D_g^c Q$. In this case, both G-quarks decay is as in the broken $U(1)$ case above, with multijets and many b-tags.

Another interesting case is when the $U(1)_s$ is identified with some SM global charge. As we discuss in the appendix, the operator $U D_s^c D_s^c$ naturally identifies $U(1)_s$ with baryon number, but does not break R-parity. In this case, while the G-quarks without sister charge would decay into three jets, the sister-charged G-quark would decay $\tilde{D}_s^c \to \tilde{u} D_s \to U + D + b\bar{b} + \chi_0$, yielding eight-jet plus $\not{E}_T$ signals, and sister charged G-squarks would decay $\tilde{d}_g^{(s)} \to u D_g^c \to u d \bar{b} b$, i.e., we would have eight jet events, with four jet subresonances. (Here, the G-squarks of different sister charge have different R-parity under the intact R-parity after ϕ, Σ acquire vevs.)

In summary, the presence of G-quarks can radically alter our expectations for phenomenology. In particular, if the G-quarks are lighter than the gluinos, it can dramatically enhance the number of jets (in particular b-jets) in the event, and suppress the $\not{E}_T$ signal.

VII. CONCLUSIONS

If there is new physics at the weak scale, the Higgs boson and its properties are the first indication of what that may be. The value of 125 GeV constrains a wide range of supersymmetric models, and may be pointing us to a new contribution to the Higgs quartic. Models such as the NMSSM can raise it, but at the cost of including pure singlets, and only with very large couplings.

In the context of the NMSSM, however, it is possible to replace H_d with a new field Σ_d,

a "Sister Higgs," and still yield the usual tree-level corrections to the Higgs mass. Such a field need not couple to ordinary matter, and, indeed, would not if it carries a charge under some new sister group G_s. If that group is gauged, new contributions to the NMSSM-like superpotential term $\Phi H_u \Sigma_d$ can drive it large, easily boosting the Higgs mass up to 125 GeV, or beyond.

The presence of this new Higgs has a number of consequences: as with a Type I 2HDM, a light charged Higgs is now possible. The overall consequences allow a modification to the Higgs properties, such as a boosted VBF $\gamma\gamma$ signal, as well as a boosted inclusive $\gamma\gamma$ signal, with suppressed decays to fermions [53].

Once we have expanded the theory to include the sister Higgs, we could also expect yet more Higgses when G_s is non-Abelian. With these fields we can write down the R-parity violating operator $H_d \Sigma_d E^c$ once G_s is broken. Because this operator couples only to right handed fermions, it is far less constrained that traditional lepton-number violating RPV operators. While it changes the phenomenology of SUSY decays, if the finals state contains τ's, the current search strategies should find it. At the same time, it motivates searches for leptoquarks, with squarks decaying to $j\ell^{\pm}$, but importantly without an associated missing energy signal (as there is no accompanying channel with a neutrino).

The Z' in these models is natural, and can be quite light without significant constraints from the ρ parameter or searches for dilepton resonances. Corrections to Zh production can be sizeable at the Tevatron but are not expected to be large at the LHC without additional couplings to SM fermions (although such corrections could come from mixing with the G-quarks at the cost of flavor constraints).

Appeals to grand-unification motivate us to consider the presence of colored fields which are charged under G_s as well. Unlike conventional fourth-generations, these vectorlike quarks decay dominantly through ϕ emission (or $\phi - h$ mixing induced h emission). These G-quarks can have important effects on phenomenology. In particular, if they are lighter than the gluino, they can open up new decay channels with lower missing energy and more jets than conventional gluino decays. The phenomenology of G-quarks (and G-squarks) is very rich and we have only begun the discussion in this work.

In the presence of a sister charge, the lightest sister-charged particle (LSiP) would be stable. If it is a Dirac fermion, constraints from direct detection searches would have excluded it in all but the narrowest corners of parameter space where the WIMP is very under

abundant or is very light. It may thus be most natural that symmetry should be broken, at least to a Z_2, or become identified with B or L so that the LSiP can decay into it.

Regardless, if the LSP also carries sister charge, squarks may only decay via three-body decays. Such decays have reduced missing energy and H_T, and are consequently less constrained by existing jets+$\not{E}_T$ searches.

In summary - very complex phenomenology can arise by the simple extension of the MSSM by the presence of G_s-charged sister-Higgs fields and their related G-quarks. A low energy theory that is not the MSSM, nor even the NMSSM, but one expanded with additional Higgs fields easily provides for a Higgs mass as large as 125 GeV and a profoundly changed LHC phenomenology. As more data accumulate, we shall see soon if such a rich Higgs sector is realized in nature.

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