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Citation: Cheung, Clifford, Yasunori Nomura, and Jesse Thaler. "Goldstini." *Journal of High Energy Physics* 2010.3 (2010) : 1-27-27.

As Published: [http://dx.doi.org/10.1007/jhep03\(2010\)073](http://dx.doi.org/10.1007/jhep03(2010)073)

Publisher: Springer Berlin / Heidelberg

Persistent URL: <http://hdl.handle.net/1721.1/64478>

Version: Author's final manuscript: final author's manuscript post peer review, without publisher's formatting or copy editing

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Goldstini

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Supersymmetric phenomenology has been largely bound to the hypothesis that supersymmetry breaking originates from a single source. In this paper, we relax this underlying assumption and consider a multiplicity of sectors which independently break supersymmetry, thus yielding a corresponding multiplicity of goldstini. While one linear combination of goldstini is eaten via the super-Higgs mechanism, the orthogonal combinations remain in the spectrum as physical degrees of freedom. Interestingly, supergravity effects induce a universal tree-level mass for the goldstini which is exactly twice the gravitino mass. Since visible sector fields can couple dominantly to the goldstini rather than the gravitino, this framework allows for substantial departures from conventional supersymmetric phenomenology. In fact, this even occurs when a conventional mediation scheme is augmented by additional supersymmetry breaking sectors which are fully sequestered. We discuss a number of striking collider signatures, including various novel decay modes for the lightest observable-sector supersymmetric particle, gravitinoless gauge-mediated spectra, and events with multiple displaced vertices. We also describe goldstini cosmology and the possibility of goldstini dark matter.

I. INTRODUCTION

Supersymmetry (SUSY) is a theoretically motivated and well-studied framework which solves the hierarchy problem and offers a rich phenomenology [1]. Of course, if SUSY is to be realized in nature, then it must be spontaneously broken. To this end, it is conventionally assumed that SUSY breaking originates from the dynamics of a single hidden sector.

While the notion of single sector SUSY breaking is convenient as a simplifying premise, it is not very generic in light of top-down considerations. In particular, string theoretic constructions routinely predict a multiplicity of geographically sequestered sectors [2], any number of which could independently break SUSY. In this paper we will explore the generic implications of multiple sector SUSY breaking.

Consider the low energy effective field theory describing N such sequestered sectors. In the limit in which these sectors are completely decoupled—even gravitationally—they enjoy an N -fold enhanced Poincaré symmetry because energy and momentum are *separately* conserved within each sector. Likewise, if SUSY is a symmetry of nature then it is similarly enhanced, such that

$$\text{SUSY} \xrightarrow{\text{decoupled}} \text{SUSY}^N \equiv \otimes \prod_{i=1}^N \text{SUSY}_i. \quad (1)$$

Because this enhancement is an accidental consequence of the decoupling limit, gravitational interactions explicitly break SUSY^N down to a diagonal combination corresponding to the genuine supergravity (SUGRA) symmetry. Consequently, the “orthogonal” SUSY^{N-1} are only approximate global symmetries.

In the event that F -term breaking occurs independently in each sector, each SUSY_i will be spontaneously broken at a scale F_i , yielding a corresponding goldstino

η_i .¹ In unitary gauge, one linear combination of goldstini, η_{long} , is eaten by the gravitino via the super-Higgs mechanism, leaving $N - 1$ goldstini in the spectrum. We denote these fields by ζ_a , where $a = 1, \dots, N - 1$.

Since the remaining $N - 1$ goldstini correspond to the approximate SUSY^{N-1} which are explicitly broken by SUGRA, one should not expect these goldstini to remain exactly massless. In fact, we will show that they acquire a *tree-level* mass

$$m_a = 2m_{3/2}, \quad (2)$$

induced by gravitational effects. As we will see, the curious factor of 2 is ultimately fixed by the symmetries of SUGRA, and we will robustly derive it in a number of different ways.

Up to now, SUSY phenomenology has been almost exclusively devoted to a scenario in which the gravitino and the goldstino are effectively one and the same.² In the context of multiple sector SUSY breaking, however, this corresponds to a rather privileged arrangement in which the dominant contributions to SUSY breaking in the supersymmetric standard model (SSM) sector arise from the SUSY breaking sector with the highest SUSY breaking scale. In any other situation, the SSM fields will actually couple more strongly to the goldstini than to the gravitino, and this will have a significant impact on collider physics and cosmology. A simple context in which

¹ Throughout the paper we take a field basis where F_i are all real and positive, and assume that $F_i \geq F_{i+1}$ without loss of generality. We will also focus on the case where SUSY breaking still occurs in the $M_{\text{Pl}} \rightarrow \infty$ limit, and only briefly comment on “almost no-scale” SUSY-breaking sectors in the Appendix. The possibility of D -term breaking will be left to future work.

² To our knowledge, the only mention of multiple goldstini in the literature appears in Ref. [3].

this occurs is when a conventional SUSY breaking scenario is augmented by additional SUSY breaking sectors which are fully sequestered (see Fig. 1 in Sec. V A).

This paper is organized as follows. In Sec. II, we review an analogous construction for Goldstone bosons arising from multiple symmetry breaking. The goldstini case of multiple SUSY breaking is then presented in Sec. III. We derive the relation $m_a = 2m_{3/2}$ in Sec. IV, using both a Stückelberg method and a conformal compensator method. A direct SUGRA calculation of the factor of two appears in the Appendix. Corrections to this mass relation are given in Sec. V, and the couplings to the SSM are given in Sec. VI. Possible LHC signatures of this scenario—including wrong mass “gravitinos”, gravitinoless gauge mediation, smoking gun evidence for the factor of two, three-body neutralino decays, and displaced monojets—are presented in Sec. VII. Goldstini cosmology is described in Sec. VIII, including scenarios that yield goldstini dark matter. We conclude in Sec. IX.

II. GOLDSTONE ANALOGY

Because the notion of multiple sector SUSY breaking is not a familiar one, it is instructive to analyze an analogous construction involving multiple $U(1)$ symmetry breaking. Consider a scenario in which ϕ_1 and ϕ_2 are complex scalar fields which enjoy separate global symmetries $U(1)_1$ and $U(1)_2$. Furthermore, assume that the diagonal $U(1)_V$ is gauged and that ϕ_1 and ϕ_2 have no direct couplings except for gauge interactions.

A. Fields and Couplings

If ϕ_1 and ϕ_2 separately acquire vacuum expectation values (vevs), then we can non-linearly parameterize the Goldstone modes as

$$\phi_i = f_i e^{i\pi_i/\sqrt{2}f_i}, \quad (3)$$

for $i = 1, 2$. One linear combination of π_1 and π_2 is eaten via the Higgs mechanism. The orthogonal combination, φ , corresponds to a physical pseudo-Goldstone boson that arises from the spontaneous breaking of a global $U(1)_A$ axial symmetry. Concretely, go to a basis

$$\begin{pmatrix} \pi_1 \\ \pi_2 \end{pmatrix} \xrightarrow{\text{unitary gauge}} \begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} \pi_{\text{long}} \\ \varphi \end{pmatrix}, \quad (4)$$

where $\tan\theta = f_2/f_1$ and $f_{\text{eff}} = \sqrt{f_1^2 + f_2^2}$. In unitary gauge, π_{long} becomes the longitudinal mode of the $U(1)_V$ gauge boson.

The interactions of φ with other fields can be obtained from plugging the parameterization of Eqs. (3) and (4) into couplings involving ϕ_i and those fields. Note a

crucial difference between the couplings of π_{long} and φ . While one can always do field redefinitions such that π_{long} couples only derivatively, $\mathcal{L}_{\text{int}} = (1/f_{\text{eff}})(\partial_\mu\pi_{\text{long}})J^\mu$ where J^μ is the $U(1)_V$ current, there is no guarantee that the same can be done for φ .

B. Masses

As is well known, π_1 and π_2 are exactly massless in the limit in which $U(1)_1 \times U(1)_2$ is an exact symmetry of the Lagrangian. One way of understanding this fact is to consider the unitary gauge Lagrangian for the massive $U(1)_V$ gauge boson,

$$\mathcal{L}_{\text{unit}} = -\frac{1}{4g^2}F_{\mu\nu}F^{\mu\nu} - f^2 A_\mu A^\mu. \quad (5)$$

For the moment, let us assume that $U(1)_1$ is broken but $U(1)_2$ is preserved. As a consequence, there is a single eaten Goldstone mode, π_1 . Using the Stückelberg replacement, we can reinstate π_1 as a propagating degree of freedom by applying a gauge transformation

$$A_\mu \rightarrow A_\mu + \frac{1}{\sqrt{2}f}\partial_\mu\pi_1, \quad (6)$$

and promoting π_1 to a dynamical field. Doing so yields

$$\mathcal{L} = -\frac{1}{2}\partial_\mu\pi_1\partial^\mu\pi_1 + \text{terms involving } A_\mu. \quad (7)$$

Obviously, the exact same argument can apply in the case in which $U(1)_2$ is broken and $U(1)_1$ is preserved. Thus, if $U(1)_1$ and $U(1)_2$ are independently broken, then the Lagrangian must take the form

$$\mathcal{L} = -\frac{1}{2}\sum_i\partial_\mu\pi_i\partial^\mu\pi_i + \text{terms involving } A_\mu, \quad (8)$$

and so a mass term is forbidden for either π_i . Said another way, either π_i could have been eaten by A_μ , so both are required to be massless. This implies that the uneaten Goldstone mode, φ , is massless.

If it is not the case that $U(1)_1 \times U(1)_2$ is an exact symmetry, then the above argument is only approximate. In particular, any explicit $U(1)_A$ violating, $U(1)_V$ preserving operators will provide a mass term for the uneaten mode, φ , at tree level. Moreover, even if such operators are missing, they can be generated radiatively. For example, this occurs in a non-Abelian Goldstone theory in which ϕ_1 and ϕ_2 are in fundamental representations of $SU(k)_1$ and $SU(k)_2$ global symmetries, respectively, of which the diagonal $SU(k)_V$ combination is gauged. Since the gauge interactions explicitly violate the $SU(k)_A$ global symmetry, radiative corrections will generate operators of the form

$$|\phi_1^\dagger\phi_2|^2, \quad (9)$$

which induce a mass for φ , albeit at loop level. As we will see shortly, the non-Abelian theory provides the closest analogy to multiple sector SUSY breaking—SUGRA,

which is precisely the gauged diagonal SUSY, explicitly violates the orthogonal SUSY $^{N-1}$ and thus induces nonzero masses for the uneaten goldstini. The important difference in the case of SUSY is that these masses will arise at tree level rather than at loop level.

III. GOLDSTINI FIELDS AND COUPLINGS

The discussion of multiple sector SUSY breaking exactly parallels that of the previous section. We will focus here on the case of F -term breaking, and imagine that there exist two chiral superfields, X_1 and X_2 , that reside in two sequestered sectors. In the absence of direct couplings, gravitational or otherwise, these fields enjoy an enhanced SUSY $_1 \otimes$ SUSY $_2$ symmetry. Assuming that the highest component of X_i acquires a vev equal to F_i , then SUSY $_i$ is broken and we can use the non-linear parameterization

$$\begin{aligned} X_i &= e^{Q\eta_i/\sqrt{2}F_i}(x_i + \theta^2 F_i) \\ &= x_i + \eta_i^2/2F_i + \sqrt{2}\theta\eta_i + \theta^2 F_i, \end{aligned} \quad (10)$$

for $i = 1, 2$, where $Q = \partial/\partial\theta$ is the generator of SUSY transformations and we have neglected all derivatively coupled terms.³ Here η_i is the goldstino corresponding to the F -term breaking of SUSY $_i$. Note that this form is identical to the usual linear parameterization of a chiral superfield except for the presence of η_i^2 in the lowest component of X_i .

In the presence of SUGRA, the diagonal combination of SUSY $_1$ and SUSY $_2$ is gauged, and thus one of the goldstini is eaten. As before, it is convenient to work in a basis

$$\begin{aligned} \begin{pmatrix} \eta_1 \\ \eta_2 \end{pmatrix} &= \begin{pmatrix} \cos\theta & -\sin\theta \\ \sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} \eta_{\text{long}} \\ \zeta \end{pmatrix} \\ &\xrightarrow{\text{unitary gauge}} \begin{pmatrix} -\sin\theta\zeta \\ \cos\theta\zeta \end{pmatrix}, \end{aligned} \quad (11)$$

where $\tan\theta = F_2/F_1$ and $F_{\text{eff}} = \sqrt{F_1^2 + F_2^2}$. Thus, η_{long} is eaten by the gravitino, and ζ remains a propagating degree of freedom.

The interactions of ζ with other fields can be obtained using the parameterization of Eqs. (10) and (11). Since X_i is a true chiral superfield, the couplings of ζ can be obtained directly in superspace. While one can always work in a field basis where η_{long} couples only derivatively, $\mathcal{L}_{\text{int}} = (1/F_{\text{eff}})(\partial_\mu\eta_{\text{long}})\tilde{J}^\mu$ where \tilde{J}^μ is the supercurrent, the same cannot be done in general for ζ .

If the number of sequestered SUSY breaking sectors is greater than two, then there will be multiple uneaten goldstini ζ_a , which are related to η_i by

$$\eta_i = V_{ia}\zeta_a, \quad (12)$$

where V_{ia} is the $N \times (N-1)$ part of the unitary matrix which goes from the η_i basis to the $\{\eta_{\text{long}}, \zeta_a\}$ basis. The ζ_a fields are orthogonal to the eaten mode. Since

$$\eta_{\text{long}} = \frac{1}{F_{\text{eff}}} \sum_i F_i \eta_i, \quad (13)$$

this implies $\sum_i F_i V_{ia} = 0$. The form of V_{ia} is determined by the mass matrix of ζ_a , which we will now discuss.

IV. GOLDSTINI MASSES

In Sec. IIB, we saw that uneaten Goldstone bosons typically acquire masses from loops of non-Abelian gauge bosons. SUGRA effects similarly induce masses for the goldstini—only this happens at tree level! More precisely, in the limit in which each sector couples only through SUGRA, all goldstini acquire a tree level mass which is universal and given by $m_a = 2m_{3/2}$.⁴ While the factor of 2 may be verified explicitly by considering the explicit SUGRA Lagrangian (see the Appendix), we find it more illuminating to derive it in two separate but more direct ways. Collider and cosmological implications of this universal mass will be discussed in Secs. VII and VIII.

A. Two via Stückelberg

The simplest way of understanding $m_a = 2m_{3/2}$ is in analogy with the logic of Sec. IIB. We start from a unitary gauge SUGRA Lagrangian, where the quadratic action for the gravitino is [5]

$$\mathcal{L}_{\text{unit}} = \epsilon^{\mu\nu\rho\sigma} \bar{\psi}_\mu \bar{\sigma}_\nu \partial_\rho \psi_\sigma - m_{3/2} (\psi_\mu \sigma^{\mu\nu} \psi_\nu + \bar{\psi}_\mu \bar{\sigma}^{\mu\nu} \bar{\psi}_\nu), \quad (14)$$

where $\sigma^{\mu\nu} \equiv (\sigma^\mu \bar{\sigma}^\nu - \sigma^\nu \bar{\sigma}^\mu)/4$ and $m_{3/2} \simeq F_{\text{eff}}/\sqrt{3}M_{\text{Pl}}$.⁵ Now consider the scenario in which SUSY has been broken *only* in sector 1, and the corresponding goldstino η_1 has been eaten. We can reinstate the η_1 degree of freedom by applying the Stückelberg construction—that is, applying a SUGRA transformation on the unitary gauge Lagrangian, and then promoting the SUGRA parameter to a dynamical field. In particular, we apply the SUGRA transformation [5]

$$\psi_\mu \rightarrow \psi_\mu + \sqrt{\frac{2}{3}} m_{3/2}^{-1} \partial_\mu \eta_1 + \frac{i}{\sqrt{6}} \sigma_\mu \bar{\eta}_1, \quad (15)$$

⁴ It may appear contradictory that the goldstini acquire a tree-level mass, since they are derivatively coupled in the limit of global SUSY. Nevertheless, for finite M_{Pl} , the goldstini couple not just as $\partial_\mu \eta_i$ but also as $\sigma_\mu \bar{\eta}_i$, so a mass term is not forbidden.

⁵ In this paper we assume that SUSY is broken in the global limit and that the SUSY breaking vacuum is unaffected by finite M_{Pl} effects. Some of the equations below, e.g. Eq. (15), do not hold if we relax this assumption (see Ref. [6] for a clear discussion of the general SUSY transformation laws for the gravitino). A more general case is discussed briefly in the Appendix.

³ A similar non-linear parametrization was considered in Ref. [4] for a single goldstino.

to the Lagrangian, yielding

$$\mathcal{L} = -i\bar{\eta}_1\bar{\sigma}^\mu\partial_\mu\eta_1 - \frac{1}{2}(2m_{3/2})(\eta_1^2 + \bar{\eta}_1^2) + \dots, \quad (16)$$

where the ellipses denote all terms involving the ψ_μ , including the mixing terms between the gravitino and the goldstino. Note how the kinetic term for η_1 is generated by the cross term obtained from Eq. (15). Given the normalization of a Majorana fermion, this implies a goldstino mass of $m_1 = 2m_{3/2}$.

Now of course if there is only one goldstino, then this mass is not physical, since η_1 is eaten via the super-Higgs mechanism. However, if there is multiple sector SUSY breaking, then there will be several goldstini η_i . Since any one of the X_i could have broken SUSY on its own and been eaten by the gravitino, all of them must take the form of Eq. (16). Thus, with multiple goldstini, the Stückelberg Lagrangian becomes

$$\mathcal{L} = \sum_i \left\{ -i\bar{\eta}_i\bar{\sigma}^\mu\partial_\mu\eta_i - \frac{1}{2}(2m_{3/2})(\eta_i^2 + \bar{\eta}_i^2) \right\} + \dots, \quad (17)$$

where the ellipses include the mixing between the gravitino and the eaten goldstino, which is now some linear combination of the η_i .⁶ We can now rotate the fermions by an orthogonal matrix, to isolate the eaten goldstino mode, and then go to unitary gauge. Since the η_i mass matrix is proportional to the identity, the leftover goldstini, ζ_a , will all have mass $m_a = 2m_{3/2}$.

B. Two via the Conformal Compensator

An alternative way of understanding the relation $m_a = 2m_{3/2}$ is to use the conformal compensator formalism [7]. Morally, the factor of 2 arises because the conformal compensator couples to mass dimension, and F_i has mass dimension 2. To see this in a simple example, consider the case of several sectors which independently break SUSY via a Polonyi superpotential

$$\begin{aligned} \mathcal{L} = & \int d^4\theta C^\dagger C \sum_i (X_i^\dagger X_i + \dots) \\ & + \int d^2\theta C^3 \sum_i \mu_i^2 X_i + \text{h.c.}, \end{aligned} \quad (18)$$

where $C = 1 + \theta^2 m_{3/2}$ is the conformal compensator, the dots indicate higher order terms necessary to stabilize the scalar components of X_i , and we have chosen a sequestered form for the Kähler potential. By rescaling

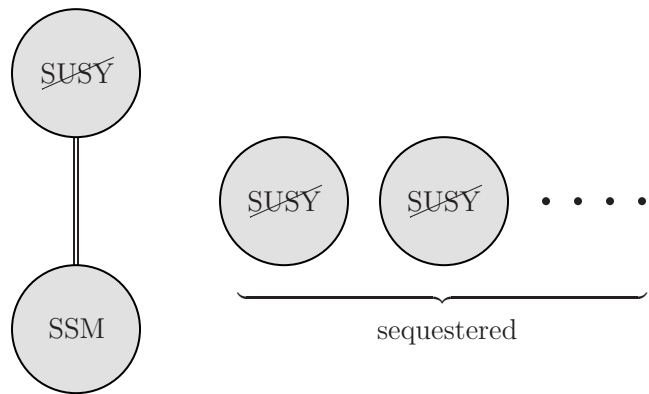


FIG. 1: A schematic depiction of a scenario in which the SSM sector couples to only one of the SUSY breaking sectors. Note that this setup still leads to interactions between SSM fields and goldstini in the sequestered sectors, since the goldstino of the sector coupling to the SSM is in general a linear combination of the gravitino and uneaten goldstini.

$X_i \rightarrow X_i/C$, we see that C only couples to dimensionful parameters—namely, μ_i . Plugging in for the lowest component of the non-linear parameterization of X_i in Eq. (10), we obtain

$$\begin{aligned} \mathcal{L} \supset & \int d^2\theta C^2 \sum_i \mu_i^2 X_i \\ = & -\frac{1}{2}(2m_{3/2}) \sum_i \eta_i^2 + \text{const.}, \end{aligned} \quad (19)$$

where we have solved for the auxiliary fields $F_i = -\mu_i^2$ and plugged in for the conformal compensator. The fact that μ_i^2 has mass dimension 2 implies that conformal compensator couples as C^2 , yielding the important factor of 2 in the goldstini mass.

V. DEVIATIONS FROM THE SEQUESTERED LIMIT

So far, we have limited our discussion to the case where the only interactions between SUSY breaking sectors arise from SUGRA. This is certainly the case if every sector, including the SSM sector, is sequestered from one another and SUSY breaking is communicated to the SSM via SUGRA effects, i.e. anomaly mediation. In this section, we consider the case where one or more SUSY breaking sectors have direct couplings to the SSM to mediate SUSY breaking. We discuss effects of such couplings on the goldstini properties.

A. Single Sector Mediation

The simplest deviation from the fully sequestered limit is for direct couplings to exist only between the SSM and

⁶ One might think that each η_i should have a mass given by $2F_i/\sqrt{3}M_{\text{Pl}}$ (twice the gravitino mass for sector i alone) instead of $2F_{\text{eff}}/\sqrt{3}M_{\text{Pl}}$ (twice the gravitino mass for all sectors together). This, however, would not lead to the correct mass for the eaten mode, which must take the form of Eq. (16).

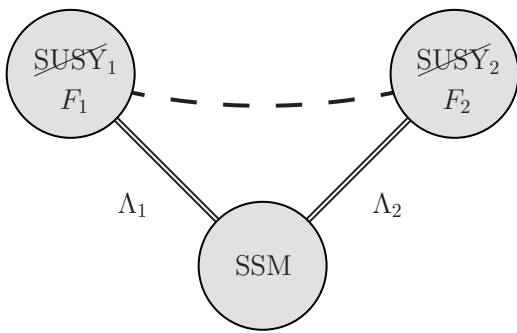


FIG. 2: A schematic depiction of a scenario in which sectors 1 and 2 have direct interactions to the SSM sector via operators suppressed by Λ_1 and Λ_2 , respectively (double lines). These interactions induce direct couplings between sectors 1 and 2 through radiative corrections (dashed line).

one of the SUSY breaking sectors, as illustrated in Fig. 1. This corresponds to the situation where a conventional SUSY breaking scenario, such as gauge mediation, is augmented by one or more fully-sequestered SUSY breaking sectors. This may easily arise in realistic top-down setups.

Despite the coupling to the SSM, the different SUSY breaking sectors themselves still interact only through SUGRA, so the analysis of the goldstini masses in the previous sections remains intact. Note that because of the mixing matrix from Eq. (12), there are still nontrivial couplings between SSM fields and the goldstini from the sequestered SUSY breaking sectors.

B. Induced Couplings between SUSY Breaking Sectors

If two or more SUSY breaking sectors have direct couplings to the SSM, a true deviation from the sequestered limit arises. To see how this happens, consider sectors 1 and 2, each of which couples to the SSM sector via an operator suppressed by Λ_1 and Λ_2 , respectively (see Fig. 2). Clearly, loops of SSM sector fields induce direct interactions between sectors 1 and 2, which may in turn modify the goldstini properties.

Direct interactions between SUSY breaking sectors can potentially modify the vacuum structure drastically so that SUSY breaking no longer occurs in some of these sectors. We assume that this is not the case, i.e. parameters take values such that the shift of the vacuum is small enough to preserve the essential structure of the sectors. (The parameter regions considered in later sections satisfy this condition.) It is then easiest to analyze the effect of direct couplings using the non-linear parameterization of Eq. (10), where x_i and F_i represent the values after the vacuum shift and η_i is the goldstino arising from sector i .

Since η_1 and η_2 have the quantum numbers conjugate to the generators of SUSY₁ and SUSY₂, respectively, they have a unit charge under the corresponding

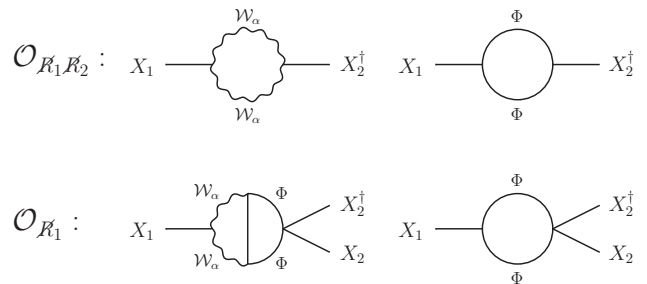


FIG. 3: Feynman diagrams which induce direct couplings between sectors 1 and 2. There is always at least one factor of $1/16\pi^2$ coming from a loop of SSM fields. Depending on the details of the underlying theory, there may be additional loop factors, for instance if the $\int d^2\theta X_1 \mathcal{W}^\alpha \mathcal{W}_\alpha / \Lambda_1$ coupling itself is generated at one loop.

R symmetries, $U(1)_{R_1}$ and $U(1)_{R_2}$, rotating these generators. Consequently, any deviations of the goldstini Majorana masses from $2m_{3/2}$ require an additional R -symmetry breaking transmission between sectors 1 and 2 beyond that provided by SUGRA through $m_{3/2}$. Since the setup considered here has tree-level direct couplings only between the SSM and SUSY breaking sectors, such a transmission must occur through the SSM sector.

The leading R -breaking transmitting couplings between a SUSY breaking sector and the SSM sector are given by the gaugino-mass and A -term operators, $\int d^2\theta X_i \mathcal{W}^\alpha \mathcal{W}_\alpha / \Lambda_i$ and $\int d^2\theta X_i \Phi^\dagger \Phi / \Lambda_i$, which may or may not exist depending on the properties of the SUSY breaking sector. Here, \mathcal{W}_α and Φ represent the gauge field strengths and chiral superfields of the SSM. Interactions of the form $\int d^4\theta X_i^\dagger X_i \Phi^\dagger \Phi / \Lambda_i^2$ do not provide necessary R -breaking transmission, unless X_i has a lowest component vev giving effectively A -term operators. For the remainder of this section, we will absorb any vev for X_i into the coefficients of the corresponding operators. Note that R -preserving operators can still play an important role in generating relevant effects when combined with operators that do transmit R -breaking.

We can characterize the induced couplings between sectors 1 and 2 according to whether they violate $U(1)_{R_1}$, $U(1)_{R_2}$, or both. These couplings are generated by the diagrams in Fig. 3, and have the form (after absorbing any vev for X_i into the operator coefficients)

$$\begin{aligned} \mathcal{O}_{\cancel{R}_1 \cancel{R}_2} &\approx \left(\frac{1}{16\pi^2} \right)^{n_{12}} \int d^4\theta X_1 X_2^\dagger + \text{h.c.}, \\ \mathcal{O}_{\cancel{R}_1} &\approx \left(\frac{1}{16\pi^2} \right)^{n_1} \frac{1}{\max\{\Lambda_1, \Lambda_2\}} \int d^4\theta X_1 X_2^\dagger X_2 + \text{h.c.}, \\ \mathcal{O}_{\cancel{R}_2} &\approx \left(\frac{1}{16\pi^2} \right)^{n_2} \frac{1}{\max\{\Lambda_1, \Lambda_2\}} \int d^4\theta X_2 X_1^\dagger X_1 + \text{h.c.}, \end{aligned} \quad (20)$$

where we have included the coefficients in the Lagrangian terms in the definitions of \mathcal{O} 's. We now consider each of these operators in turn.

C. Effects on Goldstini

If both $U(1)_{R_1}$ - and $U(1)_{R_2}$ -breaking effects exist and are transmitted, then the kinetic mixing operator $\mathcal{O}_{\hat{F}_1, \hat{F}_2}$ will arise. Note that $n_{12} \geq 1$, since there is always at least one loop of SSM fields involved in the diagram (if the gaugino mass operators themselves are generated at one loop, for instance as in gauge mediation, then $n_{12} = 3$).⁷ However, since this operator is separately holomorphic in sector 1 and sector 2 fields, it separately preserves SUSY₁ and SUSY₂ in the limit in which derivatively coupled terms are neglected—hence, this operator does not contribute to m_a .⁸ The only effect of $\mathcal{O}_{\hat{F}_1, \hat{F}_2}$ is to modify the kinetic term of ζ by an order $(1/16\pi^2)^{n_{12}}$ fraction, inducing $\delta m_a/m_a$ of the same size. If there are more than two sectors which couple to the SSM in this way, then kinetic mixings of this order will be generated among all the ζ_a .

If only $U(1)_{R_1}$ -breaking effects are transmitted, then $\mathcal{O}_{\hat{F}_1}$ is generated, where again $n_1 \geq 1$ because there is at least one loop of SSM fields. This operator yields a contribution to the goldstini masses

$$\begin{aligned} \mathcal{O}_{\hat{F}_1} &\supset \frac{(1/16\pi^2)^{n_1} F_2}{2 \max\{\Lambda_1, \Lambda_2\}} \left(\frac{F_2}{F_1} \eta_1^2 + \frac{F_1}{F_2} \eta_2^2 - 2\eta_1 \eta_2 \right) \\ &\rightarrow \frac{1}{2} \left(\frac{1}{16\pi^2} \right)^{n_1} \frac{F_{\text{eff}}}{\cos \theta \max\{\Lambda_1, \Lambda_2\}} \zeta^2, \end{aligned} \quad (21)$$

where in the last equation we have assumed that the only sectors breaking SUSY are sectors 1 and 2, and have plugged in for the mixing angles in Eq. (11). Obviously, an identical analysis can be performed when only $U(1)_{R_2}$ is broken.

If neither of $U(1)_{R_1}$ or $U(1)_{R_2}$ breaking is transmitted, the goldstini Majorana masses cannot deviate from $2m_{3/2}$. The goldstini, however, may still obtain Dirac masses with fermions of R -charge -1 . For instance, consider $\int d^4\theta X_1^\dagger X_1 S_2^\dagger S_2$, which is an R -symmetric coupling between a SUSY breaking field in sector 1 and a spectator field in sector 2 which does not have an F -component vev. For $\langle S_2 \rangle \neq 0$, this operator induces a Dirac mass between the goldstino, ζ , in X_1 and the fermionic component of S_2 . The effect from this class of operators, however, is generically smaller than that expected from $\mathcal{O}_{\hat{F}_1}$ and $\mathcal{O}_{\hat{F}_2}$ for natural values of $\langle S_2 \rangle \sim O(\sqrt{F_2})$.

The operators $\mathcal{O}_{\hat{F}_1}$ and $\mathcal{O}_{\hat{F}_2}$ can potentially produce large corrections to the goldstini masses. However, since they are suppressed by $\max\{\Lambda_1, \Lambda_2\}$, we find that in most cases these corrections are

$$\delta m_a \lesssim \left(\frac{1}{16\pi^2} \right)^n \tilde{m}, \quad (22)$$

where $n \geq 1$ and \tilde{m} is the scale for the SSM superpartner masses, which we have taken to be common for the gauginos and scalars. Therefore, if the gravitino mass is not substantially smaller than the superpartner masses, as in the case where $\Lambda_{1,2}$ are taken near the gravitational scale, then the relation $m_a = 2m_{3/2}$ will receive only small corrections. The situation is model dependent if the gravitino is much lighter. The corrected goldstini masses, however, are still significantly smaller than \tilde{m} , so that the SSM superpartners can decay into them.

The matrix V_{ia} , defined by Eq. (12), is determined to diagonalize the goldstini mass matrix

$$\mathcal{L} = -\frac{1}{2} m_{ij} \hat{\eta}_i \hat{\eta}_j + \text{h.c.} \rightarrow -\frac{1}{2} m_a \zeta_a^2 + \text{h.c.}, \quad (23)$$

where $\hat{\eta}_i = \eta_i - (F_i/F_{\text{eff}})\eta_{\text{long}}$ is the goldstini field with the eaten mode projected out, and

$$m_{ij} = 2m_{3/2} \delta_{ij} + \delta m_{ij}, \quad (24)$$

with δm_{ij} representing the effects from the operators in Eq. (20). At the zero-th order in $\delta m_{ij}/m_{3/2}$ expansion, V_{ia} is the $N \times (N-1)$ part of an orthogonal matrix preserving the first term of Eq. (24). Since the angles of this matrix are determined by a perturbation, δm_{ij} , on the unit matrix $2m_{3/2} \delta_{ij}$, they are typically of order unity.

Finally, we note that none of the operators discussed above affects the mass of the eaten mode, η_{long} . This is consistent with the general argument in Sec. IV A.

D. Other Corrections

We have seen that the corrections to the goldstini masses from induced interactions between SUSY breaking sectors are generically small. If there are tree-level direct couplings between these sectors, their effects can be studied similarly, following the analysis above. The goldstini masses are also corrected if there is a deviation from the assumption that SUSY is broken in the global limit. This effect is discussed briefly in the Appendix.

At loop level, the goldstini masses receive corrections from anomaly mediated effects, which exist even in the sequestered limit. Using the non-linear parameterization, we can calculate the corrections and find

$$\delta m_{ij} = -\gamma_i m_{3/2} \delta_{ij}, \quad (25)$$

where γ_i is the anomalous dimension of X_i defined by $d \ln Z_{X_i} / d \ln \mu_R = -2\gamma_i$.⁹ Naturally, these contributions

⁷ The loop factor may not exist if the SSM sector contains a singlet that directly mixes with SUSY breaking fields. We assume that such a singlet does not exist.

⁸ Incidentally, the same argument also shows that any operators of the form $\int d^4\theta f_1(X_1) f_2(X_2^\dagger)$ do not contribute to m_a .

⁹ If the X_i vev is nonzero in the basis where the X_i linear term vanishes in the Kähler potential, there is an additional contribution $\delta m_{ij} = \dot{\gamma}_i x_i^* m_{3/2}^2 / 2F_i$, where $\dot{\gamma}_i = d\gamma_i / d \ln \mu_R$. This contribution is generically much smaller than that in Eq. (25).

are loop suppressed.¹⁰ Note that the eaten mode, η_{long} , does not receive such a correction.

VI. INTERACTIONS WITH THE SSM SECTOR

In this section, we show how the goldstini couple to the SSM. As per usual, the gravitino couples to the SSM fields through its eaten goldstino component, η_{long} , whose interactions to a chiral multiplet take the form

$$\mathcal{L}_{\text{int}} \supset \frac{1}{F_{\text{eff}}} \left(\sum_i \tilde{m}_i^2 \right) \eta_{\text{long}} \psi \phi^\dagger, \quad (26)$$

where ψ and ϕ are the fermionic and bosonic components of a chiral superfield Φ of the SSM, and \tilde{m}_i is the soft mass contribution to this field from sector i . The interactions to a vector multiplet are given by

$$\mathcal{L}_{\text{int}} \supset -\frac{i}{\sqrt{2}F_{\text{eff}}} \left(\sum_i \tilde{m}_i \right) \eta_{\text{long}} \sigma^{\mu\nu} \lambda F_{\mu\nu}, \quad (27)$$

where λ is the gaugino, and \tilde{m}_i is the contribution to its mass from sector i .

The couplings of the uneaten goldstini to the SSM fields are different from those of the gravitino. We first consider those to chiral multiplets. The couplings of η_i to the SSM states can be obtained by using Eq. (10) in

$$\mathcal{L} = \sum_i \frac{1}{\Lambda_i^2} \int d^4\theta X_i^\dagger X_i \Phi^\dagger \Phi, \quad (28)$$

giving scalar mass contributions $\tilde{m}_i^2 = -F_i^2/\Lambda_i^2$. The interactions of the uneaten goldstini are then

$$\mathcal{L}_{\text{int}} \supset \frac{1}{F_{\text{eff}}} \sum_{i,a} \frac{\tilde{m}_i^2 V_{ia}}{r_i} \zeta_a \psi \phi^\dagger, \quad (29)$$

where $F_i \equiv r_i F_{\text{eff}}$ ($\sum_i r_i^2 = 1$), and we have used Eq. (12). In the case where there are only two SUSY breaking sectors, these interactions become

$$\begin{aligned} \mathcal{L}_{\text{int}} &\supset -\frac{1}{F_{\text{eff}}} (\tan \theta \tilde{m}_1^2 - \cot \theta \tilde{m}_2^2) \zeta \psi \phi^\dagger \\ &\approx -\left(\frac{F_2}{F_1^2} \tilde{m}_1^2 - \frac{1}{F_2} \tilde{m}_2^2 \right) \zeta \psi \phi^\dagger + \dots, \end{aligned} \quad (30)$$

where in the last equation we have assumed $F_1 \gg F_2$ and approximated F_{eff} by F_1 .

In the two sector case, it is useful to define the quantity

$$R = \left| \frac{\text{coefficient of } \zeta \psi \phi^\dagger}{\text{coefficient of } \eta_{\text{long}} \psi \phi^\dagger} \right|, \quad (31)$$

¹⁰ Corrections of similar size may also be induced by direct couplings between the SSM and SUSY breaking sectors. For example, loops of SSM states may generate holomorphic operators like $\int d^4\theta X_1 X_2$, giving corrections loop suppressed compared with $m_{3/2}$.

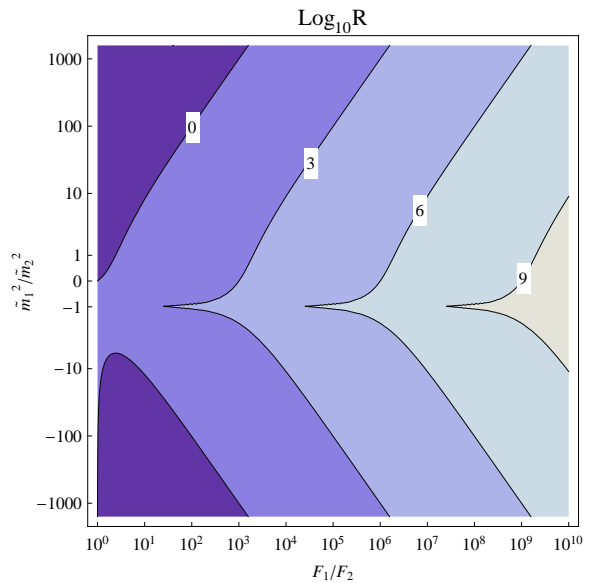


FIG. 4: A contour plot of R in Eq. (31) as a function of $\tilde{m}_1^2/\tilde{m}_2^2$ and F_1/F_2 . When $|\tilde{m}_1^2| \lesssim |\tilde{m}_2^2|$, R is greater than unity for a wide range of F_1/F_2 , so that the SSM sector fields couple more strongly to the uneaten goldstino than to the gravitino.

which characterizes the relative interaction strength of the SSM sector fields to the uneaten goldstino versus the gravitino. In Fig. 4, we plot R as a function of $\tilde{m}_1^2/\tilde{m}_2^2$ and $F_1/F_2 = \cot \theta$. We find that $R > 1$ more or less whenever $|\tilde{m}_1^2| \lesssim |\tilde{m}_2^2|$ —the SSM fields generically couple more strongly to the uneaten goldstino in this case.

The couplings of the goldstini to vector multiplets can be worked out similarly, and are given by

$$\mathcal{L}_{\text{int}} \supset -\frac{i}{\sqrt{2}F_{\text{eff}}} \sum_{i,a} \frac{\tilde{m}_i V_{ia}}{r_i} \zeta_a \sigma^{\mu\nu} \lambda F_{\mu\nu}. \quad (32)$$

If there are only two SUSY breaking sectors,

$$\begin{aligned} \mathcal{L}_{\text{int}} &\supset \frac{i}{\sqrt{2}F_{\text{eff}}} (\tan \theta \tilde{m}_1 - \cot \theta \tilde{m}_2) \zeta \sigma^{\mu\nu} \lambda F_{\mu\nu} \\ &\approx \frac{i}{\sqrt{2}} \left(\frac{F_2}{F_1^2} \tilde{m}_1 - \frac{1}{F_2} \tilde{m}_2 \right) \zeta \sigma^{\mu\nu} \lambda F_{\mu\nu} + \dots, \end{aligned} \quad (33)$$

where we have set $F_{\text{eff}} \approx F_1$ in the last line. As in the case of chiral multiplets, the couplings to the uneaten goldstino are generically stronger than those to the gravitino for $|\tilde{m}_1| \lesssim |\tilde{m}_2|$.

VII. COLLIDER PHENOMENOLOGY

Goldstini may be probed directly or indirectly at the LHC. In what follows, we consider a minimal setup in which SUSY is broken in two separate sectors, yielding a gravitino \tilde{G} and a single uneaten goldstino ζ . This scenario preserves most of the salient features of our general framework.

We focus our analysis on the regime in which $|\tilde{m}_1| \lesssim |\tilde{m}_2|$, so that the SSM fields couple more strongly to ζ than to \tilde{G} . This includes the case from Fig. 1 where a conventional SUSY breaking scenario is augmented by an additional, completely sequestered SUSY breaking sector with a higher SUSY breaking scale. Below we explore five classes of novel LHC signatures which can occur within our framework. We assume R -parity conservation throughout.

A. “Gravitino” with a Wrong Mass-Interaction Relation

Suppose that sector 2 which has F_2 ($\ll F_1$) gives masses to all the SSM superpartners. In this case, ζ couples more or less universally to all the SSM states, so that ζ looks like the “gravitino” when interpreted in the conventional framework. This apparent “gravitino”, however, has a wrong mass-interaction relation. Indeed, its interactions are controlled by F_2 (cf. Eqs. (30) and (33) when $|\tilde{m}_1| \lesssim |\tilde{m}_2|$), but its mass is controlled by F_1 (since $m_\zeta \simeq 2F_1/\sqrt{3}M_{\text{Pl}}$). This is different from the true gravitino, whose interactions and mass are controlled by a single parameter F_{eff} . Said another way, the goldstino has a mass which is a factor of $\simeq 2F_1/F_2$ larger than that of a conventional gravitino with a comparable interaction strength.

Suppose that ζ (and \tilde{G}) is lighter than all of the SSM superpartners, which we assume throughout this subsection. In this case, all the SUSY cascade will terminate with the lightest observable-sector supersymmetric particle (LOSP) decaying dominantly into ζ .¹¹ As in conventional gauge mediation, if $\sqrt{F_2} \lesssim 10^7$ GeV this decay may occur inside the detector; in particular, for small $\sqrt{F_2} \sim O(10 - 100$ TeV) it is prompt. Such a decay can provide a distinct signature at the LHC [8]. A unique aspect in our framework is that the mass of the escaping state can be significant, e.g. $\gtrsim O(10$ GeV) for $\sqrt{F_1} \approx O(10^9 - 10^{10}$ GeV), which cannot be the case in conventional gauge mediation. Therefore, if we can somehow measure a nonzero mass of this state, perhaps using methods similar to those discussed in Ref. [9], we can discriminate the present scenario from the usual one. These signals will be especially distinct if the LOSP is the bino (yielding two photons in the final state) or if a charged slepton LOSP decay leaves a displaced kink in the tracking detector. For massive escaping particles, such signals are hardly obtained in the conventional framework.¹²

If the LOSP is charged, then there can be a striking signature arising from a long-lived charged state. For $\sqrt{F_2} \gtrsim 10^6$ GeV, the LOSP may still live long enough that its mass and lifetime can be precisely determined by, e.g., velocity measurements and by observing decays of stopped LOSPs either inside a main detector [10] or in a proposed stopper detector [11]. Measurement of LOSP decays also allows us to determine the mass of the invisible state to which the LOSP decays, as long as it is larger than $O(10$ GeV). In fact, the charged LOSP arises naturally in many theoretical constructions. For example, the right-handed stau can easily be the LOSP if SUSY breaking is transmitted from sector 2 to the SSM sector via gauge or gaugino mediation. The LOSP may also be a selectron or smuon if there is a controlled source of flavor violation, which leads to a spectacular signal of monochromatic electrons or muons [12].

In the conventional scenario, the charged LOSP decays into the gravitino. Since the lifetime of the LOSP and the gravitino mass are related by F_{eff} , one can indirectly measure the Planck scale [13]

$$\begin{aligned} \Gamma_{\tilde{l} \rightarrow l \tilde{G}} &\simeq \frac{m_{\tilde{l}}^5}{16\pi F_{\text{eff}}^2}, & m_{3/2} &\simeq \frac{F_{\text{eff}}}{\sqrt{3}M_{\text{Pl}}} \\ \implies M_{\text{Pl}}^2 &\simeq \frac{m_{\tilde{l}}^5}{48\pi \Gamma_{\tilde{l} \rightarrow l \tilde{G}} m_{3/2}^2}, \end{aligned} \quad (34)$$

where we have adopted notation appropriate for a slepton LOSP. However, this is not the case if the LOSP instead decays into the uneaten goldstino ζ , since the goldstino mass and decay constant are controlled by separate parameters and thus a priori unrelated. Specifically, for $F_1 \gg F_2$, we will mismeasure M_{Pl} by a factor of $F_2/2F_1$ if we misinterpret ζ as a conventional gravitino

$$\begin{aligned} \Gamma_{\tilde{l} \rightarrow l \zeta} &\simeq \frac{m_{\tilde{l}}^5}{16\pi F_2^2}, & m_\zeta &\simeq \frac{2F_1}{\sqrt{3}M_{\text{Pl}}} \\ \implies M_{\text{Pl}}^2 &\simeq \left(\frac{2F_1}{F_2}\right)^2 \frac{m_{\tilde{l}}^5}{48\pi \Gamma_{\tilde{l} \rightarrow l \zeta} m_\zeta^2}, \end{aligned} \quad (35)$$

which would reveal that the particle to which the LOSP is decaying is not the gravitino.¹³

B. Gravitinoless Gauge Mediation

Thus far we have considered a case where ζ and \tilde{G} are lighter than the LOSP. However, since the masses of ζ and \tilde{G} are both controlled by the largest SUSY breaking scale F_1 , these states can be heavier than all the SSM superpartners even if F_2 (and the corresponding mediation scale Λ_2) is small. As a consequence, the LOSP may

¹¹ The goldstino ζ will decay further into the gravitino through intermediate SSM states. As we will see in Sec. VIII A, this decay is very slow, so that ζ can be regarded as a stable particle.

¹² The signals cannot be mimicked by a LOSP decay into the QCD axino either, since given an axion decay constant avoiding laboratory and astrophysical bounds, the decay occurs always outside the detector.

¹³ The LOSP decay product, however, may be the QCD axino \tilde{a} . Discriminating between ζ and \tilde{a} using the lifetime measurement will be difficult because values of $\Gamma_{\tilde{l} \rightarrow l \zeta}$ and $\Gamma_{\tilde{l} \rightarrow l \tilde{a}}$ mostly overlap in relevant parameter regions, especially if we allow the axion decay constant to be in the so-called anthropic range. The discrimination, however, may be possible by studying detailed structures of radiative three-body decays [14].

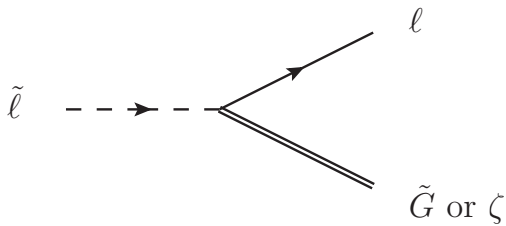


FIG. 5: If $F_1 \approx F_2$ and $\tilde{m}_1 \approx \tilde{m}_2$, then the SSM states couple to ζ and \tilde{G} with similar strengths. In particular, if ζ and \tilde{G} are lighter than all the SSM superpartners, then the LOSP decays into ζ or \tilde{G} with non-negligible branching ratios. This allows for the possibility of measuring the masses of both ζ and \tilde{G} , providing smoking gun evidence for multiple sector SUSY breaking.

be stable even if SSM superpartners obtain their masses primarily from a sector having low SUSY breaking and mediation scales.

This allows for a canonical gauge mediation spectrum without a light gravitino, and hence with neutralino dark matter. A scenario with similar phenomenology was considered before in Ref. [15]. In our context, it arises as a special case of the general framework of multiple SUSY breaking.

C. Measuring the “Two”

We have seen that the uneaten goldstino ζ may appear as a “gravitino” with a wrong mass-interaction relation, or may be heavier than the LOSP, making it irrelevant for collider experiments. Is there a situation in which we might directly observe both ζ and \tilde{G} and measure their detailed properties, in particular their mass ratio? The answer to this question is yes.

Suppose that two SUSY breaking sectors have comparable SUSY breaking strengths, $F_1 \approx F_2$, and contribute comparably to the masses of SSM superpartners, $\tilde{m}_1 \approx \tilde{m}_2$.¹⁴ In this case, ζ and \tilde{G} couple to SSM states with similar strengths. Therefore, if both ζ and \tilde{G} are lighter than all the SSM states, then the branching ratios of the LOSP to ζ and \tilde{G} are both non-negligible, as illustrated in Fig. 5 for the case of the slepton LOSP.

If $m_\zeta, m_{3/2} \gtrsim O(10 \text{ GeV})$, these masses can be determined by measuring the decays of long-lived charged LOSPs, using the same techniques as in Sec. VII A. This mass range corresponds to $\sqrt{F_1} \approx \sqrt{F_2} \approx O(10^9 - 10^{10} \text{ GeV})$, so that the LOSP is long lived. In the case that direct interactions between SUSY breaking sectors are small, this measurement will find two invisible

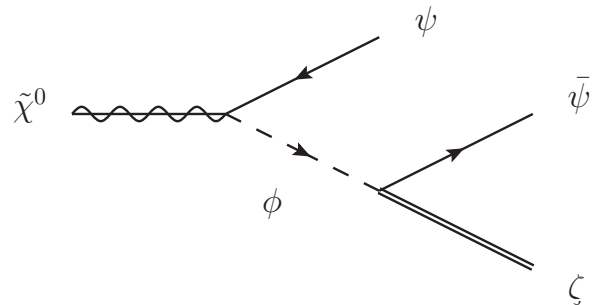


FIG. 6: If $F_1 \gg F_2$ and the SSM gaugino masses arise from sector 1 alone, then a bino-like LOSP can decay into ζ and two standard model fermions $\psi\bar{\psi}$ through an off-shell scalar ϕ , which is the superpartner of ψ . For $m_{\tilde{\chi}^0}^2 \ll m_\phi^2$ and $\tilde{m}_1^2 \ll \tilde{m}_2^2$, the branching fraction into each $\psi\bar{\psi}$ is entirely determined by the hypercharge of this field.

states $X_{1,2}$ whose masses satisfy

$$m_{X_1}/m_{X_2} \approx 2. \quad (36)$$

This would be an unmistakable signature of the uneaten goldstino ζ (or goldstini ζ_a with a degenerate mass), and hence smoking gun evidence for multiple sector SUSY breaking.

D. Difermions with Fixed Ratios

Distinct signatures may also arise if sectors 1 and 2 couple to the SSM in a more elaborate fashion. In particular, if one of these sectors preserves an (approximate) R symmetry, then the SSM gaugino masses are entirely generated by the other sector. This will affect the couplings of the SSM states to ζ , and can substantially change phenomenology.

Consider a situation that the two sectors have $F_1 \gg F_2$ and contribute comparably to the scalar masses, but that the gaugino masses arise solely from sector 1. This is true if sector 2 preserves an R symmetry. In this setup, the SSM scalars couple strongly to ζ , while the gauginos do so only very weakly. Therefore, if the LOSP is a bino-like neutralino, it decays either via $\tilde{\chi}^0 \rightarrow Z\zeta, h\zeta$ through its Higgsino fraction, or via $\tilde{\chi}^0 \rightarrow \zeta\psi\bar{\psi}$ through the off-shell SSM scalar ϕ which is the superpartner of a standard model fermion ψ (see Fig. 6). If $\tilde{\chi}^0$ has a significant Higgsino fraction, $\gtrsim O(0.1)$, and its decay into Z or h is not kinematically suppressed, then the former modes dominate. In this case the signature would look like the Higgsino LOSP decaying into ζ , even if the LOSP is bino-like.

If the above conditions are not met, the three-body decay $\tilde{\chi}^0 \rightarrow \zeta\psi\bar{\psi}$ dominates. In the limit that $m_{\tilde{\chi}^0}^2 \ll m_\phi^2$, the amplitude of this decay is proportional to $Y\tilde{m}_2^2/(\tilde{m}_1^2 + \tilde{m}_2^2)$, where Y is the hypercharge of ψ/ϕ and $\tilde{m}_{1,2}^2$ are the contributions to the ϕ mass-squared from each sector. Interestingly, for $\tilde{m}_1^2 \ll \tilde{m}_2^2$, the dependence on the ϕ mass

¹⁴ Such a situation may naturally be realized if environmental selection acts on superpartner masses through the requirement on the weak scale, and the two SUSY breaking sectors have comparable mediation scales, e.g., around the string scale.

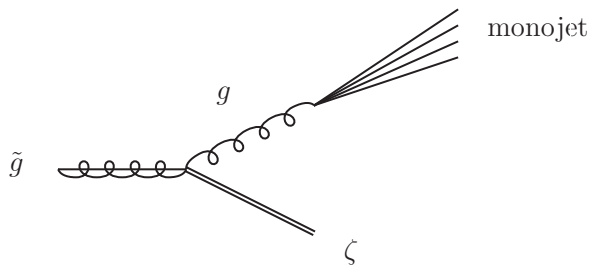


FIG. 7: If the squarks are sufficiently heavy, then the dominant gluino decay channel may be $\tilde{g} \rightarrow g\zeta$, which appears as a displaced gluino decaying into a monojet recoiling off of missing energy. If ζ is not at the bottom of the SUSY spectrum, then the decay of the ζ will produce a second displaced vertex.

drops out completely due to a cancellation between the propagator and the vertex factor. Therefore, in this parameter region, the ratios to various final states $\psi\bar{\psi}$ are entirely fixed by Y , giving

$$q\bar{q} : b\bar{b} : t\bar{t} : e\bar{e} : \mu\bar{\mu} : \tau\bar{\tau} \simeq 44 : 5 : 17 : 15 : 15 : 15, \quad (37)$$

where $q = u, d, s, c$. (There is also a completely invisible mode to neutrinos, and the $t\bar{t}$ mode may have a kinematic suppression. If $m_{\tilde{\chi}^0} > m_\zeta + 2m_h$, then $\tilde{\chi}^0 \rightarrow \zeta hh$ is also possible, whose rate depends on the masses of the Higgs/Higgsino.) This provides a unique signature of the setup considered here. Note that the decay of $\tilde{\chi}^0$ may also occur with a displaced vertex, since the $\tilde{\chi}^0$ lifetime can be long in some regions of parameter space.

E. Displaced Monojets

Another spectacular signal may arise if the SSM scalars are much heavier than the gauginos, as in split SUSY [16]. In particular, suppose that sector 2 provides weak scale masses to all of the SSM superpartners, while sector 1 does so only for the scalars—this can easily occur if sector 1 preserves an R symmetry. We also assume that the scalar masses from sector 1 are much greater than the weak scale.

If $m_\zeta < m_{\tilde{g}}$ and the squark masses are sufficiently large, $m_{\tilde{q}}^2 \gtrsim F_2/4\pi$, then the gluino prefers to decay directly into ζ and a gluon instead of cascade decaying through an off-shell squark. While $\tilde{g} \rightarrow g\zeta$ will generically be slow, for $\sqrt{F_2} \lesssim 10^7$ GeV it may occur within the detector. This gives a distinct signal of a displaced gluino decaying into a monojet recoiling off of missing energy (see Fig. 7).

Furthermore, if ζ is not at the very bottom of the superpartner spectrum, it will further decay into lighter SSM states. If the initial gluino decay occurs within the detector, then the ζ decay will also likely occur within the detector. This provides a spectacular signature of a secondary displaced vertex corresponding to the decay of

the uneaten goldstino ζ .¹⁵

VIII. COSMOLOGY

As one might expect, goldstini cosmology is not very dissimilar from gravitino cosmology. However, there are important differences arising from the fact that, unlike the gravitino, the goldstini have masses and couplings which are parametrically unrelated. This affects cosmology especially when these fields are lighter than the LOSP, which we will focus in this section.

As with the collider signatures in the previous section, we focus on the case of two SUSY breaking sectors with $|\tilde{m}_1| \lesssim |\tilde{m}_2|$. We also assume that deviations from the sequestered limit are small: the uneaten goldstino ζ has a mass $m_\zeta \simeq 2F_1/\sqrt{3}M_{\text{Pl}}$ and couplings to SSM fields proportional to $1/F_2$.

We assume “standard” cosmological history throughout this section. Many of the constraints discussed below can be avoided if we deviate from this assumption, e.g., if there is late time entropy production at temperature significantly below the weak scale.

A. Goldstini are Cosmologically Stable

If the goldstino ζ is lighter than the LOSP, it decays into the gravitino via $\zeta \rightarrow \tilde{G}\psi\bar{\psi}$, where ψ is a standard model fermion (arguments similar to the ones below will also hold for decays into photons). As we will see, this is longer than the age of the universe, so we can treat both goldstino and gravitino as stable particles.

In the conventional SUSY picture, low energy theorems dictate that the contact interaction $\tilde{G}\tilde{G}\psi\bar{\psi}$ is controlled by E^4/F_{eff}^2 , where E is a typical energy scale of the reaction [17]. While a complete description of goldstini low energy “theorems” is beyond the scope of this work, we note that $\zeta\tilde{G}\psi\bar{\psi}$ also scales like E^4/F_{eff}^2 , albeit with a prefactor that depends on \tilde{m}_i and F_i . Consequently, the width of the goldstino is given parametrically by

$$\Gamma_{\zeta \rightarrow \tilde{G}\psi\bar{\psi}} \approx \frac{1}{128\pi^3} \frac{m_\zeta^9}{F_{\text{eff}}^4} \left(\frac{F_1}{F_2} \frac{\tilde{m}_2^2}{\tilde{m}_1^2 + \tilde{m}_2^2} \right)^2. \quad (38)$$

The shortest reasonable lifetime is then

$$\tau_{\zeta \rightarrow \tilde{G}\psi\bar{\psi}} \approx 10^{22} \text{ sec} \left(\frac{\sqrt{F_2}}{100 \text{ TeV}} \right)^4 \left(\frac{100 \text{ GeV}}{m_\zeta} \right)^7, \quad (39)$$

so the goldstino is cosmologically stable. In theories of multiple sector SUSY breaking, decay transitions among the goldstini will take even longer, since they are nearly degenerate in mass.

¹⁵ While the signal of displaced monojets may be mimicked by conventional gauge mediation models with the gluino LOSP, the signal of a secondary displaced vertex cannot.

B. Late Decaying LOSP

Late decays of the LOSP to the goldstino will produce electromagnetic and/or hadronic fluxes which can alter the abundances of light elements and ruin the successful predictions of big bang nucleosynthesis (BBN) [18]. To safely evade such bounds, one either needs a small relic density of LOSPs, or the LOSP must have a lifetime shorter than ~ 100 sec.

For a conventional gravitino, BBN typically imposes a severe constraint $m_{3/2} \lesssim (10^{-2} - 1)$ GeV [19], where the precise values depend on the identity, mass, and abundance of the LOSP. In our case, however, the mass and coupling strengths of the uneaten goldstino are parametrically unrelated. Thus, the LOSP decay rate to the goldstino is a factor of $(F_1/F_2)^2$ greater than what one would expect for a comparable mass gravitino. Said another way, the goldstino behaves like a “gravitino” to which the LOSP decays faster than it should. Note that the usual LOSP to gravitino decay is now irrelevant, since the LOSP will primarily decay into the goldstino. As a consequence, a goldstino (and gravitino) mass in the range of $(1 - 100)$ GeV is easily compatible with BBN constraints in wide regions of parameter space.

C. Overproduction in the Early Universe

Another issue of a stable goldstino is that it may be overproduced in the early universe. For a comparable mass, the goldstino couples more strongly to the SSM states than the gravitino. This property has helped to avoid the BBN problem, as discussed above, but may hurt the overproduction problem. (We will see a way to sidestep this conclusion in the next subsection.)

Suppose that sector 2 provides sizable contributions to all of the SSM superpartners. The goldstino will then couple to the SSM much like a conventional gravitino. As in usual gravitino cosmology [20], the bound from overproduction is avoided for $m_\zeta \lesssim 0.2$ keV, since then the relic goldstino abundance from early thermal plasma is sufficiently small.¹⁶ For larger goldstino masses, there are upper bounds on the reheating temperature T_R in order for the relic goldstino not to overclose the universe.

It is relatively straightforward to translate the usual bounds for a gravitino, \hat{T}_R^{\max} , into corresponding bounds for an uneaten goldstino, T_R^{\max} . Since ζ has interaction strengths controlled by F_2 , its *number* density n_ζ is (approximately) the same as that one would have computed for a gravitino with $m_{3/2} = F_2/\sqrt{3}M_{\text{Pl}}$. The *energy* density $m_\zeta n_\zeta$, however, is larger than that of a gravitino with

the same mass by $m_\zeta/(F_2/\sqrt{3}M_{\text{Pl}}) = 2F_1/F_2$, implying

$$T_R^{\max}(m_\zeta, F_2) = \frac{F_2}{2F_1} \hat{T}_R^{\max} \left(m_{3/2} = \frac{F_2}{\sqrt{3}M_{\text{Pl}}} \right). \quad (40)$$

Note that this expression is not valid if T_R is sufficiently, typically $O(10)$, smaller than the superpartner mass scale, since then processes of goldstino generation are not active. Using the result for the standard gravitino scenario [22], we then find¹⁷

$$T_R^{\max} \approx 100 \text{ GeV} \left(\frac{1 \text{ GeV}}{m_\zeta} \right) \left(\frac{\sqrt{F_2}}{10^8 \text{ GeV}} \right)^4, \quad (41)$$

for $T_R^{\max} \gtrsim O(100 \text{ GeV})$; for $T_R \lesssim O(100 \text{ GeV})$, the bound disappears. The bound of Eq. (41) can also be written as $T_R^{\max}(m_\zeta, F_2) = (F_2/2F_1)^2 \hat{T}_R^{\max}(m_{3/2} = m_\zeta)$, so for $F_2 \ll F_1$ the reheating bound for the uneaten goldstino is significantly stronger than that for a comparable mass gravitino.

D. Goldstini Dark Matter

As we have seen, the constraint from BBN is avoided if the LOSP lifetime is sufficiently short, corresponding to

$$\sqrt{F_2} \lesssim (10^8 - 10^9) \text{ GeV}. \quad (42)$$

Then if T_R saturates the bound of Eq. (41), $T_R \simeq T_R^{\max}$, the uneaten goldstino will comprise all of dark matter. (Here we have assumed that the ζ abundance generated by possible late LOSP decays is small.) The required reheating temperature, however, is generically small in this case.

The strong bound of Eq. (41) on the reheating temperature was obtained by assuming that ζ couples to all the SSM states with the strengths $\approx 1/F_2$. However, this need not be the case. Consider, for example, that sectors 1 and 2 contribute comparably to the SSM scalar masses, but the gauginos obtain masses only from sector 1. This is the setup considered in Sec. VII D, and occurs naturally if sector 2 preserves an R symmetry. In this case, ζ couples to the scalars with the strengths $\approx 1/F_2$, but to the gauginos with $\approx F_2/F_1^2$, which are much weaker for $F_2 \ll F_1$.

The absence of strong ζ -gaugino interactions drastically changes the constraint from overproduction, since the standard reheating bound, Eq. (41), is dominated by ζ production from scattering involving the gluino. In the absence of these interactions, the constraint comes from

¹⁶ Structure formation, however, provides a stronger bound of $m_\zeta \lesssim O(10 \text{ eV})$ in this case [21].

¹⁷ This bound assumes a gluino mass of 1 TeV. In general, T_R^{\max} scales as $m_{\tilde{g}}^{-2}$.

ζ production from early scalar scatterings and decays, which will be satisfied for

$$\sqrt{F_2} \gtrsim 10^8 \text{ GeV} \left(\frac{m_\zeta}{1 \text{ GeV}} \right)^{1/4}. \quad (43)$$

Therefore, if Eqs. (42) and (43) are simultaneously satisfied, and if the LOSP is a scalar, then the constraints from both BBN and ζ overproduction can be avoided even for very large T_R .¹⁸ Whether this is indeed possible, however, will require a more detailed analysis because of $O(1 - 10)$ uncertainties in our estimates of the constraints.

If $\sqrt{F_2}$ saturates Eq. (43), the generated ζ can comprise all of dark matter without any additional contributions. Assuming that Eq. (42) is satisfied, the bound on T_R comes only from the usual gravitino overproduction, which is rather weak if $m_\zeta \simeq 2m_{3/2}$ is not much smaller than the weak scale, e.g. if $\sqrt{F_1} \approx (10^9 - 10^{10}) \text{ GeV}$. If $\sqrt{F_2}$ satisfies but does not saturate Eq. (43), then the ζ abundance must be dominated by late LOSP decays in order for ζ to be dark matter [23]:

$$\Omega_\zeta \simeq \frac{m_\zeta}{m_{\text{LOSP}}} \Omega_{\text{LOSP}}, \quad (44)$$

where Ω_{LOSP} is the fractional contribution of the LOSP to the critical density if it did not decay into ζ . Since Ω_{LOSP} is controlled by the standard WIMP parametrics, so is Ω_ζ if m_ζ is not significantly below m_{LOSP} .

IX. DISCUSSION

The hypothesis of single sector SUSY breaking has by and large dictated the standard picture of SUSY phenomenology at colliders and in cosmology. In the conventional scenario, the (only) goldstino is eaten by the gravitino, whose mass and coupling strength to SSM fields are inextricably and sometimes problematically related.

Motivated by top-down considerations, we have relaxed this underlying assumption and considered the possibility that a multiplicity of sectors break SUSY, yielding a corresponding multiplicity of goldstini. Intriguingly, *even when these additional sectors are completely sequestered from the SSM, this can have a drastic effect on LHC collider phenomenology*. Ultimately this occurs because the gravitino eats a linear combination of the goldstini, and in a curious twist on the conventional narrative, what would have been our gravitino is replaced by a linear combination of the uneaten goldstini.

A key result of this paper is that all of the uneaten goldstini receive an irreducible and universal mass $m_a = 2m_{3/2}$ from SUGRA effects, as long as SUSY is broken

in the global limit. As a consequence, the SSM fields can have sizable couplings to the goldstini, whose masses and decay constants are a priori unrelated. This greatly expands the realm of phenomenological possibilities. In particular, we considered a number of novel collider signatures, including anomalous neutralino and slepton decays, gravitinoless gauge mediated spectra, and monojet signals from (multiple) displaced vertices.

A true smoking gun signature of multiple sector SUSY breaking will exist if a charged LOSP has sizable branching ratios to both the gravitino and at least one goldstino. In this case, the mass ratio between the gravitino and goldstino may be accurately measured in a stopper detector, and a ratio of 2 would give dramatic evidence towards the scenario considered in this paper.

There are many possible directions for future work. While we have concentrated on the scenario where each SUSY breaking sector is F -term dominated, there is of course the possibility that one or more sectors experience D -term or “almost no-scale” SUSY breaking. In the latter case, there is significant mixing between gravitational modes and SUSY breaking fields, and as previewed in the Appendix, the goldstini masses can deviate significantly from $2m_{3/2}$. Moreover, while most of the phenomenological analyses in this work have focused on the two sector case for simplicity, it would be interesting to complete a more thorough analysis of the case of multiple goldstini. Finally, we hope to explore more fully the cosmological implications of this large class of theories.

Acknowledgments

We thank N. Arkani-Hamed, A. Arvanitaki, N. Craig, S. Dimopoulos, D. Freedman, M. Schmaltz, and D. Shih for interesting discussions. The work of C.C. and Y.N. was supported in part by the Director, Office of Science, Office of High Energy and Nuclear Physics, of the US Department of Energy under Contract DE-AC02-05CH11231, and in part by the National Science Foundation under grants PHY-0555661 and PHY-0855653. J.T. is supported by the U.S. Department of Energy under cooperative research agreement DE-FG0205ER41360.

Appendix A: Explicit SUGRA Calculation

The goldstini mass spectrum derived in Sec. IV can also be derived by explicit computation, using the SUGRA formalism of Ref. [5]. The simplest case to consider is N sequestered sectors labeled by i that each contain only a single light chiral multiplet X_i . That is, we assume that any other multiplets in sector i have a supersymmetric mass term and can be integrated out of the effective SUGRA Lagrangian. In particular, this means that all moduli must be stabilized in the supersymmetric limit.

We start from a Kähler potential and superpotential

¹⁸ If the LOSP is a gaugino, the dominant decay is the three-body decay mode from Sec. VIID, which faces more stringent BBN constraints because of phase space suppression.

of the sequestered form [24]

$$K = -3M_{\text{Pl}}^2 \ln \left(\frac{-1}{3M_{\text{Pl}}^2} \sum_i \Omega^{(i)}(X_i, X_i^\dagger) \right), \quad (\text{A1})$$

$$W = W_0 + \sum_i W^{(i)}(X_i), \quad (\text{A2})$$

where each $\Omega^{(i)}$ and $W^{(i)}$ is only function of a single X_i . Here, W_0 is a constant that must be tuned to make the cosmological constant zero, and we can take W_0 to be real without loss of generality. It is convenient to define the modified Kähler potential

$$G = \frac{K}{M_{\text{Pl}}^2} + \ln \frac{W}{M_{\text{Pl}}^3} + \ln \frac{W^*}{M_{\text{Pl}}^3}, \quad (\text{A3})$$

and its derivatives $G_i = \partial_i G$, $G_{j^*} = \partial_{j^*} G$, $g_{ij^*} = \partial_i \partial_{j^*} G$, where

$$\partial_i \equiv M_{\text{Pl}} \frac{\partial}{\partial X_i}, \quad \partial_{j^*} \equiv M_{\text{Pl}} \frac{\partial}{\partial X_j^\dagger}. \quad (\text{A4})$$

The Kähler metric g_{ij^*} and its inverse $g^{ij^*} = (g^{-1})_{ji}$ can be used to raise and lower indices, such that $G^i = g^{ij^*} G_{j^*}$. With this notation, the scalar potential is

$$V = M_{\text{Pl}}^4 e^G (G_i G^i - 3). \quad (\text{A5})$$

The condition for vanishing cosmological constant (and hence flat space) is

$$G_i G^i = 3, \quad (\text{A6})$$

and the minimum of the potential satisfies

$$\partial_i V = 0, \quad \partial_{j^*} V = 0. \quad (\text{A7})$$

After SUSY is broken, one linear combination of the fermionic components of X_i is the true goldstino and is eaten to form the longitudinal component of the gravitino

$$\eta_{\text{long}} = \frac{1}{\sqrt{3}} G_i \psi^i. \quad (\text{A8})$$

The gravitino mass is

$$m_{3/2} = M_{\text{Pl}} e^{G/2}. \quad (\text{A9})$$

In unitary gauge, the remaining fermions have a quadratic Lagrangian of the form

$$-i \tilde{g}_{ij^*} \bar{\psi}^j \bar{\sigma}^\mu \partial_\mu \psi^i - \frac{1}{2} m_{ij} \psi^i \psi^j - \frac{1}{2} m_{i^* j^*}^* \bar{\psi}^i \bar{\psi}^j, \quad (\text{A10})$$

where \tilde{g} is the Kähler metric with the true goldstino direction removed. The mass matrix is

$$m_{ij} = m_{3/2} \left(\nabla_i G_j + \frac{1}{3} G_i G_j \right), \quad (\text{A11})$$

where $\nabla_i G_j = \partial_i G_j - \Gamma_{ij}^k G_k$ depends on the Christoffel symbol Γ_{ij}^k derived from the Kähler metric. Note that the direction corresponding to the eaten goldstino has a zero mass eigenvalue (assuming vanishing cosmological constant). The remaining $N-1$ uneaten goldstini masses can be determined by the physical mass-squared matrix

$$M^2 = A A^*, \quad A_i^{j^*} = m_{ik} g^{kj^*}, \quad (\text{A12})$$

where A^* is the complex conjugate of the matrix (not the Hermitian conjugate). In A , it is possible to use g instead of \tilde{g} since the true goldstino direction is zeroed out by m . Note that for the mass-squared matrix M^2 (unlike for m), we need not assume the X_i have canonically normalized kinetic terms.

The key assumption of this paper is that SUSY is broken in the global limit $M_{\text{Pl}} \rightarrow \infty$. Moreover, we assume that any mixing between the chiral multiplets X_i and the gravity multiplet is a subdominant effect, meaning that at the minimum of the potential

$$\epsilon_i \equiv \sqrt{\frac{1}{3M_{\text{Pl}}^2} \frac{\partial_i \Omega \partial_{i^*} \Omega}{\partial_i \partial_{i^*} \Omega}} \ll 1. \quad (\text{A13})$$

This corresponds to the assumption that there are no large linear terms in the Kähler potential, and in particular implies that Polonyi-like fields must have vevs $\langle X_i \rangle \ll M_{\text{Pl}}$.

It is now a straightforward exercise to calculate the eigenvalues of M^2 as a series expansion in ϵ_i . Using Eqs. (A1) and (A2), one finds

$$A_i^{j^*} = \delta_i^{j^*} \left(2m_{3/2} + \frac{\partial_i V}{\partial_i W} \right) e^{2i\theta_i} - \frac{2}{3} m_{3/2} G_i G^{j^*} + O(\epsilon_i), \quad (\text{A14})$$

where

$$\theta_i = \arg(\partial_i W). \quad (\text{A15})$$

By the condition in Eq. (A7), the $\partial_i V$ term in Eq. (A14) vanishes, and because the uneaten goldstini are all orthogonal to η_{long} , the $G_i G^{j^*}$ term is irrelevant. The θ_i phases in A are also irrelevant, since $M^2 = A A^*$. So as advertised, one finds that the $N-1$ uneaten goldstini all have masses of $2m_{3/2}$ with corrections of order ϵ_i .

One can also use the mass-squared matrix M^2 to calculate the eigenvalues for more general scenarios where ϵ_i is not small. One amusing example is to consider $N-1$ sectors with $\epsilon_i \ll 1$, and an additional ‘‘almost no-scale’’ sector with arbitrary ϵ_N but $W^{(N)} = 0$. In that case, one can show that of the $N-1$ goldstini, one is massless to all orders in ϵ_i (it only gets a mass proportional to $\partial_N W$). The other $N-2$ goldstini get a mass

$$2m_{3/2} \left(\frac{1}{1 + \epsilon_N^2} \right) + O(\epsilon_i). \quad (\text{A16})$$

Note that when $\epsilon_N = 0$, this reduces to the previous result, since in that limit X_N is simply an extra massless

mode that does not contribute to SUSY breaking. We will explore these and other cases in future work. As a preview, the result in Eq. (A16) is equal to $2F_C + O(\epsilon_i)$,

where F_C is the highest component of the conformal compensator.

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