Split supersymmetry and dark matter generation

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ABSTRACT: We analyze a simple Split Supersymmetry scenario where fermion masses come from anomaly mediation, yielding $m_s \sim 1000\,\mathrm{TeV}$, $m_{3/2} \sim 100\,\mathrm{TeV}$, and $m_f \sim 1\,\mathrm{TeV}$. We consider non-thermal dark matter production in the presence of moduli, and we find that the decay chains of moduli $\to$ LSPs and moduli $\to$ gravitinos $\to$ LSPs generate dark matter more efficiently than perturbative freeze-out, allowing for a light, LHC visible spectrum. These decaying moduli can also weaken cosmological constraints on the axion decay constant. With squark masses of order $1000\,\mathrm{TeV}$, LHC gluinos will decay millimeters from their primary vertices, resulting in a striking experimental signature, and the suppression of Flavor Changing Neutral Currents is almost sufficient to allow arbitrary mixing in squark mass matrices.

KEYWORDS: Cosmology of Theories beyond the SM, Supersymmetric Effective Theories, Supersymmetry Phenomenology.

Contents

1. Introduction

With the Large Hadron Collider only a few years away, a new division has arisen in the particle physics community — is the unnaturalness of the standard model a problem to be solved through clever model building, or is it a hint that physics at the TeV scale is different than we have imagined, so that the values of dimensionful parameters are determined by anthropic fine tuning? As physicists, we would prefer a 'physical' explanation of small numbers such as the cosmological constant and the higgs mass, but the success of Weinberg's prediction of the cosmological constant [1] and the level of fine tuning necessary even in our best theories forces us to take the anthropic argument seriously. We will be taking it seriously for the present work.

Yet the mind of an anthropically-motivated model builder is a troubled one, for it seems that this profound shift in our worldview has had the unfortunate side effect of putting our whole enterprise out of business. If the weak scale is determined by anthropic selection, perhaps there is no new physics at the TeV scale, and the LHC and even the ILC will be colossal disappointments.

There are two simple reasons to remain hopeful about new TeV scale physics: gauge coupling unification and dark matter. The former requires new particles charged under

the standard model gauge group, and the latter requires a new stable particle with the correct relic density. If we assume for maximal simplicity that these two problems have a common solution, then it is reasonable to assume that our dark matter candidate will have a weak interaction cross section. Furthermore, if we assume that the relic dark matter density is determined by perturbative freeze-out, then we can expect new physics at the TeV scale, tightly constrained by gauge coupling unification. Some of these assumptions may be wrong, but that is not necessarily a bad thing — the minimal model that follows from this reasoning will probably be invisible to the LHC [22].

A very popular set of models that result from this methodology fall under the heading of Split Supersymmetry [2], where the scalar superpartners necessitated by supersymmetry are very heavy, while the fermions, protected by chiral symmetry, lie near the TeV scale. The purpose of the present work is to show that in one particularly simple model of Split Supersymmetry, perturbative freeze-out is not the dominant mechanism for generating dark matter. However, the mechanism that will replace it is more efficient, allowing for an even lighter, more LHC-visible spectrum. This seems like good news, because in standard Split Supersymmetry scenarios [3, 18], dark matter is expected to be too heavy to be seen at the LHC; however we will see that in our case there is a danger of overclosure. One of the most elegant scenarios is ruled out, but we explore several ideas that can rescue it, and the scenario of Moroi and Randall [13] remains a viable possibility.

To understand our mechanism, we first need to explain the spectrum of the model. As shown in [6, 19], wherever supersymmetry is broken, there will be visible sector supersymmetry breaking effects from Anomaly Mediation. Furthermore, the methodology of effective field theory requires that we include in the Lagrangian all higher dimension operators allowed by symmetry, with appropriate suppression by inverse powers of the cutoff. When supersymmetry is broken in a hidden sector, F and D term VEVs in these operators will induce supersymmetry breaking in the visible sector. These two effects are completely generic, requiring no theoretical gymnastics, so a model where supersymmetry breaking arises only in this way would be particularly elegant. This is the model that we will analyze; it was studied for related reasons in [18].

With Anomaly Mediated Supersymmetry Breaking, gaugino masses are given by a loop suppression factor times $m_{3/2}$, so we expect that $m_{3/2} \sim 100 \, \text{TeV}$ so that new fermions are near a TeV. Contributions from higher dimension operators suggest that the scalar mass scale $m_s \sim 1000 \, \text{TeV}$. In supersymmetric theories, there generically exist weakly interacting moduli that get masses of order m_s after supersymmetry breaking. Now we see the complexity of the cosmology — not only do we have to account for the perturbative freeze out of the LSPs, but we also need to consider processes involving late-decaying moduli and gravitinos. These late decaying particles produce a great deal of entropy, potentially weakening cosmological constraints on the properties of axions.

We will see that with our mass spectrum, dark matter production is dominated by two potential decay chains: moduli → LSPs and moduli → gravitinos → LSPs. Only the first process is available for $m_\phi < 2m_{3/2}$,[1] whereas both decay chains are open for $m_\phi > 2m_{3/2}$

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[1]Of course this depends on the available gravitino-producing modulus decay channels — the strict limit

— thus the physics is qualitatively different for different values of the modulus mass. The first process was studied in [13], and it naturally leads to the correct relic density in our scenario when $m_\phi \lesssim m_{3/2}$. The most naive version of the second decay chain is too efficient [33 – 35], so we consider a variety of mechanisms to alleviate this problem, including symmetries in the modulus sector, scenarios with multiple moduli, KKLT type moduli, and supersymmetry breaking dominated by D-terms. Of course we do not need these mechanisms for a successful cosmology if $m_\phi \lesssim m_{3/2}$, but it is interesting to explore all reasonable options.

Our model also has an exciting LHC signature. The gluino must decay through a virtual squark, so with 1000 TeV squarks, it is very likely that the LHC will see displaced gluino vertices if gluino production is kinematically allowed. In fact, we expect that gluinos will be copiously produced, since cosmology suggests a light spectrum.

We are also in an interesting region for Flavor Changing Neutral Currents. The strongest constraints from FCNC come from the ϵ parameter of the K-$\bar{\text{K}}$ system, which depends on the imaginary parts of the mass insertion parameters. Setting all mass insertions equal to a common value δ (there are no significant cancellations), we find from [7] that

$$\epsilon = 3 \times 10^{-3} \delta \left(\frac{1000 \text{ TeV}}{m_s} \right)^2 \tag{1.1}$$

Experiments constrain $\epsilon < 2 \times 10^{-3}$, so for $m_s \approx 1000$ TeV, we find that as long as $\delta \lesssim 1/2$, our scalar mass scale does not conflict with observations of FCNC. We also expect that electric dipole moments induced by new interactions may be visible at next generation experiments, as shown in [3], and our spectrum may help to explain neutrino masses, as examined in [18].

The outline of this paper is as follows. In section 2 we display the mass spectrum of our model. In section 3 we consider the cosmological implications of moduli, gravitinos, and LSP (Wino) dark matter. In section 4 we show that displaced gluino vertices will be a generic LHC signature for our model, and with section 5 we conclude.

2. The mass spectrum

We obtain the fermion mass spectrum

$$m_{3/2} = \frac{\langle W \rangle}{M_p^2} = \langle F_\phi \rangle \sim 100 \text{ TeV} \tag{2.1}$$

$$m_{\tilde{g}} = \frac{\beta(g)}{2g} \langle F_\phi \rangle \sim 1 \text{ TeV} \tag{2.2}$$

from anomaly mediated supersymmetry breaking [6, 19]. The details of hidden sector supersymmetry breaking are unimportant because the hidden sector only communicates with standard model particles through the auxilary field of supergravity and through higher

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is $m_\phi < m_{3/2}$ — but under reasonable assumptions $\phi \to 2\psi_{3/2}$ is the only available decay mode.

dimension operators in the Kahler potential. We assume that the scalars acquire masses from operators such as

$$L \supset \int d^4\theta \frac{1}{M_{GUT}^2} X^\dagger X Q^\dagger Q \tag{2.3}$$

where X is a hidden sector field, because these terms cannot be forbidden by any symmetry. We expect these operators to arise generically when we integrate out GUT or string scale particles, making them a factor of 10 - 100 times larger than $m_{3/2}$. Thus we can estimate that

$$m_s = m_{3/2} \frac{M_{pl}}{M_{GUT}} \sim 1000 \text{ TeV}, \tag{2.4}$$

completing our rough picture of the mass spectrum.

It is also important to consider the generation and the effects of the μ and $B\mu$ terms. Will μ and $B\mu$ be near m_s, $m_{3/2}$, or at the TeV scale? First, there can be contributions from R-symmetry breaking spurionic operators $X = 1 + \theta^2 m_s$ such as

$$L \supset \int d^4\theta \epsilon X^\dagger X H_1 H_2 \tag{2.5}$$

$$L \supset \int d^4\theta \epsilon X^\dagger H_1 H_2 \tag{2.6}$$

where the factor of ϵ is included to parameterize a possible approximate PQ symmetry. This gives rise to $B\mu \sim \epsilon m_s^2$ and $\mu \sim \epsilon m_s$. In most split supersymmetry scenarios we use R-symmetry to prohibit such spurionic contributions because they produce a nearly degenerate spectrum. However, in our scenario $m_{3/2}/m_f \sim 100$, which can be conveniently explained as a loop factor from anomaly mediation. Thus such operators may be permitted. If supersymmetry breaking does not lead to R-symmetry breaking we only have spurions[2] such as

$$L \supset \int d^4\theta \epsilon Y H_1 H_2 \tag{2.7}$$

with $Y = 1 + \theta^4 m_s^2$, which contribute to $B\mu$ only. Finally, there are contributions from the conformal compensator of supergravity

$$L \supset \int d^4\theta \epsilon \phi^\dagger \phi H_1 H_2 \tag{2.8}$$

where $\phi = 1 + \theta^2 m_{3/2}$, so that $\mu \sim \epsilon m_{3/2}$. Thus we find that $B\mu \approx \epsilon m_s^2$ in all cases, and either $\mu \approx \epsilon m_s$ or $\mu \approx \epsilon m_{3/2}$. We can take $\epsilon \approx 1/100$ to explain the top-bottom mass hierarchy, giving a μ term at the TeV scale.

The presence of the μ term modifies the gaugino masses [6, 19] so that

$$m_{\tilde{b}} = 8.9 \times 10^{-3} \left(1 - \frac{f(\mu^2/m_A^2)}{11}\right) m_{3/2} \tag{2.9}$$

$$m_{\tilde{w}} = 2.7 \times 10^{-3} \left(1 - f(\mu^2/m_A^2)\right) m_{3/2} \tag{2.10}$$

$$m_{\tilde{g}} = 2.6 \times 10^{-2} m_{3/2} \tag{2.11}$$

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[2]As explained in [3], this does not require that supersymmetry breaking comes from gauge superfields, but only that supersymmetry breaking is not accompanied by R-symmetry breaking.

where $m_A^2 \sin(2\beta) = B\mu$ and the function

$$f(x) = \frac{2x \ln(x)}{x - 1} \tag{2.12}$$

For small values of x (the expected case) this is a small effect. It increases the separation between the wino, bino, and gluino mass scales, but there is no real qualitative change in the spectrum. For large x we can obtain a very light Wino LSP, although we would need to abandon the PQ symmetry; this will be breifly considered in section 3.3.3.

3. Cosmology

Our variant of the Split Supersymmetry mass spectrum includes 100-$10,000$ TeV moduli field(s) ϕ and a 100 TeV gravitino, so we must check that these new ingredients do not disturb Big Bang Nucleosynthesis, and that the correct relic abundance of LSP dark matter candidates obtains. We begin with some general issues about our universe's history, and then we consider the possibility that our moduli fields weaken cosmological bounds on the axion. Finally, we explain the details of our proposed LSP generation mechanism.

We can think of our modulus field as though it were an inflaton — it begins with a VEV ϕ_0 that can be of order M_{pl}. Once the hubble constant decreases to $H \sim m_\phi$, the modulus begins to oscillate and its energy density red-shifts like that of matter. This occurs very early in the universe's history, with $T_{roll} \sim \sqrt{M_{pl} m_\phi} \sim 10^9$ TeV. At a time t_{eq} shortly after ϕ begins to roll (assuming ϕ begins with a Planckian VEV), $\rho_\phi = \rho_{Rad}$, and from this point until $t_{decay} = \Gamma_{tot}^{-1}$ the universe is modulus dominated. We take the modulus decay rate as[3]

$$\Gamma_{tot} = \frac{N}{16\pi} \frac{m_\phi^3}{M_{pl}^2} = \left(10^{-5} \text{ sec}\right)^{-1} \left(\frac{N}{10}\right) \left(\frac{m_\phi}{1000 \text{ TeV}}\right)^3 \tag{3.1}$$

where $M_{pl} = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. We can think of N as the number of available light decay modes, although it is strongly dependent on unknown, $O(1)$ couplings. We see that the modulus decays well before BBN. The decaying modulus will dilute particles left over from the original inflaton decay, reheat the universe, and re-populate the universe with gravitinos.

Any thermal gravitinos produced during a prior inflationary reheating are diluted by the entropy from the modulus decay, which is given by

$$\frac{s_{after}}{s_{before}} = \left(\frac{\rho_\phi(t_{decay})}{\rho_R(t_{decay})}\right)^{3/4} \sim \frac{M_{pl}}{m_\phi} \left(\frac{\phi_0}{M_{pl}}\right)^2 \tag{3.2}$$

where the ratio of the energy densities at the time of modulus decay simply comes from the two different equations of state. If we assume ϕ_0 is in the neighborhood of M_{pl}, then the

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[3]There is certainly a significant uncertainty because we do not know any of the $O(1)$ couplings involved. However, we do know that modulus decays to matter in the complex representation of a symmetry group [13, 33, 35] receive a suppression proportional to a power of m_{mat}/m_ϕ, which means that moduli decay dominantly to gauge bosons, gauginos, and possibly the heavy scalars.

entropy released is very large. Modulus decay reheats the universe to a temperature [11]

$$T_R \approx 1.1 g_*^{-1/4} \sqrt{\Gamma_{tot} M_{pl}} = 180 \text{ MeV} \times \left(\frac{N}{10}\right)^{1/2} \left(\frac{g_*}{10.75}\right)^{-1/4} \left(\frac{m_\phi}{1000 \text{ TeV}}\right)^{3/2} \qquad (3.3)$$

which is very low, so we certainly do not thermally regenerate heavy particles. To avoid disturbing BBN, we need $T_R \gg 1 \text{ MeV}$, so $m_\phi \gg 30 \text{ TeV}$. Similarly, the gravitino lifetime is [32]

$$\tau_{3/2} = 5 \times 10^{-2} \text{ sec} \left(\frac{100 \text{ TeV}}{m_{3/2}}\right)^3 \qquad (3.4)$$

so for $m_{3/2} > 60 \text{ TeV}$ gravitinos do not disturb BBN [12, 33] even if they are re-introduced by modulus decay. Note, however, that the lifetime is increased in split supersymmetry because the gravitino cannot decay to states involving the heavy scalars.

Dark matter is produced in two different ways: directly from decaying moduli, and indirectly from decaying gravitinos. Thermal relic gravitinos are diluted to negligible levels by modulus decay, so gravitinos can only arise from the decay of heavy moduli with $m_\phi > 2m_{3/2}$. Thus there are two qualitatively different types of moduli — those with $m_\phi > 2m_{3/2}$ and those with $m_\phi < 2m_{3/2}$. In the following subsection we consider direct modulus-LSP interactions, focusing on moduli with $m_\phi < 2m_{3/2}$, and commenting on heavier moduli. Then in the next subsection we consider gravitino production from moduli with $m_\phi > 2m_{3/2}$. It turns out that dark matter production from gravitinos is more efficient than direct production from modulus decay; this is essentially because the dark matter from gravitinos is produced later in the universe's history, so it accounts for a larger fraction of the universe's energy density.

3.1 Direct modulus decay to dark matter

We expect that the modulus will have generic string or Planck suppressed couplings to MSSM fields. For example, we may have interactions such as

$$\int d^2\theta \frac{\Phi}{M_{pl}} W_\alpha W^\alpha \quad \Longrightarrow \quad L \supset \frac{\phi}{M_{pl}} F_{\mu\nu}^2 + \frac{F_\phi}{M_{pl}} \lambda\lambda, \qquad (3.5)$$

where λ may be any gaugino[4]. If Φ has a supersymmetric mass, then $F_\phi \propto m_\phi \phi$, and both terms give large decay rates, but in other situations where m_ϕ only arises after supersymmetry breaking, we may have $\frac{dF_\phi}{d\phi} < m_\phi$ and the fermionic decay mode will be suppressed by the fermion mass. Due to R-Parity conservation, this implies that an order one fraction of the 'modulus particles' may eventually produce an LSP. The number density of 'modulus particles' relative to the entropy just after reheating is roughly

$$Y_\phi = \left(\frac{m_\phi}{T_R}\right)^{-1} \sim \left(\frac{m_\phi}{M_{Pl}}\right)^{1/2} \qquad (3.6)$$

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[4]Note that if Φ acquires a large F-component VEV, then this operator could give dangerously large gaugino masses, so we must assume that either F_ϕ is small or that the coefficient of this operator is suppressed.

If this were the whole story, we would overclose the universe. If the branching ratio to fermions is suppressed by m_f^2/m_ϕ^2 an acceptable dark matter density might be possible.

However, the LSPs produced by modulus decay pair annihilate [13] until their annihilation rate is less than H, after which they freeze out. Although the Wino's produced by the decaying modulus are ultra-relativistic, they acquire a thermal distribution with temperature T_R from interactions with the radiation bath. For instance, the process $\tilde{W}^0\nu \to \tilde{W}^0\nu$ proceeds more quickly than H for $T_R \gg 1\,\text{MeV}$, rendering the Winos non-relativistic. The annihilation cross section for Wino dark matter [13] into $W^\pm$ pairs is

$$\langle v_{rel}\sigma \rangle = \frac{g^2}{2\pi}\frac{1}{m_{LSP}^2}\frac{(1-x_W)^{3/2}}{(2-x_W)^2} \tag{3.7}$$

in the non-relativistic limit, with $x_W = m_W^2/m_{LSP}^2$. Thus the maximal dark matter relic density is reduced to

$$n_{LSP} \sim \left.\frac{3H}{2\langle v\sigma\rangle}\right|_{T_R} \sim \frac{T_R^2}{M_{pl}\langle v\sigma\rangle}. \tag{3.8}$$

Note that this will apply to any late decay that produces dark matter, including gravitino decay. Thus we find [13]

$$\frac{\rho}{s} \approx 4\times 10^{-10}\,\text{GeV} \times \left(\frac{10}{N}\right)\left(\frac{g_*}{10.75}\right)^{-1/4}\frac{(2-x_W)^2}{(1-x_W)^{3/2}}\left(\frac{m_{LSP}}{100\,\text{GeV}}\right)^3\left(\frac{100\,\text{TeV}}{m_\phi}\right)^{3/2} \tag{3.9}$$

where we have taken $g_* = 10.75$. We need $\rho/s \approx 4\times 10^{-10}\,\text{GeV}$ [36] to account for dark matter, so if there is a single light modulus with $m_\phi < 2m_{3/2}$, we can easily obtain the correct relic abundance [13]. Since the decay of this modulus will not produce gravitinos, we have a viable cosmology.

If $m_\phi \gg m_{3/2}$ then after modulus decay we are left with a negligable density of LSPs, so next we consider the consequences of gravitino production in scenarios with such a heavy modulus.

3.2 Moduli-gravitino interactions and indirect dark matter production

There are two potential channels in which modulus decay produces gravitinos,

$$\phi \to \tilde{\phi} + \psi_{3/2} \tag{3.10}$$

$$\phi \to 2\psi_{3/2} \tag{3.11}$$

where $\tilde{\phi}$ is the modulino, the superpartner of the modulus field. We do not expect the modulus to decay to a gravitino and another fermion because the other fermion would reside in the hidden sector, and we expect that hidden sector fields are very massive. For now we will assume that the modulino is very heavy, so that the first channel is also forbidden by kinematics — later we will see that relaxing this assumption does not help us to obtain a viable cosmology. In this section we consider the limit that $m_\phi \gg m_{3/2}$, and we begin with the two gravitino channel.

Gravitinos interact primarily through their longitudinal (helicity $\pm 1/2$) components, which correspond to the goldstino, the field that is 'eaten' in the superhiggs mechanism.

Thus the goldstone boson equivalence theorem becomes the goldstino equivalence theorem, which simply tells us that amplitudes involving longitudinal gravitinos can be computed using the goldstino instead. In the appendix it is shown that the only two derivative operator involving a scalar and two non-linearly realized, derivatively coupled goldstinos (longitudinal gravitinos) is

$$\frac{1}{F^\dagger}\partial_\nu s_\chi^\dagger \chi \sigma^\nu \bar{\sigma}^\mu \partial_\mu \chi - \frac{1}{F}\partial_\nu s_\chi \chi^\dagger \bar{\sigma}^\mu \sigma^\nu \partial_\mu chi^\dagger. \tag{3.12}$$

where the 'sgoldstino' field is

$$s_\chi = \sum_i \frac{F_i}{F}\phi_i, \tag{3.13}$$

and $F^2 = \Sigma_i F_i^2 + \Sigma_a D_a^2$ gives the goldstino decay constant. The sgoldstino is generally not a mass eigenstate — it is a blend of many different scalar fields. Thus the rate for moduli decay to goldstino pairs is determined entirely by the overlap of the mass eigenstate modulus with the sgoldstino.

In supergravity one expects higher dimension operators to be present in the Kahler potential, with suppression of order the GUT, string, or Planck scale. These operators tend to shift the VEVs of scalar fields and auxiliary fields. We would expect the effect on a modulus to be proportional to $1/m_\phi$ because the mass tends to stabilize the field, so on dimensional grounds we can estimate that

$$F_\Phi \sim \frac{\mu_{susy}^4}{m_\phi M_{pl}} \sim \frac{m_{3/2}}{m_\phi}F \tag{3.14}$$

for the auxiliary field of a heavy modulus. A more rigorous argument for this conclusion was given in [34], and it seems to be generically true in specific models in the absence of special symmetries[5]. Assuming that the ϕ field really is a mass eigenstate, this implies a rate for $\phi \to 2\psi_{3/2}$

$$\Gamma_{3/2} = \frac{C}{16\pi}\frac{m_\phi^3}{M_{pl}^2} \tag{3.15}$$

where C is a dimensionless constant. The branching fraction for modulus decay to gravitino pairs is therefore C/N.

In the approximation of instantaneous modulus decay, the ratio of modulus particles before the decay to entropy immediately after the decay is simply $Y_\phi = T_R/m_\phi$. Thus we estimate that

$$Y_{3/2} = \frac{C}{N}\frac{T_R}{m_\phi} = 6 \times 10^{-8}\frac{C}{\sqrt{N}}\left(\frac{m_\phi}{1000 \text{ TeV}}\right)^{1/2} \tag{3.16}$$

Now conservation of R-parity implies that each gravitino will eventually decay into at least one LSP, so we can approximate $Y_{LSP} = Y_{3/2}$. If the LSPs do not pair annihilate, then we find

$$\frac{\rho_{LSP}}{s} = 6 \times 10^{-6} \text{ GeV}\frac{C}{\sqrt{N}}\left(\frac{m_{LSP}}{100 \text{ GeV}}\right)\left(\frac{m_\phi}{1000 \text{ TeV}}\right)^{1/2} \tag{3.17}$$

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[5]It continues to be true when the modulus has a shift symmetry, although as we will discuss later there are exceptions in some models, such as specific versions of KKLT [35].

so we must take $C/\sqrt{N} \lesssim 10^{-4}$ to avoid overclosing the universe. This seems to be a very stringent constraint — it rules out the most elegant scenario with $C \sim 1$ – so in the next section we will consider whether a small C can be achieved naturally, or if other processes can reduce the relic density of LSPs. Note that the relic density is proportional to m_{LSP}, so if this scenario can be made viable, then we expect a light spectrum.

One might hope that the LSPs produced by gravitino decay pair annihilate, decreasing their abundance to within acceptable levels. We can use equation (3.8) together with the spectrum given in equation (2.9) to write the resulting LSP density in terms of $m_{3/2}$ alone [33, 34]

$$\frac{\rho_{ann}}{s} \approx 8 \times 10^{-9} \text{ GeV} \frac{(2-x_W)^2}{(1-x_W)^{3/2}} \left(\frac{m_{3/2}}{100 \text{ TeV}} \right)^{3/2}, \tag{3.18}$$

and this would require $m_{3/2} < 20 \text{ TeV}$ to obtain the correct relic density. Such a light gravitino would cause severe problems with BBN, and a Wino with $m_{LSP} < 50 \text{ GeV}$ would already have been detected. Thus if the gaugino masses are given by anomaly mediation (with $\mu \ll m_A$), then we cannot rely on Wino pair annihilation, and instead we need a mechanism for suppressing C if this scenario is to be viable. However, because this result seems tantalizingly close to giving the correct relic density, we will consider below whether deviations from the basic anomaly mediated predictions are possible.

3.3 Possibilities for suppressing moduli-induced gravitinos

3.3.1 Approximate symmetries

If the modulus is charged under an approximate symmetry, then its decay rate to gravitino pairs may be naturally suppressed. For this to make sense, the modulus cannot have a shift symmetry, so it cannot be a volume modulus as in KKLT, but it may be a shape modulus. It is also important that there exists a sector of light particles with the same symmetry properties, so that modulus decay to this sector is unsuppressed — otherwise the overall modulus decay rate would be small, but the branching fraction to gravitinos would be unchanged.

For instance, as a simple toy model we could consider

$$L = \int d^4\theta \left[X^\dagger X + \Phi^\dagger \Phi + \Phi_c^\dagger \Phi_c + \epsilon \frac{\Phi^\dagger X^2}{M_{pl}} + \text{H.C.} - \frac{(X^\dagger X)^2}{M^2} + \cdots \right]$$
$$+ \int d^2\theta \left[\Lambda + \mu^2 X + m_\phi \Phi \Phi_c \right] \tag{3.19}$$

where Φ and Φ_c have opposite charge under an approximate U(1) symmetry whose violation is parameterized by ϵ. In the limit that $\epsilon \to 0$, the modulus cannot decay to gravitinos by charge conservation, so for small ϵ the decay rate is suppressed. We could easily include a sector of light 'charged' particles so that the overall modulus decay rate is not suppressed. Whether or not such a situation could arise depends on the details of very high-energy physics, but there is no reason to expect that such a setup is impossible.

3.3.2 Light moduli

If the modulus is light, with mass close to $2m_{3/2}$, then the branching fraction to gravitinos will be phase-space suppressed. Thus parametrically we would have

$$C \propto \frac{\sqrt{m_\phi^2 - 4m_{3/2}^2}}{m_\phi}, \tag{3.20}$$

but obtaining $C \sim 10^{-3}$ in this way would require a large, poorly motivated fine-tuning.

A more interesting possibility is that there are two or more moduli with different masses. If the lightest modulus is lighter than $2m_{3/2}$, then it will not produce gravitinos, but its decay will dilute the gravitinos from earlier processes. If the two moduli have masses $m_\phi < 2m_{3/2} \ll m_\Phi$, then the lighter modulus begins rolling later in the universe's history, so that $\rho_\phi \approx \rho_\Phi$ while they are both oscillating. After Φ decays, the lighter field dominates the energy density of the universe until it in turn decays, releasing more entropy. Between the time of Φ decay and the time of ϕ decay, radiation red-shifts as R^{-4} but the energy in ϕ only red-shifts as R^{-3}, so ϕ decay dilutes heavy relics such as gravitinos by a factor

$$\Delta \approx \left(\frac{R_\phi}{R_\Phi} \right)^{3/4} \approx \left(\frac{T_{RH-\Phi}}{T_{RH-\phi}} \right)^{3/4} \approx \left(\frac{m_\Phi}{m_\phi} \right)^{9/8} \tag{3.21}$$

which could optimistically be as large as 1000. For instance, with $m_\Phi = 5 \times 10^4 \, \text{TeV}$, $m_\phi = 100 \, \text{TeV}$, $m_{3/2} = 60 \, \text{TeV}$, and $N \sim 10$, one would need $C \sim 1/20$ to obtain the correct relic abundance of dark matter (with these values direct production of dark matter from ϕ decay is not a problem). We will see that the presence of such a light modulus is also useful for alleviating cosmological bounds on the axion.

Finally, we could simply return to the original scenario of Moroi and Randall [13], with a single modulus with $m_\phi < 2m_{3/2}$. As they showed, the correct relic density can be obtained rather easily by relying on the pair annihilation studied in section 3.1 (in addition, they claimed that the modulus branching fraction to Winos was chirally suppressed, but this is false [33, 35]). It is certainly possible that moduli with masses of order m_s simply do not exist, but this is a UV sensitive issue.

3.3.3 Model dependence

Although models with $C \sim 1$ are generic, there are models where C is parametrically different. For instance, as shown in [35], in specific supergravity realizations of the KKLT scenario, one finds

$$C \sim \left(\frac{m_{3/2}}{m_\phi} \right)^2. \tag{3.22}$$

In this case $F_\phi \sim \frac{m_{3/2}}{m_\phi} F$ as usual, but 'ϕ' is not a mass eigenstate. The mass eigenstate modulus mixes with the supersymmetry breaking sector, so that it effectively has $F_{eigen} \sim \frac{m_{3/2}^2}{m_\phi^2} F$ giving a suppressed $\phi \to 2\psi_{3/2}$ decay rate. This result depends on specific properties of the supersymmetry breaking sector (the Polonyi field must be light, with mass of order $m_{3/2}$) that need not be true in more general KKLT-type scenarios, but these theories are

technically natural, and in some cases the terms that would give $C \sim 1$ can be forbidden by global symmetries.

We have been assuming that the gauginos get there mass from anomaly mediation alone, but as mentioned above, it is possible that the same operators

$$\int d^2\theta \frac{\Phi}{M_{pl}} W_\alpha W^\alpha \quad \Longrightarrow \quad L \supset \frac{F_\phi}{M_{pl}} \lambda\lambda \tag{3.23}$$

that allow moduli to decay to gauginos also contribute to the mass of the gauginos. If $F_\phi/M_{pl} \sim m_{3/2}^2/m_\phi$ as expected generically, then the Wino mass can be altered, and for a light gravitino and heavy modulus the Wino could become lighter. In that case, Wino pair annihilation could be sufficient to achieve the correct relic density; for instance, with $m_{\tilde{W}} = 70\,\text{GeV}$ and $m_{3/2} = 100\,\text{TeV}$, one obtains the correct abundance. Unfortunately this is a UV sensitive question, and it seems to require fine tuning to make the Wino light. Furthermore, there is the possibility of introducing CP violation through these types of interactions. Note that these operators are negligable when F_ϕ is small and the $\phi \to 2\psi_{3/2}$ decay rate is small, which is the domain of the section 3.3.1.

A better solution is to use a large μ term in equation (2.9) to make the Wino light. We would need to abandon the approximate PQ symmetry explaining the top-bottom hierarchy, but with $\mu \sim m_A$ one obtains a lighter Wino, making it possible to explain the dark matter relic abundance purely in terms of pair annihilation.

A final possibility is that supersymmetry breaking is dominated by D terms. As shown in the appendix, in this case C would be suppressed as

$$C \propto \left(\frac{F}{D}\right)^2. \tag{3.24}$$

This occurs because only the sgoldstino decays to gravitino pairs, and in the limit of D-breaking, there is no sgoldstino (instead the goldstino has a vector partner, which is not of interest here). As shown in [37] it is possible to obtain parametrically large D terms, and D-breaking naturally gives a split supersymmetry type spectrum. This possibility may warrant further investigation, as it may help to alleviate the cosmological moduli/gravitino problem in more general circumstances.

3.4 The $\phi \to \tilde{\phi} + \psi_{3/2}$ channel

For completeness, we consider the decay channel $\phi \to \tilde{\phi} + \psi_{3/2}$ [11], although we will see that it is dangerous, so we will need to forbid it kinematically. It comes from operators of the form

$$\begin{aligned} L \supset \int d^4\theta \frac{a}{M_{pl}} X \Phi^\dagger \Phi \\ = a\frac{m_{\tilde{\phi}}}{M_{pl}} \chi \tilde{\phi}\phi - \sqrt{3} a m_{3/2} m_{\tilde{\phi}} \phi\phi + \cdots \end{aligned} \tag{3.25}$$

where χ is the goldstino, $X = x + \theta\chi + \theta^2 m_{3/2} M_{pl}$ is the field that dominantly breaks supersymmetry, and a is an $O(1)$ coupling constant that must be relatively large so that

the modulino will be lighter than the modulus. As a complex scalar field, ϕ actually has two real modes, with masses

$$m_{\phi\pm} = \sqrt{m_{\tilde\phi}^2 \pm \sqrt{3}am_{3/2}m_{\tilde\phi}} \approx m_{\tilde\phi} \pm \frac{\sqrt{3}}{2}am_{3/2} \tag{3.26}$$

Generically, both modes will be present, but ϕ_- cannot decay to a modulino and a gravitino. If we require that $m_{\phi+}$ is greater than the sum of the modulino and gravitino masses, then we must have $a > (2 + m_{3/2}/m_{\tilde\phi})/\sqrt{3}$; if this relation is violated, then the decay will be kinematically prohibited.

With these assumptions, one can calculate the modulus decay rate and branching fraction into modulinos and gravitinos [11], obtaining

$$\begin{aligned}\Gamma &= \frac{a^2 m_\phi m_{\tilde\phi}^2}{8\pi M_{pl}^2}\left(1 - \frac{(m_{\tilde\phi} + m_{3/2})^2}{m_\phi^2}\right)^{3/2}\left(1 - \frac{(m_{\tilde\phi} - m_{3/2})^2}{m_\phi^2}\right)^{1/2} \\ &\approx \frac{3\sqrt{2}a^4 m_\phi m_{3/2}^2}{8\pi M_{pl}^2}\end{aligned} \tag{3.27}$$

for the decay rate into a gravitino and a modulino, where in the second line we take $m_{3/2} \ll m_\phi$. Thus we find a relic abundance [11]

$$Y_{3/2} \gtrsim 10^{-9}\left(\frac{m_{3/2}}{100 \text{ TeV}}\right)^2\left(\frac{1000 \text{ TeV}}{m_\phi}\right)^{3/2} \tag{3.28}$$

We know that $Y_{LSP} \approx Y_{3/2}$, so we obtain

$$m_{LSP}Y_{LSP} \gtrsim 10^{-7} \text{ GeV} \times \left(\frac{m_{LSP}}{100 \text{ GeV}}\right)\left(\frac{m_{3/2}}{100 \text{ TeV}}\right)^2\left(\frac{1000 \text{ TeV}}{m_\phi}\right)^{3/2} \tag{3.29}$$

for the dark matter energy density of the universe. We expect that $m_{LSP}Y_{LSP} = 4 \times 10^{-10}$ GeV [36] if the LSP accounts for all of the dark matter in the universe, so this decay mode will overproduce LSPs. Fortunately, we can prohibit this decay kinematically without introducing any fine-tuning.

3.5 Weakening cosmological bounds on the axion

The axion [26] decay constant F is bounded from below due to astrophysical constraints, and in generic cosmological scenarios it is bounded from above by the requirement that the axion does not overclose the universe. The moduli fields in our model decay when the universe is at a temperature near Λ_{QCD}, releasing a significant amount of entropy, and potentially relaxing constraints on F by diluting the axions. Throughout we will be considering the relic axion energy density from misalignment production, which is by far the dominant method of production for large F.

As shown in [27], diluting relic axions is a bit delicate. This is because there are three processes that need to be considered: a particle or field is decaying, releasing entropy and making the universe cool more slowly; H is decreasing, alleviating axion hubble friction;

and the axion mass is increasing, since m_a is strongly dependent on temperature because it arises from instanton effects. The axion only begins to roll as H drops below m_a, so if our modulus decays before this point then the entropy released does not decrease the final axion density. In general, particles decaying after the universe has cooled below $1\,\text{MeV}$ are dangerous to BBN, leaving a narrow window of $1\,\text{MeV} < T_R < T_h$ for reheating temperatures of decays that can dilute axions. Note that T_R should be interpreted as the approximate temperature of the universe after modulus decay — the temperature of the universe never increases.

First we estimate T_h, the reheating temperature above which the modulus does not dilute the axion density at all. As the modulus decay completes, the universe will be radiation dominated with temperature T_R, so $H \sim T_R^2/M_{pl}$ at this time. The axion has mass

$$m_a(0) = 13 \text{ MeV} \left(\frac{1 \text{ GeV}}{F}\right) \tag{3.30}$$

at temperatures less than about Λ_{QCD}/π and

$$m_a(T) = 0.1 m_a(0) \left(\frac{\Lambda_{QCD}}{T}\right)^{3.7}. \tag{3.31}$$

at temperatures above Λ_{QCD}/π. If $T_R = T_h$, then $3H(T_h) = m_a(T_h)$ just as the modulus decays, so we find that

$$T_h = 1.5 \text{ GeV} \left(\frac{10^{12} \text{ GeV}}{F}\right)^{0.18} \tag{3.32}$$

where we have taken $\Lambda_{QCD} = 200\,\text{MeV}$. We see that modulus decay will certainly dilute axions with F near the current bound, but for $F = 10^{15}\,\text{GeV}$, $T_h \sim T_R$ for the $1000\,\text{TeV}$ modulus field. However, we can do much better with a lighter modulus.

In order to obtain a relic axion energy density less than the current matter density in the universe, it was shown in [27] that we must have

$$F < 2 \times 10^{14} \text{ GeV} \left(\frac{100 \text{ MeV}}{T_R}\right)^{1/2} \left(\frac{m_a(T_a)}{m_a(0)}\right)^{1/2} \tag{3.33}$$

where T_a is the temperature of the universe when $m_a(T) = 3H$. We can estimate T_a as [27]

$$T_a = 200 \text{ MeV} \left(\frac{T_R}{100 \text{ MeV}}\right)^{0.26} \left(\frac{10^{15} \text{ GeV}}{F}\right)^{0.13} \tag{3.34}$$

thus as anticipated above, $m_\phi = 1000\,\text{TeV}$ leads to no improvement. However, if we consider moduli with $m_\phi \sim 100\,\text{TeV}$ as considered in section 3.1 and 3.3.2 then we obtain $T_R \sim 6\,\text{MeV}$, giving $F < 7 \times 10^{14}\,\text{GeV}$.

4. Displaced gluino vertices at the LHC

We know that cosmology favors a light anomaly mediated spectrum, so we can be justifiably optimistic that our model will be tested at the LHC. In fact, we expect that gluinos will

be copiously produced, and that gluino pairs will decay at secondary vertices separated by distances of order a millimeter or more. To begin to analyze this process, note that the resolution of the LHC's ATLAS detector [8] will be

$$\Delta d_0 = 11 + \frac{73}{(p_T/\text{GeV})\sqrt{\sin(\theta)}} \ \mu m \tag{4.1}$$

$$\Delta z_0 = 87 + \frac{115}{(p_T/\text{GeV})(\sin(\theta))^{\frac{3}{2}}} \ \mu m \tag{4.2}$$

Since gluinos are produced by QCD reactions such as $gg \to \tilde{g}\tilde{g}$, we can expect a large p_T and an order one θ. Thus we will need gluino decay vertices displaced by $\gtrsim 10 \ \mu m$ to have a chance of distinguishing them from primary vertices. There are four factors relevant to the displacement distance: the gluino production rate, the gluino lifetime, the relativistic time dilation, and the possibility of detecting gluinos from the tail of their distribution, which decays exponentially with distance from the primary vertex.

We will not attempt to compute the the gluino production rate precisely [21, 28 – 30], but we know that the answer is large, and our results will be very insensitive to the details. It was found in [9] that for $m_{\tilde{g}} \sim 350 \, \text{GeV}$, the LHC will produce about one gluino per second, or about 3×10^7 per year of operation. For $m_{\tilde{g}} \sim 2 \, \text{TeV}$, the LHC will produce at least a thousand per year. Thus we make the conservative assumption that the LHC will produce

$$N = 10^3 \left(\frac{2 \ \text{TeV}}{m_{\tilde{g}}} \right)^4 \tag{4.3}$$

gluinos per year.

Various groups [5, 20] have examined the decay of the gluino and computed its lifetime in Split Supersymmetry, including the relevant one-loop operator running from the SUSY breaking scale to the TeV scale. The result is

$$\tau_g = \frac{4 \ \text{sec}}{K} \left(\frac{m_s}{10^6 \ \text{TeV}} \right)^4 \left(\frac{1 \ \text{TeV}}{m_{\tilde{g}}} \right)^5 \tag{4.4}$$

where $K \lesssim 1$ is a weakly varying function of m_s, $m_{\tilde{g}}$, and $\tan(\beta)$. This is extended by a factor of approximately

$$\gamma \approx \frac{3 \ \text{TeV}}{m_{\tilde{g}}} \tag{4.5}$$

due to relativistic time dilation. We translate the gluino lifetime into a displacement by multiplying by a factor of $c\sqrt{1 - m^2/(3 \ \text{TeV})^2} \approx c$.

Putting these results together, we find that the number of gluinos with displacement in the interval $(R, R + dR)$ is

$$N(R) = 1000 \left(\frac{2 \ \text{TeV}}{m_{\tilde{g}}} \right)^4 \exp \left(-\frac{R}{D_0} \right) \frac{dR}{D_0} \tag{4.6}$$

where the canonical distance D_0 is given by

$$D_0 = (7.2 \ \text{mm}) \left(\frac{1 \ \text{TeV}}{m_{\tilde{g}}} \right)^6 \left(\frac{m_s}{1000 \ \text{TeV}} \right)^4 \tag{4.7}$$
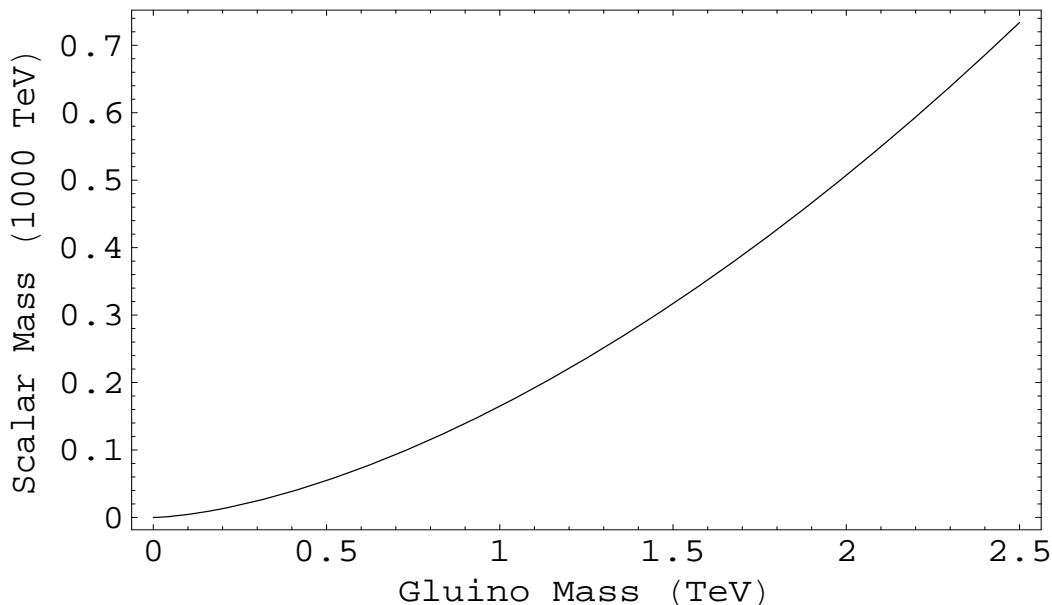
Figure 1: We will observe displaced gluino vertices at the LHC if m_s and $m_{\tilde{g}}$ lie in the region above the curve.

If we require the observation of at least a few displaced gluino vertices in one year's worth of LHC data, we find the limit

$$10^{-2} < \left(7 + 4\log\left(\frac{2\,\text{TeV}}{m_{\tilde{g}}}\right)\right)\left(\frac{1\,\text{TeV}}{m_{\tilde{g}}}\right)^6\left(\frac{m_s}{1000\,\text{TeV}}\right)^4 \tag{4.8}$$

on the mass parameters of the model. This result is only accurate to within a factor of a few, but this is unimportant because displaced gluino vertices will be visible for virtually the entirety of our anomaly-mediated parameter space, as seen in figure 1. For comparison, if we require the observation of 100 displaced gluino vertices each year at the LHC, then the digit 7 in the above equation changes into a 2, but this is the only alteration. Note that backgrounds for this signal from muons and b-quarks will be small due to the large jet energies, large missing energy, and especially because of the large vertex displacements. It was shown in the detailed study of [23] that simply using cuts on E_T and missing E_T, for $m_{\tilde{g}} < 1.4\,\text{TeV}$, the signal from R-Hadronized gluinos could be seen in 30 fb^{-1} of LHC data.

5. Conclusions

We have shown that a simple Split Supersymmetry spectrum based on anomaly mediation can satisfy cosmological constraints with a light spectrum. This is possible because the successive decays moduli $\to$ LSPs and moduli $\to$ gravitinos $\to$ LSPs generate dark matter very efficiently, in contrast with most anthropically motived models [3, 18, 22] based on perturbative freeze out, which often require 1-2 TeV LSP masses ([24] is an exception). Unfortunately, in one of the most elegant scenarios dark matter is overproduced, but we

found several mechanisms that can lead to the correct relic density, and the original mechanism of Randall and Moroi [13] remains a very viable possibility. We have also seen that our model will have a striking LHC signature — displaced gluino vertices.

As outlined in the introduction, there are only a few clues for would-be anthropic model builders. The two most tantalizing are probably gauge coupling unification and dark matter, and we must assume that these two issues are resolved in concert if we are to avoid an almost infinite set of possibilities for new physics. Here we have shown that retracting the usual assumptions about the generation of dark matter do not necessarily make models irrelevant for experimental collider physics, as we might have feared[6]. Furthermore, we noticed a fortuitous accident — as a side effect, our mechanism weakens cosmological constraints on axions, expanding the rather narrow window on the axion decay constant. With so little information about new physics, we should be appreciative when a small piece of the strong CP problem falls into our lap.

The success of anthropic predictions of the cosmological constant, combined with the derth of electric dipole moment signatures and new flavor changing neutral current effects, the smallness of the S and T parameters, and the lack of new physics at LEP point toward an anthropic solution to the two naturalness problems of high-energy physics and cosmology. If the world is supersymmetric at high energy, then the spectrum of new particles that we have considered is an excitingly predictive model for LHC physics.

A. Moduli-gravitino decay rates

These results were obtained in collaboration with Aaron Pierce and Jesse Thaler. Similar methods can be found in [31].

Consider an arbitrary, globally supersymmetric lagrangian $L(\phi, \psi, F, \lambda, A_\mu, D, \partial_\mu)$ that is a function of off-shell chiral multiplets and vector multiplets. Now assume that supersymmetry is spontaneously broken by some combination of D component and F component VEVs. The sgoldstino[7] is the linear combination

$$\phi_\chi = \sum_i \frac{F_i}{F} \phi_i, \tag{A.1}$$

where $F^2 = \Sigma_i F_i^2 + \Sigma_a D_a^2$ is the goldstino decay constant. Note that there is no reason to expect that ϕ_χ is a mass eigenstate. We will show that the only two derivative operator involving a scalar and two goldstinos is

$$\frac{1}{F^\dagger} \partial_\nu \phi_\chi^\dagger \chi \sigma^\nu \bar\sigma^\mu \partial_\mu \chi - \frac{1}{F} \partial_\nu \phi_\chi \chi^\dagger \bar\sigma^\mu \sigma^\nu \chi^\dagger. \tag{A.2}$$

Thus the rate of moduli decays to gravitino pairs is determined by the overlap of the mass-eigenstate moduli with ϕ_χ. As a simple corollary, we see that if supersymmetry were

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[6]However, it is worth noting that when dark matter is not generated by perturbative freeze out, we lose the elegant parametric prediction that the weak scale is the geometric mean of the Planck and Cosmological Constant scales [2]. The numerical prediction for the value of the cosmological constant is unchanged.

[7]If there are non-vanishing D component VEVs then there will also be a vector partner of the goldstino, but since we are only interested in moduli-goldstino interactions, we are ignoring these fields.

broken entirely by D terms, then there would not be a sgoldstino, so the rate for modulus decay to gravitino pairs would be zero.

To prove this assertion, we need to do a field redefinition to introduce the non-linearly realized goldstino χ. As in the conceptually simpler case of a goldstone boson, we can do this by choosing a 'vacuum alignment', and then parameterizing a supersymmetry transformation with the goldstino $\chi(x)$. For instance, if the theory consisted of a single chiral multiplet $\Phi = (\phi, \psi, F)$, we would represent $\Phi = e^{i\delta_\chi}(\phi, 0, F)$.

In the general case, when supersymmetry is linearly realized the goldstino is the combination

$$\psi_g = \sum_i \frac{F_i}{F} \psi_i + \sum_a \frac{D_a}{F} \lambda_a. \tag{A.3}$$

To isolate it, we begin by rotating the fermions ψ_i, λ_a into a mass eigenstate basis, so that the fermion lagrangian becomes

$$L_f = i\psi_g^\dagger \bar\sigma^\mu \partial_\mu \psi_g + \sum_\alpha \left[i\psi_\alpha^\dagger \bar\sigma^\mu \partial_\mu \psi_\alpha + \frac{m_\alpha}{2}(\psi_\alpha \psi_\alpha + \psi_\alpha^\dagger \psi_\alpha^\dagger) \right] + \cdots \tag{A.4}$$

where the elipsis denotes interaction terms. Note that the superpartners of the fermions ψ_α have no auxiliary component VEVs.

Now we perform a χ parameterized supersymmetry transformation on L, and then choose the 'vacuum alignment' condition $\psi_g = 0$. We are only interested in two operators, the interaction term mentioned above and the kinetic term for χ, which we require for canonical normalization. Both of these operators involve two goldstinos, so they can only come from $C\delta_\chi^2 A$ or $C\delta_\chi A \delta_\chi B$ where A, B, and C are some combination of fields in the lagrangian. But the first type of term would be $\frac{\Delta L}{\Delta A}\delta_\chi^2 A$, and this vanishes on-shell because it is proportional to the A equation of motion. Thus only terms of the second kind are relevant.

The supersymmetry transformation rules are the familiar

$$\delta_\chi \phi = \sqrt{2}\chi\psi \tag{A.5}$$

$$\delta_\chi \psi = i\sqrt{2}\sigma^\mu \chi^\dagger D_\mu \phi + \sqrt{2}\chi F \tag{A.6}$$

$$\delta_\chi F = i\sqrt{2}\chi^\dagger \bar\sigma^\mu D_\mu \psi \tag{A.7}$$

$$\delta_\chi \lambda = \sigma^{\mu\nu}\chi F_{\mu\nu} + i\chi D \tag{A.8}$$

$$\delta_\chi A_\mu = -i\lambda^\dagger \bar\sigma_\mu \chi + i\chi^\dagger \bar\sigma_\mu \lambda \tag{A.9}$$

$$\delta_\chi D = -\chi\sigma^\mu D_\mu \lambda^\dagger - D_\mu \lambda \sigma^\mu \chi^\dagger. \tag{A.10}$$

where D_μ is the appropriate gauge covariant derivative. The χ kinetic term can only come from the transformation of the ψ_g kinetic term, since the transformation of all other operators involve too many fields. Furthermore, it is remarkable that two derivative interactions between a scalar and two goldstinos also only come from the ψ_g kinetic term, because the other terms either give too many fields or vanish on the equations of motion. Thus as claimed we find

$$L \supset i\chi^\dagger \bar\sigma^\mu \partial_\mu \chi + \frac{1}{F^\dagger}\partial_\nu \phi^\dagger \chi \sigma^\nu \bar\sigma^\mu \partial_\mu \chi - \frac{1}{F}\partial_\nu \phi \chi^\dagger \bar\sigma^\mu \sigma^\nu \chi^\dagger + \cdots \tag{A.11}$$

after canonically normalizing χ, where the elipsis denotes other interactions, including a two goldstino interaction with the vector superpartner of the goldstino. Although from the transformation rules it might seem that χ has non-derivative interactions, these all cancel — χ has a shift symmetry, just like a goldstone boson. Thus yukawa couplings with χ cannot be generated.

We should note that this is an effective field theory, and it is only valid for $E^2 \ll F = m_{3/2}M_{pl}$. In most examples of interest (KKLT moduli, or moduli that get masses from supersymmetry breaking), the modulus mass easily satisfies this criterion. In any case, above this scale supersymmetry breaking is a small effect, and in place of a single goldstino one would need to consider the detailed dynamics of the hidden sector.

Acknowledgments

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