

Axion mediation

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ABSTRACT: We explore the possibility that supersymmetry breaking is mediated to the Standard Model sector through the interactions of a generalized axion multiplet that gains a F-term expectation value. Using an effective field theory framework we enumerate the most general possible set of axion couplings and compute the Standard Model sector soft-supersymmetry-breaking terms. Unusual, non-minimal spectra, such as those of both natural and split supersymmetry are easily implemented. We discuss example models and low-energy spectra, as well as implications of the particularly minimal case of mediation via the QCD axion multiplet. We argue that if the Peccei-Quinn solution to the strong-CP problem is realized in string theory then such axion-mediation is generic, while in a field theory model it is a natural possibility in both DFSZ- and KSVZ-like regimes. Axion mediation can parametrically dominate gravity-mediation and is also cosmologically beneficial as the constraints arising from axino and gravitino overproduction are reduced. Finally, in the string context, axion mediation provides a motivated mechanism where the UV completion naturally ameliorates the supersymmetric flavor problem.

KEYWORDS: Supersymmetry Phenomenology

Contents

1	Introduction	1
1.1	Motivation and UV completions	2
2	Effective field theory of supersymmetric axions	6
3	Visible sector soft terms	9
3.1	Sfermions	9
3.2	Gauginos and gaugino screening	11
3.3	Higher order contributions	13
4	Notable models	14
4.1	Split supersymmetry	14
4.1.1	Scalars and gauginos	14
4.1.2	EWSB and Higgsinos	17
4.2	Split families and ‘Natural SUSY’	18
4.3	Split gaugino masses	20
4.4	QCD axion mediation	20
5	Axino as goldstino and cosmological constraints	22
5.1	The axino and the goldstino	22
5.2	Cosmology	23
6	Discussion and conclusions	24

1 Introduction

If supersymmetry (SUSY) is realized at the TeV scale and is wholly or in part responsible for the solution to the hierarchy problem then, given experimental constraints, its low-energy realization is almost certainly more complex than has typically been assumed in much of the literature. For instance, the constrained minimal supersymmetric standard model (CMSSM), even with altered assumptions in the Higgs sector, seems unlikely to be close to the truth. If one closes one’s eyes to the question of a consistent and motivated UV realization, a number of low-energy spectra can be constructed which lead to a sufficient weakening of the LHC search limits so that SUSY is still relevant for the hierarchy problem. Assuming for definiteness that R-parity is conserved in the low-energy theory, so-called “natural”, “compressed” and “stretched” SUSY spectra have been argued to significantly reduce the LHC limits while maintaining relatively low fine-tuning [1–5] (possibly with a

modified NMSSM-like Higgs sector [6], either in the large λ limit [7–10], or with an altered R-symmetry structure [11]).¹

However, the best-known and most studied mediation mechanisms for supersymmetry breaking, namely gauge mediation, anomaly mediation, and minimal sequestered gravity mediation, while consistent with the severe indirect flavor constraints on supersymmetry breaking (by construction in the gravity case) have trouble realizing the modified spectra. It is therefore worthwhile to look further for motivated mediation schemes which lead to one or more of the variant spectra. In this paper we argue that an axion multiplet, possibly the QCD axion multiplet itself, is a natural candidate to mediate supersymmetry breaking and can lead to unusual and phenomenologically attractive patterns of soft terms.

The QCD axion provides the best known solution to the strong-CP problem and so is a very well-motivated extension to the Standard Model (SM) [14–16]. The axion mechanism requires a new scale $f \ll M_{\text{Pl}}$ at which Peccei-Quinn (PQ) symmetry is spontaneously broken, and results in a new light particle which interacts with the SM. In the supersymmetric extension of such models the axion multiplet can mediate supersymmetry breaking to the visible sector and the separation of scales of the axion and M_{Pl} ensures that axion mediation will dominate over uncontrollable gravity-mediated effects. Moreover, the axion sector itself provides a natural candidate for the origin of supersymmetry breaking. Thus, the QCD axion is a promising yet relatively unexplored candidate to dominate supersymmetry breaking.²

The structure of the paper is as follows: in the remainder of the Introduction, we discuss how a UV theory typically results in one or more axion multiplets, and mechanisms by which axions may gain significant F-terms. In section 2, we set our notation and define the effective interactions of an axion multiplet consistent with shift symmetry, paying particular attention to invariance under local chiral field redefinitions. Following this section 3 gives expressions for the soft terms generated in the MSSM sector when a generic axion multiplet gains the dominant F-term, including typical hierarchies of sfermion and gaugino mass terms. In section 4 we specialize the generic couplings to some well motivated cases which lead to non-standard spectra of soft terms, and briefly discuss the resulting phenomenology. Finally, in section 5 we show there are some cosmological advantages to such a set-up, and conclude in section 6.

1.1 Motivation and UV completions

In light of the strong CP problem and the necessity of a QCD axion, we note that there are a variety of models that can generate axions with couplings to SM fields. In particular, we consider both string models in which axion multiplets are the moduli multiplets of the underlying compactification, and purely field-theoretic models. For both possibilities we

¹Relaxing the assumption of R-parity conservation weakens LHC limits only by a modest amount, with gluinos often still heavier than a TeV; it may even increase the limits in the case of leptonic R-parity violation [12, 13]. Significantly reducing the LHC limits requires SUSY events with nearly no missing energy and few leptons, both of which appear in generic spectra from decays of charginos and tops.

²Models in which the axion couples to the SUSY breaking sector, and the messengers both mediate SUSY breaking and generate the anomalous QCD coupling were studied in [17].

discuss mechanisms by which the axion multiplet may acquire a large F-term and hence dominate the mediation of supersymmetry breaking to the visible sector.

In the field theory case an anomalous global $U(1)$ symmetry — the PQ symmetry — is broken and an axion arises as a pseudo-goldstone boson. In KSVZ models the axion sector is not coupled to the visible sector chiral multiplets, and the anomalous coupling between the axion multiplet and the standard model gauge groups is a result of integrating out new heavy vector-like pairs of chiral multiplets charged under the PQ symmetry and the Standard Model gauge groups [18, 19]. In DFSZ models the MSSM matter itself is charged under the PQ symmetry and generates the anomalous coupling [20, 21]. A realistic UV theory may be expected to be some mixture of these two cases. Traditionally QCD axion models have been considered in the ‘axion window’, $10^9 \text{ GeV} \lesssim f \lesssim 10^{12} \text{ GeV}$: the parameter regime in which both experimental search limits are satisfied (the lower bound) and axions are not more than the observed amount of dark matter for generic values of the axion field in the early universe (the upper bound).

In field theoretic models the axion will mediate supersymmetry breaking if a single sector spontaneously breaks a global symmetry leading to a Nambu-Goldstone boson (the axion) and simultaneously breaks supersymmetry such that the axion gains an F-term, F_A . It is quite plausible that one sector can accommodate both of these roles in global supersymmetry, as any sector that breaks a global symmetry and has no run-away directions will necessarily also break supersymmetry so that the Nambu-Goldstone boson can remain massless while the moduli direction gains a mass. This, in fact, is used as a criterion in the search for theories that break supersymmetry in strongly coupled regimes (for a review of dynamical supersymmetry breaking see, for example, [22] and the references therein).

Moreover, our requirement that the interactions of the axion dominate the mediation of supersymmetry breaking to the SM sector can be naturally accommodated in this scenario. For example, in an extra-dimensional context, if the supersymmetry breaking sector is physically sequestered from the SM sector, the multiplets that remain light are the dominant source of mediation. This fact is due to the exponential mass suppression of wave functions along the extra dimension, or in 4D language, the statement that heavy fields gain large anomalous dimensions and therefore have suppressed interactions with SM sector. The axion multiplet, protected by Goldstone’s theorem, is precisely such a light field: the axion remains essentially massless and the other fields in the multiplet only gain masses of order $\sqrt{F_A}$, while other multiplets can gain masses of order the axion decay constant f and for $\sqrt{F_A} \ll f$ (i.e. MSSM soft terms at the TeV scale) the axion dominates the mediation. A model of supersymmetry breaking where the low energy degrees of freedom include an axion has been studied for example in [23].

In the string context, the topological complexity of a typical compactification gives rise to many axion-like multiplets. Each non-trivial cycle of the internal manifold can lead to a state with axion-like couplings from the zero mode of an asymmetric tensor field integrated on the cycle. The number of such cycles is typically $\mathcal{O}(10^3)$ or as large as $\mathcal{O}(10^5)$, giving in principle a comparable number of axion-type fields [24, 25]. For instance, in type IIB string theory, the integral of a rank-4 antisymmetric tensor field C_4 over every independent 4-cycle Σ_i^4 of the internal manifold gives rise to an independent pseudo-scalar field a_i in

4-dimensions,

$$a_i = \frac{1}{2\pi} \int_{\Sigma_i^4} C_4. \tag{1.1}$$

In addition, the underlying Abelian gauge symmetry in 10D, $C_4 \rightarrow C_4 + d\lambda_3$, implies that each field a_i inherits a global ‘PQ’ shift symmetry. Generally speaking, for a 10D rank- n antisymmetric field, every independent n -cycle of the compactification gives rise to an independent pseudoscalar mode in 4D. Therefore, a plethora of 4D axion-like fields results.

Not all these potential axion-like modes survive to the low-energy theory. As we discuss below, depending on the process of stabilization moduli can be projected out of the light spectrum. Typically, however, many remain massless at the perturbative level, acquiring a (possibly ultra-light) mass due only to non-perturbative effects which violate the shift symmetry. Therefore a multitude of axions can in general survive in the low-energy theory.³

The ability of the axion to solve the strong-CP problem, i.e. set $\bar{\theta}_{CP}$ to less than 10^{-9} , requires that non-perturbative effects which explicitly break PQ symmetry are dominated by QCD dynamics and not other sources such as string instantons. Numerically, this suppression of string instantons translates to the requirement $S \gtrsim 200$ where $\Lambda^4 = \mu^4 e^{-S}$ with Λ the scale that appears in the axion potential and μ the UV energy scale, in this case the string scale. As argued by Svrcak and Witten [24] one then quite generally finds that the axion scale f is polynomially suppressed $f \sim M_{\text{pl}}/S$. Therefore the QCD axion in string theory may be expected to have $f \lesssim 10^{16}$ GeV. Even if the QCD axion is not responsible for mediation of SUSY breaking, once the existence of a suitable QCD axion is imposed on the string compactification, there is naturally a plentitude of light axion fields — the string axiverse — with $f \sim M_{\text{pl}}/S$ [25], and one or more of these can take part in SUSY-breaking dynamics and transmit this breaking to the visible sector.

Obtaining string axions in the axion window requires warped Randall-Sundrum-like throats in which the axion degree of freedom is IR localized [24, 29] and whose throat parameters generate an IR scale $f \lesssim 10^{12}$ GeV. Although such throats are not unlikely given our limited knowledge of realistic compactifications [30], it is another requirement that one must impose. As emphasized in [31–34], however, the upper bound of the axion window is sensitive to untested assumptions about early universe dynamics, and in particular the over-closure limit is easily relaxed by a mild environmentally-selected tuning of the post-inflationary mis-alignment angle of the axion field. We thus consider that axion scales in the string-motivated range, $f \sim 10^{16}$ GeV, are allowed and are particularly interesting to explore (for search possibilities in this regime see, e.g., [25, 33, 35]).

String axions may naturally gain an F-term in the process of moduli stabilisation: as we now briefly review, this is the classic scenario of soft terms from moduli fields in string compactifications (see, for example [36–42]). In heterotic models this process can occur through gaugino condensation: a hidden gauge group coupling runs strong at a scale Λ below the string scale, $\Lambda \ll m_S$. While gaugino condensation itself does not break supersymmetry, the couplings to the axion field in the gauge kinetic function can induce an F-term [43]. There is a well-known issue of stabilizing the resulting run-away

³ Such axiverse scenaria are discussed in a variety of UV contexts [26–28].

moduli direction towards infinite field values [44], which may be addressed, for example, through multiple gaugino condensates. It has been argued that M-theory compactifications may lead to an axion multiplet gaining the dominant F-term [45–47]. Alternatively, flux compactifications of IIB string theories can lead to moduli with F-terms, either in the process of stabilization as in large volume scenarios [48], or as a result of uplifting from a SUSY preserving anti-de-Sitter vacuum to a Minkowski vacuum, for example in KKLT compactifications [49] (a general review can be found in [50]).

Despite much work studying string theory compactifications, there is still considerable uncertainty in how realistic low energy behavior is realized. However, since the axion decay constant f must be below M_{pl} , the moduli multiplets that form axions can naturally dominate the mediation. Additionally, if the QCD axion originates from string theory, it is very natural that such a multiplet gains a large F-term: it was shown in [51] that if a string axion is to be the QCD axion, the associated modulus field must be stabilized non-supersymmetrically. On the other hand, if stabilization occurs supersymmetrically (though difficult to realize in explicit models), the QCD axion must come from a field theory sector which, as discussed, is also a good candidate to dominate supersymmetry breaking in the MSSM.

A final attractive scenario that can apply to either string or field theory constructions is if the sector that generates the axion also strongly stabilises the flat direction of the modulus but does not generate the dominant supersymmetry breaking of the theory. This can occur with no SUSY breaking since supergravity in AdS allows supersymmetric mass splittings within multiplets (before uplifting). Now if, in addition to this sector, the axion couples to a sector which experiences gaugino condensation at a scale $\Lambda_h \ll f$ such that

$$|\langle \bar{\lambda}_h \lambda_h \rangle| \sim \Lambda_h^3, \tag{1.2}$$

then a standard anomalous coupling $\int d^2\theta \frac{A}{f} W_\alpha W^\alpha$ (directly present or induced through fermion anomaly diagrams) will lead to an axion F-term

$$F_A \sim \frac{\langle \bar{\lambda}_h \lambda_h \rangle}{f}. \tag{1.3}$$

Unlike the case of heterotic compactifications, there is no danger of a run-away direction since the moduli direction is stabilized in the axion sector. In this model the axion multiplet naturally gains the dominant F-term, although the axion itself will not be light since $\Lambda_h > \Lambda_{QCD}$ and cannot act as the QCD axion.

Hence, while we do not select a unique possibility, the wide range of UV completions motivates our detailed study of the IR effects of SUSY breaking through an axion F-term. In particular, we consider the extent to which it is possible to realize non-minimal patterns of soft terms such as natural supersymmetry, non-universal gaugino masses, and split supersymmetry from the effective interactions of such a multiplet. A single axion multiplet coupled to the SM captures these effects, so in the spirit of minimality we will assume from now on that just one axion multiplet gains an F-term and its couplings are the dominant source of SUSY breaking in the MSSM. It should be kept in mind, however, that more than one axion multiplet may play a role in SUSY breaking and its mediation.

2 Effective field theory of supersymmetric axions

As previously discussed, in field theory constructions an axion arises as a pseudo-Nambu-Goldstone boson from a spontaneously broken anomalous global U(1) symmetry (the Peccei-Quinn symmetry), while in string models axions can appear as fundamental moduli fields which automatically have approximate shift symmetries. However, the origin of the axion is not important from an effective field theory perspective: we simply consider an axion to be a chiral multiplet, A , which respects a shift symmetry at the perturbative level and may acquire a small mass via non-perturbative effects. The axion superfield A contains the pseudoscalar axion a as well as the saxion s , axino ψ_A , and auxiliary component F_A ,

$$A = \frac{(s + ia)}{\sqrt{2}} + \sqrt{2}\theta\psi_A + \theta^2 F_A. \tag{2.1}$$

At low energies the global PQ symmetry is realized non-linearly as a shift symmetry

$$A \mapsto A + i\alpha f, \tag{2.2}$$

where f is the axion decay constant. The low energy theory may also contain matter charged under PQ; chiral multiplets Φ_i with PQ charge x_i transform as

$$\Phi_i \mapsto e^{ix_i\alpha}\Phi_i. \tag{2.3}$$

The general low energy effective theory for an axion multiplet has been discussed in [52]. We summarize the results here for completeness and to establish notation. The effective interactions in a supersymmetric theory can be written as a sum of the gauge-kinetic interaction $\Delta\mathcal{L}_G$, superpotential interactions $\Delta\mathcal{L}_S$, and Kähler potential interactions $\Delta\mathcal{L}_K$ in the Wilsonian effective action,

$$\mathcal{L}_{\text{int}} = \Delta\mathcal{L}_G + \Delta\mathcal{L}_S + \Delta\mathcal{L}_K. \tag{2.4}$$

The gauge-kinetic coupling

$$\Delta\mathcal{L}_G = - \sum_n \int d^2\theta \frac{C_n}{32\pi^2} \frac{A}{f} \text{tr}(W_\alpha^n W^{n\alpha}) + \text{h.c.} \tag{2.5}$$

defines the standard anomalous coupling to the MSSM gauge field strengths W_α^n , where n labels the gauge group G_n .⁴ While the term (2.5) is frequently considered to give rise to the leading interactions of the axion multiplet, there are additional super-potential and Kähler potential interactions in the presence of chiral fields Φ transforming under PQ. These couplings are *not* sub-leading and are in fact crucial for the invariance of physical observables under field redefinitions of the chiral multiplets.

Specifically, the renormalizable holomorphic superpotential couplings are given by

$$\Delta\mathcal{L}_S = \sum_{ijk} \int d^2\theta \left(m_{ij} e^{-(x_i+x_j)A/f} \Phi_i \Phi_j + \lambda_{ijk} e^{-(x_i+x_j+x_k)A/f} \Phi_i \Phi_j \Phi_k \right) + \text{h.c.}, \tag{2.6}$$

⁴We use a gauge kinetic normalisation $\int d^2\theta \frac{1}{4g^2} W_\alpha W^\alpha$.

while the Kähler couplings to leading order in $1/f$ are⁵

$$\Delta\mathcal{L}_K = \sum_i \int d^4\theta \left(1 + y_i \frac{(A + A^\dagger)}{f} + z_i \frac{(A + A^\dagger)^2}{f^2} \right) \Phi_i^\dagger \Phi_i. \quad (2.7)$$

The fields Φ_i may be light MSSM fields or additional heavy vector-like pairs, with supersymmetric masses m_{ij} and trilinear Yukawa couplings λ_{ijk} . The axion multiplet can have further interactions with one or more hidden sectors, and possibly a (small) mass arising from, for example, string instanton effects, but we will not need to be explicit about their details here.

There is a large number of parameters in (2.5)–(2.7): superpotential couplings x_i , Kähler couplings y_i, z_i , and anomalous couplings C_n , which together specify all axion supermultiplet couplings up to order f^{-2} . However, these parameters are not independent: under a chiral rotation

$$\Phi_i \mapsto e^{k_i A/f} \Phi_i, \quad (2.8)$$

they transform as

$$\begin{aligned} x_i &\mapsto x_i - k_i \\ y_i &\mapsto y_i + k_i \\ z_i &\mapsto z_i + y_i k_i + \frac{k_i^2}{2} \\ C_n &\mapsto C_n - 2 \sum_i k_i T_n^{\Phi_i}, \end{aligned} \quad (2.9)$$

where $T_n^{\Phi_i}$ is the Dynkin index of the chiral fields Φ_i under the group n . The transformation of C_n is due to the Konishi anomaly [53]. After such a redefinition the PQ symmetry is still realized by (2.2) but with the new values of charges x_i as given in (2.9). In particular, there is a set of combinations invariant under field redefinitions,

$$\begin{aligned} x_i + y_i \\ y_i^2 - 2z_i \\ C_n - 2 \sum_i x_i T_n^{\Phi_i} \\ C_n + 2 \sum_i y_i T_n^{\Phi_i}. \end{aligned} \quad (2.10)$$

For a theory with N chiral superfields $2N+3$ of these combinations are linearly independent; as we will see, soft SM sector masses will be proportional to these combinations when the

⁵We work with a field expansion such that the saxion has zero VEV, in contrast to much of the string literature which is concerned with moduli stabilization where the minimum of the modulus (i.e. saxion/axion) potential is not known initially. This explains the difference between our form of the Kähler potential and that commonly given in string models.

axion sector acquires an F-term. A UV theory determines a particular set of couplings x_i, y_i, z_i, C_n , often in terms of a much smaller number of charges or couplings, and the equivalence class of theories defined by the transformations (2.9) must lead to identical expressions for physical observables. The choice of basis can be made on grounds of clarity and convenience.

In anticipation of our study of the phenomenology of particular models, it is interesting to consider the values such parameters may typically take. Firstly the x_i occur as a result of the charges of fields under some group. Therefore these are expected to be either natural numbers or exactly zero, with exponential separations or non-integer values unlikely. They may be universal between different generations, depending on the underlying theory (for example, as discussed later due to brane localisation of matter), and additionally may be universal within a generation, or vary for different fields. For example in a GUT compatible model, within a generation the fields Q_L, u^c and e^c are expected to have equal x , which may be different to the value of x for d^c and L .

In our convention the parameters C_n are order one when generated at loop level in either field or string theory models. Like the coefficients x , these depend on the charges of fields under symmetries, hence are typically natural numbers. An alternative scenario occurs in some string models with a tree level coupling between the modulus multiplet and gauge fields, in which case C_n is $\mathcal{O}(32\pi^2)$.

Finally, the parameters y and z have a rather different physical interpretation. These are not PQ symmetry breaking, and not necessarily linked to charges under some group. In string models they may be related to the modular weights of the fields, in which case they are also typically natural numbers; however in the field theory case they may be generated by integrating out additional matter. Depending on the generation mechanism they may be universal between generations (for example if they occur through integrating out gauge fields), or alternatively may be able to vary over a very large range. For example if the wavefunctions of different generations are localized differently in a warped dimension, there could an exponential variation in the magnitude of these parameters. Alternatively if they occur by integrating out matter of mass of the Planck mass they could be of magnitude $f/M_{\text{pl}} \sim 1/100$ in the string axion case.

To clarify the meaning of a physical observable in the context of these basis dependent parameters, consider the divergence of the current associated to the PQ symmetry:

$$\partial_\mu J_{PQ}^\mu = \frac{g^2}{16\pi^2} C_{PQn} F^{a\mu\nu} \tilde{F}_{\mu\nu}^a. \tag{2.11}$$

Unlike the Wilsonian couplings C_n which are changed by the field redefinition (2.8), the current divergence and therefor the anomaly coefficient C_{PQn} are physical quantities and must be basis-independent. That this is so can be checked by computing the divergence of the PQ current given by the sum of the Wilsonian coupling C_n and the anomalous diagrams containing the chiral fields Φ ,

$$C_{PQn} = -C_n + 2 \sum_i x_i T_n^{\Phi_i}, \tag{2.12}$$

which is one of the invariant combinations of (2.10) as expected. To say it another way, a chiral rotation of the fields Φ changes the coupling C_n in the Lagrangian but the change is compensated by the shift in the couplings x_i which generate an anomaly through the well-known triangle diagrams.

Another example is the 1PI coefficient of the anomalous gauge-kinetic coupling, which determines for instance the axino-gaugino-gauge amplitude at leading order. Recall that by definition the 1PI action includes, at fixed loop order, effects from integrating out all momentum scales; any symmetries of the underlying theory must be manifest in the 1PI generating functional. For the effective interactions of the axion multiplet the 1PI coefficient at one loop is given by [52]

$$C_{1PI_n}(p) = -C_n - 2 \sum_{m_i^2 < p^2} y_i T_n^{\Phi_i} + 2 \sum_{m_i^2 > p^2} x_i T_n^{\Phi_i}. \quad (2.13)$$

The details of the calculation are not important for our work, but we note the answer is indeed invariant under the transformations (2.9). This occurs for a very similar reason to the invariance of C_{PQ_n} : Field rotations lead to changes in the Lagrangian coupling C_n which are cancelled by loop diagrams containing the matter fields with charges x_i . It is interesting to note that C_{1PI_n} is a function of momentum and the behavior depends critically on the masses of the particles concerned; we will encounter this mass dependence again in the SUSY-breaking gaugino masses in section 3.2.

3 Visible sector soft terms

In sections 1.1 and 2 we have discussed the motivation for axion mediation from the top-down perspective and established an effective theory of a supersymmetric axion with an emphasis on basis-independent physical observables. Now we turn to an analysis of the soft terms induced by an axion multiplet that participates in SUSY breaking dynamics and develops an F-term expectation value F_A . We calculate the sfermion and gaugino masses explicitly; many of our results can also be understood in an elegant way from analytic continuation. Ultimately we find the soft terms obtained can interpolate between traditional mediation mechanisms such as gauge mediation and less explored possibilities such as split supersymmetry.

3.1 Sfermions

Chiral matter fields feel SUSY breaking from the axion sector due to Kähler and superpotential couplings. Depending on the supersymmetric mass of the multiplet, we find the couplings x in the superpotential (2.6) and y, z in the Kähler (2.7) lead to sfermion masses proportional to one of two invariant combinations: $x + y$ or $y^2 - 2z$. For a single pair of fields Φ_1, Φ_2 with a supersymmetric mass term $m\Phi_1\Phi_2$, the mass matrix for the scalar components is given by

$$\mathcal{L} \supset - \begin{pmatrix} \phi_1^\dagger & \phi_2^\dagger \end{pmatrix} \begin{pmatrix} m^2 + (y_1^2 - 2z_1) \frac{F_A^2}{f^2} & (x_1 + x_2 + y_1 + y_2) m \frac{F_A}{f} \\ (x_1 + x_2 + y_1 + y_2) m \frac{F_A}{f} & m^2 + (y_2^2 - 2z_2) \frac{F_A^2}{f^2} \end{pmatrix} \begin{pmatrix} \phi_1 \\ \phi_2 \end{pmatrix}, \quad (3.1)$$

and associated masses of the scalar mass eigenstates

$$\begin{aligned}
 m_s^2 &= m^2 + \frac{1}{2} (y_1^2 - 2z_1 + y_2^2 - 2z_2) \frac{F_A^2}{f^2} \\
 &\pm \frac{F_A}{f} \sqrt{(x_1 + x_2 + y_1 + y_2) m^2 + \frac{F_A^2}{4f^2} ((y_1^2 - 2z_1) - (y_2^2 - 2z_2))^2}.
 \end{aligned}
 \tag{3.2}$$

where we denote the complex scalar component of Φ_i as ϕ_i and the θ^2 component as F_i . Clearly these masses are invariant under the basis redefinition of (2.9), as required.

There are two limits of relevance to our discussion. If the matter fields are heavy, *i.e.* $m \gg (F_A/f)$, the masses of the scalars are given by

$$m_s^2 = m^2 \pm m \frac{F_A}{f} \sqrt{(x_1 + x_2 + y_1 + y_2)}.
 \tag{3.3}$$

An example are fermions charged under PQ that acquire a mass $M \lesssim \mathcal{O}(f)$ when PQ is broken, such as the extra heavy matter in KSVZ axion models [18, 19]. If the heavy fields Φ_{Mi} with couplings to the axion multiplet are also charged under the SM gauge group, they will act as messengers in gauge mediation and give soft masses to sfermions ϕ_{MSSM} at two loops,

$$\tilde{m}_{\phi,GM}^2 = 2c_n \left(\frac{g_n^2}{16\pi^2} \right)^2 \sum_i \left(T^{\Phi_{Mi}} (x_1 + x_2 + y_1 + y_2) \right)^2 \left(\frac{F_A}{f} \right)^2,
 \tag{3.4}$$

where c_n is the quadratic Casimir of ϕ under gauge group G_n . These loop-suppressed contributions from messengers to soft masses are important for MSSM fields which do not have a direct coupling to the axion multiplet (*i.e.* the Kähler contribution (3.5) $\tilde{m}_{\phi,K}^2 = 0$).

Alternatively, if the fields Φ_1, Φ_2 are MSSM fields which satisfy $m \ll (F_A/f)$ then the Kähler couplings, if present (for example as in the DFSZ case), will dominate leading to masses

$$\tilde{m}_{\phi,K}^2 = (y_i^2 - 2z_i) (F_A^2/f^2).
 \tag{3.5}$$

These Kähler -mediated masses are in general one loop larger than the gauge-mediated contributions (3.4). In the case of the third generation quarks the superpotential terms proportional to m^2 may have a small but non-negligible effect. Such terms only appear after electroweak symmetry breaking.

The superpotential couplings, if present, also induce trilinear terms,

$$\mathcal{L} \supset -\lambda_{ijk} (x_1 + x_2 + x_3 + y_1 + y_2 + y_3) \frac{F_A}{f}, \phi_1 \phi_2 \phi_3
 \tag{3.6}$$

which again depend on the combination $x + y$. These are of particular importance in the case of the third generation sfermions due to their large Yukawa couplings.

In explicit calculations, renormalization group evolution from the SUSY breaking scale is important as in general we will take the SUSY breaking scale to be relatively high. There are positive contributions to scalar masses due to non-zero gaugino masses [54] as well as negative contributions from other scalar masses as in split family scenarios [55].

3.2 Gauginos and gaugino screening

Now we turn to gaugino masses, for which there are three significant contributions: tree-level, Kähler -mediation, and gauge-mediation. First, a coupling through the gauge kinetic term C_n leads to a tree level contribution once the axion supermultiplet acquires an F-term,⁶

$$m_{C_n} = \frac{g^2 C_n F_A}{16\pi^2 f}. \tag{3.7}$$

The anomaly coefficients C_n can arise in the UV theory or from anomalous diagrams with chiral matter charged under both PQ symmetry and the gauge group G_n . In the case of the QCD axion the axion, of course, must couple to QCD so $C_3 \neq 0$. Moreover in conventional 4D GUTs there is also the condition $C_1 = C_2 = C_3$.

An additional mass contribution occurs due to the existence of a counter-term in the gauge coupling that appears from any chiral multiplet charged under the gauge group and which develops a non-zero F-term. The presence of such a counter-term was first noted due to its appearance in models of anomaly mediation [56–59], and leads to a mass contribution

$$m_{c.t.} = -\frac{g^2}{16\pi^2} \sum_i 2T_n^{\Phi_i} \frac{F_i}{\langle \Phi_i \rangle} \tag{3.8}$$

where F_i is the F-term developed by a chiral multiplet Φ_i which also has a scalar expectation value $\langle \Phi_i \rangle$. Solving for the F-terms we obtain $F_i = y_i \frac{F_A}{f} \langle \Phi_i \rangle$, and therefore the gaugino mass induced by this term is given by⁷

$$m_{c.t.} = -\frac{g^2}{16\pi^2} \sum_i 2y_i T_n^{\Phi_i} \frac{F_A}{f}. \tag{3.9}$$

The two contributions (3.7) and (3.9) balance such that the gaugino mass is unchanged under the anomalous transformations and demanding this is one of the ways of deriving the expression for the counter-term [56].

As noted in [58], at least part of what is known as anomaly mediation is completely disconnected from any sort of gravitational effect and appears in globally supersymmetric theories. The contribution to the gaugino mass we see here is exactly the result of this portion of anomaly mediation. While such a term is unimportant when considering the classic SUSY spectra with gauginos of similar mass to sfermions, we will later see that it can actually drive the creation of a split SUSY spectrum.

Finally, there is a gauge mediated contribution from any chiral fields which have direct couplings to the axion. Since the mass matrix for the scalars is invariant under field redefinitions, so is the contribution to the gaugino masses. The exact expression is complicated, but we will mostly be interested in two extreme regimes where the result simplifies. The first case is KSVZ-like [18, 19], where heavy states with $m \gg \frac{F_A}{f}$ couple to the axion;

⁶The mass includes a factor $4g^2$ due to our choice of gauge coupling normalisation and a factor of $1/2$ since this is a Majorana mass term.

⁷There is a question of the validity of this term at the origin $\langle \Phi_i \rangle = 0$; the counter-term has been shown to persist in the same form in this case [57].

these act as messengers and give a significant contribution to gaugino masses. The second case is DFSZ-like [20, 21], in which the chiral fields with axion couplings are MSSM fields.

In the KSVZ-like case the masses are given by $m^2 \pm m \frac{F_A}{f} \sqrt{(x_1 + x_2 + y_1 + y_2)}$. Since these fields form a vector-like pair, $T_n^{\Phi_1} = T_n^{\Phi_2} = T_n^\Phi$; they lead to gauge mediated contribution to the gaugino mass given by

$$m_{\text{gauge}} = -\frac{g^2}{16\pi^2} 2T_n^\Phi (x_1 + x_2 + y_1 + y_2) \frac{F_A}{f}. \quad (3.10)$$

Adding this contribution to that obtained from the counter-term due to the small F-term gained by this pair of fields we find a net mass (where we are not including the mass contributions from any light fields to be discussed shortly)

$$m_{1/2} = \frac{g^2}{16\pi^2} \left(C_n - 2T_n^\Phi (x_1 + x_2) \right) \frac{F_A}{f}. \quad (3.11)$$

This is clearly redefinition invariant as expected; the extension to more than one pair of states is straightforward.

An interesting feature of (3.11) is that all dependence on y_i cancels. This is the ‘gaugino screening’ phenomenon [57, 60], where Kähler couplings between messenger fields and the SUSY breaking sector have no effect on gaugino masses at leading order in $\frac{F}{m}$. This result can be understood from a holomorphic perspective. Gaugino masses are given by the θ^2 component of the real gauge coupling R which depends on the holomorphic gauge coupling S as

$$R_n(\mu) = S_n + S_n^\dagger - \sum_l \frac{T_n^{\Phi_l}}{8\pi^2} \log \mathcal{Z}_l, \quad (3.12)$$

at a scale μ . The holomorphic gauge coupling is given by

$$S_n(\mu) = S_n(\Lambda_{UV}) + \frac{b_n}{16\pi^2} \log \left(\frac{\mu}{\Lambda_{UV}} \right) - \sum_i \frac{T_n^{\Phi_i}}{16\pi^2} \log \left(\frac{M_i}{\Lambda_{UV}} \right), \quad (3.13)$$

where M is the physical mass of the messengers which after analytic continuation acts as a spurion, gaining a θ^2 term. Additionally M is the mass of the canonically normalized field and therefore is given by $M = \frac{M_s}{\mathcal{Z}}$ where M_s is the mass that appears in the superpotential. Hence the real gauge coupling is independent of \mathcal{Z} and thus y_i to leading order.

The DFSZ-like case where the states are light is also interesting. The gauge mediated contributions to gaugino masses are completely negligible since for all MSSM fields $m \ll \frac{F_A}{f}$ [61]. Therefore the gaugino masses are given by sum of (3.7) and the terms arising from the counter-term (where the effects of heavy fields are now not included)

$$m_{1/2} = \frac{g^2}{16\pi^2} \left(C_n + 2 \sum_i y_i T_n^{\Phi_i} \right) \frac{F_A}{f}, \quad (3.14)$$

which now does depend on the y_i but is still redefinition invariant as required. From an analytic continuation point of view it is clear what has happened in this case too: the light states are not integrated out and so the real gauge coupling retains dependence on \mathcal{Z}_i through the term $R \supset -\frac{1}{8\pi^2} \sum_i T_n^{\Phi_i} \log \mathcal{Z}_i$.

In a generic model with both heavy and light states the total gaugino mass is

$$m_{1/2} = \frac{g^2}{16\pi^2} \left(C_n - 2 \sum_{m_i \gg \frac{F_A}{f}} x_i T_n^{\Phi_i} + 2 \sum_{m_i \ll \frac{F_A}{f}} y_i T_n^{\Phi_i} \right) \frac{F_A}{f}, \quad (3.15)$$

where the sum is over heavy states for the superpotential couplings and Kähler couplings for light states, and the heavy states are assumed to form vector-like pairs; this expression is also basis independent for the same reason as the 1PI coefficient. Except for the UV-sensitive coefficients C_n , the gaugino masses are proportional to the gauge couplings squared as in gauge mediation, although particular values of the anomalous coefficients can lead to different hierarchies. Compared to the sfermion masses, the gaugino masses can be at the same scale (if there is no large Kähler contribution to the scalars as in (3.4)) or a factor of $\frac{g^2}{16\pi^2}$ lighter for scalars with mass determined as in (3.5). Altering the Kähler contribution can change the difference between gaugino and scalar masses, but does not alter the ratios of gaugino masses themselves, unlike models which interpolate between anomaly mediation and gauge mediation (as for example [62]).

3.3 Higher order contributions

There is one more contribution to gaugino masses which, while not phenomenologically important, is conceptually relevant. The real gauge coupling actually includes a term due to an anomaly when the gauge multiplet is rescaled to canonical normalisation [60]

$$R \supset \frac{g^2}{8\pi^2} T_n^G \log(S + S^\dagger). \quad (3.16)$$

where T_n^G is the Dynkin index of the adjoint representation of the gauge group n . To leading order, by inserting the leading dependence $S \supset \frac{1}{4g^2} - \frac{C_n}{32\pi^2} \frac{A}{f}$, this gives a higher-order gaugino mass contribution

$$m'_{1/2} \supset -\frac{g^4}{8\pi^2} T_n^G \frac{C_n}{16\pi^2} \frac{F_A}{f}. \quad (3.17)$$

This is smaller than the tree level contribution by a factor of $\sim \frac{g^2}{8\pi^2}$ therefore is not usually phenomenologically relevant, but superficially appears not to respect redefinition invariance. The solution is found by including the next to leading order gauge mediated contribution. This is given by [60]

$$m'_{1/2} \supset -\frac{g^2}{8\pi^2} T_n^G (m_{\text{gauge}} + m_{c.t.}). \quad (3.18)$$

Therefore this term is invariant under field redefinitions in the same way as the leading order contribution. For heavy vector-like states there is a cancellation between m_{gauge} and $m_{c.t.}$ which gives a higher-order, but still invariant, contribution to the mass,

$$m'_{1/2} \supset -\frac{g^4}{128\pi^4} T_n^G \left(C_n - 2 \sum_{m_i \gg \frac{F_A}{f}} x_i T_n^{\Phi_i} + 2 \sum_{m_i \ll \frac{F_A}{f}} y_i T_n^{\Phi_i} \right) \frac{F_A}{f}. \quad (3.19)$$

Clearly, this depends on the effective field theory parameters in the same combination as the leading contribution, $m_{1/2}$ given by (3.15), hence can be written as

$$m'_{1/2} \supset -\frac{g^2}{8\pi^2} T_n^G m_{1/2}. \tag{3.20}$$

There are additional corrections to gaugino and scalar masses from gauge mediation at $\mathcal{O}(F^3/M^5)$, where M is the messenger mass. These do not appear in the expression from analytic continuation which cannot capture higher order terms arising from super-covariant derivatives. An expression calculated directly from loop diagrams can be found in [61], however for f in the range of interest these terms are normally negligible. In a region where these corrections become important they typically increase the scalar masses and decrease gaugino masses. One possibly interesting scenario is the case where the leading contribution (3.15) vanishes identically, so the correction (3.20) also vanishes; then the dominant source of mass is the gauge mediated term of order (F^3/M^5) , resulting in highly suppressed gaugino masses and a very split spectrum.

4 Notable models

In the following subsections, we begin by exploring scenarios that lead to unusual non-minimal SUSY spectra with possible relevance for the LHC. Axion mediation naturally interpolates between a gauge-mediated type spectrum with gauginos and scalars at the same scale (when the mediation is dominated by heavy fields charged under PQ and SM fields) and a Kähler mediated split spectrum (when the dominant axion couplings are to Standard Model fields). Specifically, we consider the phenomenologically interesting case of split SUSY with a sfermions a loop factor above the gauginos in mass (the so-called mini-split scenario). In addition we look at ‘natural SUSY’ models with a mix of gauge and Kähler mediated contributions — split families or split gauginos, which ameliorate the tension between the current LHC constraints and the requirements of naturalness. We also discuss some of the special features of SUSY-breaking mediation by the QCD axion supermultiplet itself.

4.1 Split supersymmetry

4.1.1 Scalars and gauginos

With the continued null supersymmetry search results and the discovery of a relatively heavy and SM-like Higgs, a possible interpretation of current data is that the Higgs mass may be tuned, even in the presence of supersymmetry. Indeed, in the MSSM at least 10% tuning is inevitable given the high Higgs mass and stop and gluino limits from the LHC [63]. In Split Supersymmetry, the tuning of the Higgs mass is one of the central ingredients; the rest of the spectrum is minimal with scalar superpartners of the SM fermions parametrically heavy, while gauginos and possibly higgsinos near the TeV scale can preserve the success of gauge coupling unification and provide a dark matter candidate [64–66]. In axion mediation it is natural for the gauginos to be lighter than the scalars by a loop factor or more, providing an attractive mini-split spectrum, as we will now discuss.

In axion mediation, split supersymmetry is very straightforward to achieve. In the presence of Kähler couplings, y , between the axion and MSSM chiral multiplets, sfermions gain masses $m_s^2 \sim \frac{y^2 F_A^2}{2 f^2}$ while the gaugino masses are suppressed by an extra loop factor $\frac{g_n^2}{16\pi^2} \sum_i T_n^{\Phi_i} \sim 1/100$. This contribution to gaugino masses appears as part of anomaly mediation, but is independent of any supergravity effects, relying only upon the presence of light fields charged under the SM gauge groups — in this case, the MSSM fermions themselves (this mechanism is dubbed Kähler mediation in [56]).

The sfermions in this case are relatively light — 10-1000 TeV — as in the case of mini-split SUSY and in agreement with the measured Higgs mass of 125 GeV without the need for further model building to raise the Higgs mass [63, 67]. In this mini-split case it is either essential or advantageous that there is flavor structure which prevents too-large flavor- and CP-violating observables arising from these not very heavy sfermions of the first and second generations. Importantly, in axion-mediation, this issue has a natural solution as the axion can arise from Kähler moduli which couple universally to scalars and for f_a of string or GUT scale of 10^{16} GeV or lower, these universal contributions dominate over M_{pl} -suppressed contributions. This feature alleviates the flavor concerns that are present in generic anomaly mediation scenarios [63, 64, 68]. Note that in many string constructions of supersymmetry breaking, e.g., KKLT [49], the non-universal complex structure moduli are fixed by SUSY-preserving dynamics at high scale and, as explained in the introduction, are thus not usefully described as axion multiplets. There is, however, a final light Kähler modulus which is only stabilized by SUSY-breaking dynamics and which can play the role of the axion supermultiplet of axion mediation.

Phenomenologically there are a range of possibilities depending on the axion decay constant f and the couplings in the Kähler potential. First, consider an axion in the dark matter window [69–71], $f \sim 10^{11}$ GeV. Then to achieve TeV-scale gauginos with order one couplings y , we take the scale of SUSY breaking to be $\sqrt{F} \sim 10^8$ GeV, generating a typical spectrum as shown in figure 1(a). The scalars are at 40 TeV, leading to a Higgs mass of 125 GeV given $\tan\beta \sim 4$ [72, 73]. All the gauginos could potentially be discovered at the LHC. The relatively light scalars mean that the gluino, if produced, would decay promptly inside the detector and standard searches for the gluino apply [74]. Of course, increasing F or the Kähler couplings increases the overall scale of the spectrum, taking the gauginos out of experimental reach.

The gravitino in this case has a mass of ~ 2 MeV; the gauginos are stable on collider scales but will decay to the gravitino with a lifetime of ~ 0.3 s, i.e. sufficiently quickly not to pose a threat to BBN [75–78]. The gravitino will also likely be overproduced unless the reheat temperature is very low, and depending on the exact value of f the gravitino and the axion will be co-dark matter.

The spectrum is changed for a stringy axion, $f \sim 10^{16}$ GeV. For TeV-scale gauginos and order one couplings y , a SUSY breaking scale $\sqrt{F} \sim 10^{10}$ GeV gives a spectrum of the form shown in figure 1(b). The scalars are at ~ 30 TeV, with $\tan\beta = 5$. The gravitino in this case has a mass of ~ 100 GeV, comparable in mass to the gauginos, so either the gravitino or the gauginos could be the LSP depending on $\sqrt{F_A}$ and Kähler couplings y and z .

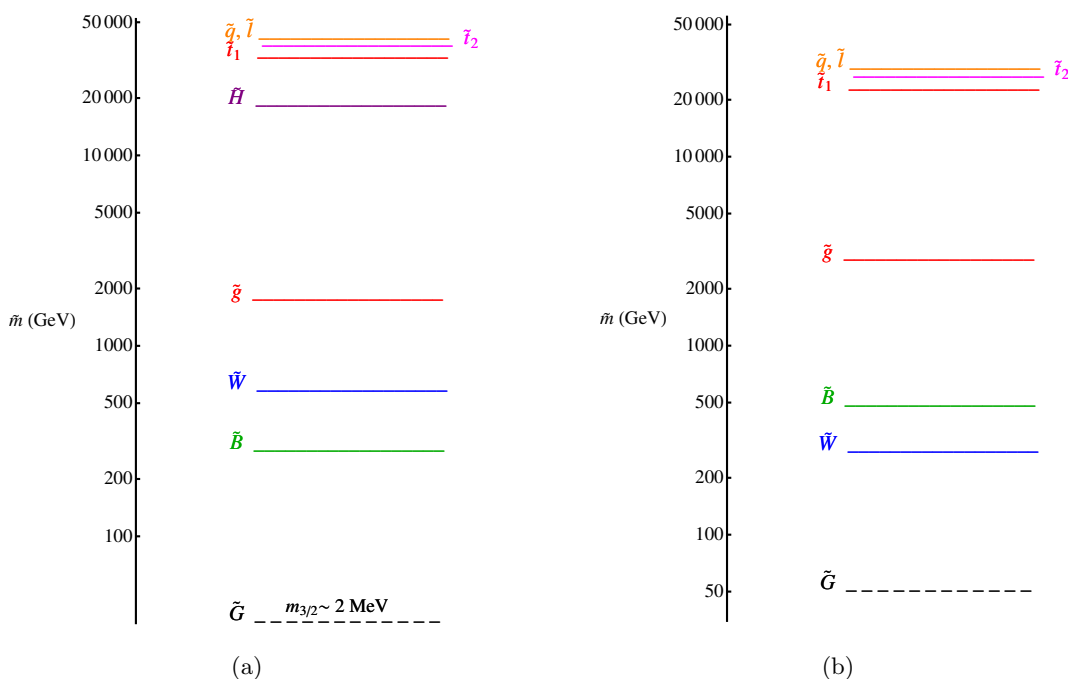


Figure 1. The spectra of MSSM soft masses, after running to the electroweak scale, in the case of Kähler mediation for (a) $f = 10^{11}$ GeV in the ‘axion window’ and (b) a ‘stringy’ axion with $f = 10^{16}$ GeV. In the former case, μ is generated through PQ symmetry breaking, the gravitino is light, and the bino is the NLSP (assuming universal gaugino masses at the mediation scale). The sfermions are at about 40 TeV, leading to a Higgs mass of 125 GeV for $\tan \beta \sim 3$. In the latter stringy case, with similarly heavy scalars giving the observed Higgs mass, the gravitino has a mass of ~ 100 GeV, comparable to the gauginos, so either the gravitino or the gauginos could be the LSP. Non-universal gaugino masses, which are a natural possibility in axion mediation, can lead to a wino NSLP avoiding potential conflict with late decay constraints in cosmology.

If a neutralino is the NLSP in this scenario, it will decay to the gravitino with a lifetime of 10^6 - 10^7 s, in conflict with element abundances in BBN as well as CMB measurement [79]. A bino NLSP or LSP is thermally overproduced, leaving little room out of these limits. However, for a stringy axion it is natural to have non-universal gaugino masses as we will discuss in section 4.3, and alternate neutralino NLSP admixtures can alleviate the tension; for example, a 100 GeV wino has thermal abundance of less than 5% of dark matter density, easing cosmological constraints.

In the case that the gravitino is heavier than one or more of the gauginos the gravitino lifetime is on the order of 10^7 - 10^8 s, also leading to tension with cosmology. Limits from BBN result in a bound on reheating temperature of less than 10^6 GeV unless the gravitino is heavier than a TeV [77]. The latter is disfavored for a spectrum with TeV-scale gauginos in axion mediation as the axion scale would be $f > 10^{17}$ GeV.

Finally, for string values of $f \sim 10^{16}$ GeV, the initial axion misalignment angle has to be small not to over-produce dark matter, with this environmental tuning becoming more severe as f increases. Since there is no reason for the angle to be tuned further than necessary, we again expect axion and neutralino or gravitino co-dark matter.

4.1.2 EWSB and Higgsinos

Even though the Higgs mass is tuned, we must still consider the details of the Higgs sector to ensure that electroweak symmetry breaking takes place. There is a connection of PQ symmetry breaking to the MSSM Higgs sector, which has a U(1) symmetry in the absence of the μ term. If the MSSM Higgses are charged under the PQ symmetry responsible for the axion, then breaking PQ will generate a μ term for the Higgs, $\mu \sim \lambda f^2/M_{\text{pl}}$, where λ is the coupling between the Higgses and the field which spontaneously breaks PQ. Otherwise, if the Higgses are charged under a different U(1) than the PQ symmetry of the axion, the μ term has to be generated independently of the axion sector.

In the ‘axion window’, $f \sim 10^{11}$ GeV, it is possible to identify the PQ symmetry of the axion with that of the Higgses. For concreteness, consider an example UV model [52, 80]

$$W = \lambda \frac{X^2}{M_{\text{pl}}} H_u H_d + \lambda_s S(XY - f^2). \quad (4.1)$$

Here X, Y are charged under PQ with charges of -1 and $+1$, respectively; the two fields acquire vacuum expectation values of order f . H_u, H_d have PQ charges of $+1$ and other MSSM particles have PQ charges set by the Yukawa interactions, e.g. $\{Q, L, \bar{u}, \bar{d}, \bar{e}\} = \{-1, -1, 0, 0, 0\}$. This generates a μ term with $\mu \sim \lambda \frac{f^2}{M_{\text{pl}}}$ when PQ is spontaneously broken. Then

$$\mu = \lambda (10^4 \text{ GeV}) \left(\frac{f}{10^{11} \text{ GeV}} \right)^2. \quad (4.2)$$

For a dark matter window axion the μ term is at the same scale as the sfermions. A large μ term also has an important effect on the gaugino spectrum: threshold corrections proportional to μ change the bino and wino masses by as much as 20% [81],

$$\delta M_1 = \frac{3}{5} \frac{\alpha_1}{4\pi} \mu \frac{2m_A^2 \sin 2\beta}{m_A^2 - |\mu|^2} \log \left(\frac{m_A^2}{\mu^2} \right); \quad \delta M_2 = \frac{\alpha_2}{4\pi} \mu \frac{2m_A^2 \sin 2\beta}{m_A^2 - |\mu|^2} \log \left(\frac{m_A^2}{\mu^2} \right). \quad (4.3)$$

We take this effect into account in the spectrum in figure 1.

The low energy effective theory below the scale of PQ breaking can be written as couplings in the Kähler and superpotential as in section 2. Keeping the same Higgs charges, we also generate B_μ through the term

$$W \supset \mu e^{-(x_{H_u} + x_{H_d})A/f} H_u H_d + \text{h.c.} \quad (4.4)$$

so

$$B_\mu = \frac{\mu}{f} F_A (x_{H_u} + x_{H_d}) = 2\mu \frac{F_A}{f}. \quad (4.5)$$

Higgs fields coupled to the axion in the Kähler potential will give masses to the scalar components of the same order as the sfermion masses, $m_{H_u}^2 \sim m_{\tilde{H}_u}^2 = y^2 (F_A/f)^2$. Then we have $\mu \sim 10$ TeV, $m_{\tilde{H}_u}^2 \sim m_{H_u}^2 \sim 40$ TeV, which results in successful electroweak symmetry breaking with $\tan \beta \sim \text{few}$. This mechanism also provides an upper bound on sparticle masses: while all the masses in figure 1(a) can be scaled up by increasing

f/F_A and take gauginos out of observable reach, the Higgsinos are constrained to be below 100 TeV by gauge coupling unification [63], so the connection with EWSB ensures a relatively light spectrum.

Of course for the string axion, this mechanism would make μ far too large unless the coupling λ is quite small, $\lambda \sim 10^{-5}$. Another possibility is to introduce heavy messengers which couple to the Higgs to generate a μ term on the order of gaugino masses. One example is coupling the Higgses directly to the messengers (heavy fermions charged under PQ and SM gauge fields) as in models of lopsided gauge mediation [82, 83].

4.2 Split families and ‘Natural SUSY’

With current limits on universal squark masses reaching close to 2 TeV, [84, 85] an attractive scenario that maintains a natural solution to the hierarchy problem is the split family class of models — also known as Natural Supersymmetry — in which the stops and sbottoms are significantly lighter than the 1st and 2nd generation squarks [1, 2, 86]. Split family models take advantage of the fact that light flavor generations contribute more to the stringent limits due to high production rates, while 3rd generation squarks contribute directly to the tuning in the Higgs mass. Thus, separating the two issues by raising the mass of the 1st and 2nd generations can relieve the strain on naturalness from experimental limits (for some example models, see [87–90]).

The most viable way of achieving a split family spectrum in axion mediation is from an axion with couplings both to MSSM and new heavy chiral multiplets. The first two generations are spatially separated from the third along an extra dimension and gain a large mass from Kähler couplings, while the third generation does not. Gauginos gain mass both from the heavy multiplets and the Kähler couplings of the first two generations. The third generation gains small masses through gauge mediation from the heavy multiplets. A common issue with split family scenarios is that heavy first and second generation sfermions in combination with light gauginos tend to run third generation scalars negative. The maximum splitting between generations is limited to be a factor of 5-10, so flavor problems arising from the first two generation sfermions are not sufficiently solved [55]. It is possible to address the flavor problem in a relatively elegant fashion in which the first two generations have an approximate symmetry protecting them against too-strong flavor violations. In particular, in the field theory case Kähler couplings arise from integrating out additional heavy fields which couple to the MSSM via new gauge interactions. This makes the coefficients discrete parameters based on charge assignments that can naturally be the same for some or all generations. Similar scenarios can also occur in the string case. The viability of a low energy natural spectrum is studied for example in [91, 92], while a model where the first two generation sfermions gain an additional mass contribution through universal couplings appears in [93]. However, to calculate the tuning of such a spectrum in axion mediation requires detailed knowledge of the priors on axion couplings expected in a string UV completion. This is a difficult task which do not attempt to address here. We content ourselves with the overall features of the resulting low-energy superpartner spectrum, a representative example of which is shown in figure 2; we use `SOFTSUSY` to compute the low-energy masses [94].

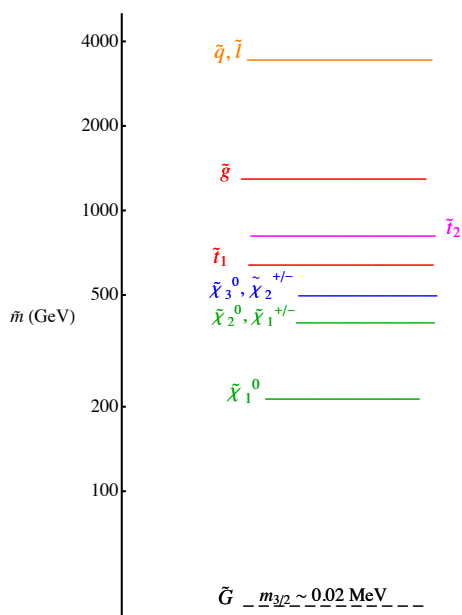


Figure 2. The low energy soft terms in the case of a split family spectrum, with $f_a \sim 10^{11}$ GeV. TeV-scale visible sparticles require a SUSY breaking scale $\sqrt{F} \sim 10^7$ GeV. The first two generations have significant Kähler couplings to the axion multiplet, while those of the third generation are suppressed. Combined with a gauge mediated contribution this leads to mass splitting of the sfermion masses of about a factor of 5, taking the light generations out of experimental reach. Universal gaugino masses lead to a bino NLSP and gravitino LSP, while in the non-universal gaugino case different NLSPs are possible. Mildly heavy higgsinos lead to some tension with naturalness.

There are some drawbacks for natural SUSY in axion mediation due to the high scale of SUSY breaking. First, in light of LHC limits, a successful split families model has light stops and a gluino above experimental limits ($m_{\tilde{g}} \geq 1.25$ TeV assuming a split family spectrum, e.g. [95–97]); so in raising the gluino mass we do not want to raise the stop masses. Unfortunately the stop mass grows significantly through RG running from a high scale. With even a somewhat low axion scale $f = 10^{11}$ GeV, two-loop contributions from the gluino increase the Higgs mass tuning by pulling up the stop mass such that $m_{\tilde{t}} \sim (2/3)m_{\tilde{g}}$, regardless of the stop soft masses at the SUSY breaking scale. This is just the general fact that theories with a low cutoff scale are better for naturalness given experimental constraints (for further discussion see e.g. [91, 92]). In addition, the constraints on models with a light gravitino and bino NLSP include charginos heavier than 450 GeV and stops above 580 GeV [98].⁸ Another possible concern for the axion decay constant in the string motivated window $f \sim 10^{16}$ GeV is that additional generic Planck-suppressed gravity mediated contributions to soft masses may be large enough to cause tension with flavor constraints, however understanding the extent to which this may be a problem requires a full string construction. On the other hand, A-terms can naturally be large in this case which can increase the Higgs mass within the MSSM.

⁸We thank T. Gherghetta for calling our attention to these limits.

4.3 Split gaugino masses

Typically, gauge or gravity mediation lead to gaugino masses falling into a pattern proportional to their coupling constants $M_i \propto g_i^2$. This leads to the well-studied case of fairly heavy gluinos and (assuming R-parity conservation) substantial missing energy signals from decays to the lightest neutralino of moderate mass. In contrast in axion mediation these patterns can be easily relaxed.

If the UV theory has an underlying GUT group and the PQ charges of fields and axion multiplet couplings respect this group then the standard pattern of gaugino masses is obtained. However, many models deviate strongly from this very specific case. If, for example, the UV completion is an ‘orbifold-GUT’ theory or a IIB type string model based on brane stacks there is often no true 4D GUT symmetry, and matter which would normally be expected to fall into a single irreducible $SU(5)$ multiplet is now localized on different branes [99, 100]. Therefore these fields can very naturally have differing charges under the PQ symmetry despite the fact that the success of supersymmetric gauge-coupling unification is still naturally preserved and explained, at least in the orbifold-GUT case. The anomaly coefficients may also naturally be non-universal in an underlying heterotic string model [101], or in other string constructions.

Non-universal gauginos can yield interesting phenomenology. Two cases to consider are $C_1 = 0$ and $C_3 = 0$. In the first case $C_1 = 0$, the hypercharge $U(1)$ is non-anomalous; then it is possible to obtain a very light Bino, $M_1 \ll M_2$ instead of the usual relation $M_1 \sim \frac{1}{2}M_2$. This loosens the indirect limits on bino mass: without the correlation to chargino limits on M_2 , M_1 can be 50 GeV or less. For high enough SUSY breaking scales (such that the gravitino is heavier than the bino), the bino can then be a viable dark matter candidate with the correct relic abundance, a case that has been studied generally (see for example [102–105]).

The second case of an anomaly-free $SU(3)$ cannot be a QCD axion but can be motivated for example in the axiverse picture of multiple axions. In axion mediation, this results in a gluino much lighter than the wino and bino at the SUSY breaking scale, provided $C_3 = 0$ and $x_i, y_i \ll C_{1,2}$. The gluino contributes significantly to the fine tuning of the Higgs mass at two loops, so a light gluino can reduce the fine tuning of the weak scale and possibly reducing missing energy signatures in decays through a squeezed spectrum [63, 93]. Further, certain mass ratios between M_1, M_2 , and M_3 have been studied and found to reduce the fine tuning of the electroweak scale provided there is no fine-tuning in the UV theory to set up these parameters [106, 107]. Not surprisingly, these mass ratios involve a gluino much lighter than $SU(2)$ gauginos at the SUSY breaking scale.

4.4 QCD axion mediation

In a minimal model, the axion multiplet provides the solution to the strong-CP problem as well as mediating supersymmetry breaking that is, the pseudo-scalar component of the multiplet which acquires the (dominant) F-term is the QCD axion. As mentioned in the Introduction, for the axion to solve the strong-CP problem its mass must arise almost entirely from non-perturbative QCD breaking of the PQ symmetry, and thus, given the constraints

on the axion scale f , the QCD axion mass must be small.⁹ Such a light axion with couplings to the SM is actively searched for, with the possibility of its direct detection in laboratory experiments as well as indirect detection via astrophysical observations [108, 109]. As we will now discuss, in addition to many aspects of the phenomenology of previous subsections, this presents the exciting possibility of correlating axion detection measurements to supersymmetric spectra; although of course making such measurements experimentally, with sufficient precision to show correlation, would be a very challenging task.

In more detail, one experimental observable is the anomalous axion-photon coupling,

$$\frac{\alpha_1 C_1}{8\pi f} a F_{\mu\nu} \tilde{F}^{\mu\nu}. \tag{4.6}$$

Searches for axions that have a two-photon vertex are especially promising, including experiments which look for axions from the Sun [110] and the galaxy and place bounds of $C_1/f < 8.8 \times 10^{-11} \text{ GeV}^{-1}$ ($f_a > 3 \times 10^8 \text{ GeV}$ for a QCD axion). Laboratory searches include ‘shining light through a wall’ [111] and microwave cavity experiments [112–114] which take advantage of the axion to photon conversion in the presence of a magnetic field and can place limits or potentially discover a light axion.

Other detectable interactions are the derivative couplings of the axion to fermions of the form $y_i(\partial_\mu a)\bar{\psi}\sigma^\mu\psi$, and $x_i am_\psi\bar{\psi}\psi$ which arise in the SUSY context from Kähler and superpotential couplings. From these couplings one derives the basis-independent 1PI interaction between the axion and fermions,

$$-(x_1 + x_2 + y_1 + y_2) \frac{m_a}{f} a \bar{\psi}\psi. \tag{4.7}$$

Interactions of the form (4.7) with electrons and quarks are experimentally relevant. The former can lead to excessive white dwarf cooling which rules out $f < 10^9 \text{ GeV}$, though there may be a possibility that axion emission can improve fits of white dwarf cooling models to the data [108]. If axions couple to quarks, cooling by emission from nuclei can also be bounded by constraints from SN 1987A to be $f > 4 \times 10^8 \text{ GeV}$.

From searches for superpartners at the LHC we can in turn learn about the scale of supersymmetry breaking and the mass spectrum by measuring the gravitino mass (in the case it is the LSP) and gaugino and sfermion masses and correlate the mass spectrum with axion couplings. An extra handle on the SUSY mediation is particularly appealing in split supersymmetry, which has a short list of observables due to very heavy scalars. Gauginos masses will depend directly on the anomalous coupling and the derivative coupling. Then we can measure the anomalous photon coupling C_1 and axion fermion couplings y (and continue to not observe flavor violations). Along with either the gravitino mass or some knowledge from the Higgs sector this would give enough information to prove mediation via this mechanism.

⁹For completeness we note that there exist exceptions to this statement that, while unlikely in our view, cannot be definitively be excluded. If, for example, the matter content of the UV theory were such that QCD becomes non-asymptotically free and strongly coupled *in the UV* then extra contributions to the axion mass could arise preserving the natural solution to the strong-CP problem.

In the case of multiple axions the multiplet with the strongest couplings to MSSM fields will be the one to dominate mediation when all F-terms are comparable. So, it may be possible more generally to discover the multiplet which dominates mediation in axion detection experiments and provide evidence for axion mediation of supersymmetry breaking.

5 Axino as goldstino and cosmological constraints

In discussing the phenomenology of axion mediation in section 4, we focus on the minimal case in which a single axion is the leading source of supersymmetry breaking; that is, the axion multiplet is the only one in the theory that gains a significant F-term. Then, the axino, as the fermionic component, is also the goldstino in the theory; it is eaten by the gravitino through the ‘super-Higgs’ mechanism once we move to supergravity. The fact that the axino is not an extra degree of freedom but part of the gravitino multiplet is an additional benefit of SUSY breaking in the axion sector, alleviating some tensions with cosmology as we discuss below.

5.1 The axino and the goldstino

As is well known, couplings of the goldstino \tilde{G} are fixed by supercurrent conservation,

$$\mathcal{L} = i\tilde{G}^\dagger \bar{\sigma}^\mu \partial_\mu \tilde{G} - \frac{1}{F} \left(\tilde{G} \partial_\mu j^\mu + c.c. \right) \quad (5.1)$$

where j^μ is the supercurrent of the other chiral and gauge superfields [54],

$$j^\mu = \sigma^\eta \bar{\sigma}^\mu \psi_i \partial_\eta \phi^{*i} - \frac{1}{2\sqrt{2}} \sigma^\eta \bar{\sigma}^\rho \sigma^\mu \lambda^\dagger F_{\eta\rho} + \dots \quad (5.2)$$

The goldstino couplings determine the dominant interactions of the longitudinal gravitino components: these are physical and are important in collider phenomenology (if the gravitino is the LSP) as well as in cosmology. However, these couplings do not readily appear to be invariant under chiral rotations discussed in section 2, as they are fixed by the form of the supercurrent and do not depend on whether the goldstino multiplet transforms non-linearly under a PQ symmetry. However, the on shell couplings depend on the superfields’ masses, making the connection between axino and goldstino interactions manifest.

This is clearest in the case of the axino-gaugino-gauge coupling; from the goldstino Lagrangian,

$$\mathcal{L} \supset -\frac{i}{F} \frac{m_{1/2}^n}{2\sqrt{2}} \tilde{G} \sigma^\eta \bar{\sigma}^\rho F_{\eta\rho}^n \lambda_n^\dagger. \quad (5.3)$$

By supersymmetry, the axino-gaugino-gauge is of the same form as the gaugino masses (3.15) as they both result from the coupling $\int d^2\theta \frac{\alpha_a}{16\pi} \frac{C_a}{f} A W_\alpha W^\alpha$ along with loop contributions from chiral fields. Then the goldstino coupling is manifestly invariant,

$$\mathcal{L} \supset -i \frac{\alpha_n}{8\sqrt{2}\pi f} \left(C_n - 2 \sum_h T_n^{\Phi_h} (x_{h1} + x_{h2}) + 2 \sum_l y_l T_n^{\Phi_l} \right) \tilde{G} \sigma^\eta \bar{\sigma}^\rho F_{\eta\rho}^n \lambda_n^\dagger \quad (5.4)$$

and since the axion has the only F-term, it is also the same in both descriptions.

For matter multiplets, couplings are between the scalar mass eigenstates and the corresponding fermion mass eigenstates and proportional to the mass difference between the two. To simplify the calculation we consider one pair of vector-like fields. The fermions have degenerate masses m and the combination that couples to the scalar mass eigenstate is the same combination of chiral multiplets. Since the masses of the matter multiplets are again invariant under chiral rotations, we find the coupling proportional to the masses,

$$\mathcal{L} \supset \frac{m}{f} (x_1 + x_2 + y_1 + y_2) (\psi_2 \phi_1 + \psi_1 \phi_2) + \frac{F}{f^2} ((\psi_1 \phi_1 (y_1^2 - 2v_1) + \psi_2 \phi_2 (y_2^2 - 2v_2))), \tag{5.5}$$

which is of course redefinition invariant. There are also non-renormalizable couplings between two or more axinos and MSSM fields; however, these are not constrained by the supercurrent coupling and also are not phenomenologically relevant.

5.2 Cosmology

A frequent issue with an axion in the context of supersymmetry is the cosmological axino problem; for a large class of models, the axino is not protected by any symmetry, and upon SUSY breaking acquires a mass of at least the gravitino mass [115, 116]. If the axion multiplet does not directly participate in SUSY breaking, it will generically acquire a mass from supergravity effects in the Kähler potential,

$$K \supset \int d^4\theta \frac{(A + A^\dagger)^2 (X + X^\dagger)}{M_{Pl}} \sim \frac{1}{2} m_{3/2} \psi_a \psi_a \tag{5.6}$$

The mass is then at least of the order of the gravitino mass, unless there are fortuitous cancellations or sequestering effects. Two weakly interacting particles with comparable masses, the axino and the gravitino, can be disastrous for cosmology: as pointed out in [116], as at least one of them will be overproduced over a large range of gravitino masses. In particular, the energy density has opposite scaling with the mass for the axino and the gravitino.

In more detail, the energy density $m_{3/2} Y_{3/2}$ of the gravitinos (where $Y_{3/2}$ is the gravitino abundance) is given by¹⁰

$$m_{3/2} Y_{3/2} = m_{3/2} K M_{Pl} T_R \langle \sigma_{3/2} v \rangle, \tag{5.7}$$

where $\langle \sigma_{3/2} v \rangle$ is the velocity-averaged scattering cross-section and $K = \frac{n_s(T) \sqrt{90} \zeta(3)}{n_s(T_R) \pi^3 \sqrt{N_*}}$ is a numerical factor that depends on the change in entropy density since reheating and the number of relativistic degrees of freedom [117]. The coupling strength of the gravitino scales inversely with its mass, and the cross-section inversely with the mass squared: $\sigma_{3/2} \sim \frac{\tilde{m}^2}{m_{3/2}^2 M_{Pl}^2}$. So the energy density is inversely proportional to the gravitino mass,

$$m_{3/2} Y_{3/2} \propto \frac{T_R \tilde{m}^2}{m_{3/2} M_{Pl}}. \tag{5.8}$$

¹⁰Here we assume that scattering dominates gravitino production for $m_{3/2} \gtrsim 10^{-4}$ GeV, the medium- to high-scale range of supersymmetry breaking relevant for axion mediation.

The abundance is similar for the axino,

$$m_{\psi_a} Y_{\psi_a} = m_{\psi_a} K M_{Pl} T_R \langle \sigma_{\psi_a} v \rangle \tag{5.9}$$

However, for the axino the cross-section is independent of its mass and depends only on f , $\sigma_{\psi_a} \sim 1/f^2$. For an axino with the same mass as the gravitino, the energy density instead grows with its mass,

$$m_{\psi_a} Y_{\psi_a} \propto \frac{m_{3/2} T_R M_{Pl}}{f^2}, \tag{5.10}$$

and the axino is overproduced at high $m_{3/2}$, while the gravitino is overproduced at low $m_{3/2}$.

This results in an upper bound on low reheating temperature of $T_R < 3 \times 10^5$ GeV for all values of $m_{3/2}$ and an axion in the axion window, 10^9 GeV $< f < 10^{12}$ GeV [116]. The bound is relaxed for higher values of f in the string axion regime, $f \sim 10^{16}$ GeV: there, $T_R \lesssim 10^9$ GeV which is less restrictive but can still run into tensions for instance with theories of high-scale baryogenesis.

The axino being eaten by the gravitino is very beneficial for cosmology: as reviewed here, a light axino and gravitino together result in low reheating temperature for much of the parameter space causing tension between cosmology and supersymmetry breaking. Of course, if the axino is the goldstino, the limit from a separate axino abundance disappears, and high reheat temperatures are allowed for a range of gravitino masses. Thus, an axion participating directly in supersymmetry breaking dynamics and acquiring an F-term is a clear and natural way to relax the friction between cosmology and a supersymmetric axion.

6 Discussion and conclusions

In this work we have considered the possibility that SUSY breaking is mediated primarily through the interactions of a generalized axion multiplet. Non-minimal spectra, such as those of both natural SUSY and split SUSY are simply realized.

We argued in section 1.1 that from a UV perspective axion-mediation of supersymmetry breaking is a very natural and attractive possibility. This is especially (but not only) true in the axiverse context, where we expect many light axion-like fields. Following a thorough review in section 2 of axion supermultiplet couplings in an effective theory language, we then show in section 3 that the SUSY breaking contributions to gaugino versus scalar masses depend on different field-redefinition-invariant combinations of the axion supermultiplet couplings. This leads to a straightforward implementation of split supersymmetry, with heavy scalars separated from a loop factor from and light gauginos, so realizing the attractive mini-split scenario, as discussed in detail in section 4.1.

A second interesting feature discussed in section 4.1 is that the Kähler couplings which tend to dominate the visible sector sfermion masses can be naturally universal, explaining the non-observation of flavor violation due to supersymmetric partners. In the string axion case this occurs when the axion is part of a modulus which couples universally, as is the case in many explicit constructions of SUSY-breaking in string theory, such as KKLT-like scenarios with SUSY-breaking dynamics stabilizing the overall Kähler modulus. Moreover it is well known that to leading order Kähler moduli know nothing about flavor

structure and therefore couple universally to all generations provided they are geometrically localized in the same place [118]. Since the axion must be part of a multiplet that is stabilized by non-SUSY-preserving dynamics and, in addition, has a scale parametrically smaller than M_{pl} (either $f \sim 10^{16}$ GeV in the string case, or $\sim 10^9 - 10^{12}$ GeV in the ‘axion window’) this implies that the universal axion couplings naturally dominate the non-universal gravitational and heavy moduli contributions to sfermion masses. In the field theory case, Kähler couplings arise from integrating out additional heavy fields which couple to the MSSM via new gauge interactions. This makes the coefficients discrete parameters based on charge assignments that can naturally be the same for some or all generations. On the other hand a natural SUSY spectrum can be realized by localizing the third generation separately from the first two, or by taking the discrete charge assignments to differ across generations; some features of this scenario were discussed in section 4.2.

In addition, if the axion couples to visible-sector fields in the superpotential, as one would expect in the general case, then there will be a characteristic pattern of (small) splittings among soft terms. The sfermion mass squared will be split around a common central value (either given by the Kähler couplings, or the gauge-mediated contributions that occur in the KSVZ-like-axion case) by $2m_i \frac{F}{f}$ where m_i is the mass of the corresponding fermion. Therefore lighter generations will be nearly degenerate and the stops will be spread further. Observation of such a pattern would give supporting evidence in favor of axion-mediated SUSY breaking.

Mediation of SUSY breaking via the QCD axion multiplet itself with either axion scale in the axion window, or at the string value $f \sim 10^{16}$ GeV, is experimentally interesting as well. The axion itself can be detected through its derivative couplings to matter or the anomalous photon coupling [108] and these couplings can be correlated with the spectrum of superpartners to confirm some features of axion mediation.

Axion mediation is also cosmologically beneficial in many cases as the axino is eaten by the super-Higgs mechanism to give rise to the massive gravitino. This was studied in section 5 where we argue that cosmological constraints arising from axino and gravitino overproduction are ameliorated. There is an alternative case of interest where, even though the axion multiplet dominates mediation to the visible sector, there is another field with a larger F-term. This may happen if the fields with the largest F-term do not have strong couplings to the visible sector; then the axion is not the goldstino and will not remain light. Even in this case there are cosmological benefits: since the axion sector is now part of the supersymmetry breaking sector the axino is expected to gain a mass of order $\sim \sqrt{F}$, sufficiently large to be cosmologically safe. This is similar to the cosmology studied in [119].

Finally, there are a number of possible extensions of this work. In this paper we focus on the case of a single axion supermultiplet which dominates the mediation of supersymmetry-breaking to the visible sector. However it is also possible that multiple axion multiplets are involved in the mediation, and, moreover, motivated by the same considerations that support the axiverse there can be multiple SUSY breaking sectors, each with its own goldstino [120] or set of goldstinos [121]. Related to this is the natural possibility that the axion under consideration is the R-axion, a case that certainly deserves study. We hope to return to these topics in a future work.

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