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On the diversity of gauge mediation: footprints of dynamical SUSY breaking

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ABSTRACT: Recent progress in realising dynamical supersymmetry breaking allows the construction of simple and calculable models of gauge mediation. We discuss the phenomenology of the particularly minimal case in which the mediation is direct, and show that there are generic new and striking predictions. These include new particles with masses comparable to those of the Standard Model superpartners, associated with the pseudo-Goldstone modes of the dynamical SUSY breaking sector. Consequently there is an unavoidable departure from the MSSM. In addition the gaugino masses are typically significantly lighter than the sfermions, and their mass ratios can be different from the pattern dictated by the gauge couplings in standard (i.e. explicit) gauge mediation. We investigate these features in two distinct realisations of the dynamical SUSY breaking sector.

Contents

1 Introduction

In the run-up to the LHC, the implementation of supersymmetry (SUSY) breaking and its mediation are coming under renewed scrutiny. Attention has recently focussed on how dynamical supersymmetry breaking (DSB) can be achieved, and how its effects can subsequently be transmitted to the Standard Model sector. This interest was stimulated by the observation of Intriligator, Seiberg and Shih (ISS) [1] that DSB readily occurs in very simple and calculable SQCD-like models.

Clearly it is the interaction of the DSB sector with the visible sector that plays a crucial role in BSM phenomenology. However, constructing a viable model that incorporates both sectors presents a twofold problem: both SUSY breaking *and* R-symmetry breaking need to be transmitted to the visible sector. The R-symmetry plays an important role because supersymmetry breaking requires unbroken R-symmetry (in a generic theory) [2], which is at odds with the fact that (Majorana) gauginos must have a mass that violates R-symmetry.

In principle the metastable models of ISS can circumvent this theorem by allowing a moderate and controlled amount of R-symmetry breaking. What ISS reminded us is that,

because the Nelson-Seiberg theorem [2] applies only to the global vacuum of the theory, we are at liberty to generate gaugino masses if we are prepared to tolerate a certain amount of metastability.

New avenues for gauge mediation were consequently opened up. One phenomenological application came shortly after with ref. [3], which noted that because the ISS model breaks supersymmetry in a magnetic Seiberg-dual formulation, the couplings of *explicit* messenger fields to the DSB sector is naturally suppressed by powers of Λ_{ISS}/M_{Pl} where Λ_{ISS} is of order the Landau pole in the theory.[1] Thus the magnetic theory can maintain an approximate R-symmetry even if the underlying electric theory has no R-symmetry and is generic. The phenomenology of this scenario is similar to standard gauge mediation although, because of the weakness of the coupling to the DSB sector, the scale of supersymmetry breaking has to be much higher than is normally assumed. An alternative method of dealing with the R-symmetry question is to assume that it is broken spontaneously. Several examples of both one-loop and tree-level R-symmetry breaking were developed in refs. [4–10] and very minimal models of *direct* mediation (i.e. where the "quarks" of the dynamical SUSY breaking sector play the role of messengers) [11–15] based on a "baryon"-deformation of the ISS model were developed in refs. [5]. These followed earlier developments in refs. [16–28].

A distinction between the phenomenology of the two kinds of model was drawn in ref. [8] where it was noted that, whereas the explicit mediation models are rather similar to standard gauge mediation, the direct mediation models can differ significantly, with much heavier scalar superpartners than usual. (Benchmark points were presented in ref. [8] to support this, and also to show that a baryon-deformed ISS model coupled to the MSSM model, provides a fully calculable system of broken supersymmetry.) Several questions remain however which we will address in this paper. At first sight, one might suspect that this kind of spectrum indicates a residual approximate R-symmetry in the model, possibly because it is broken spontaneously at one-loop — indeed this would seem to be a mildly split version of the argument presented in ref. [29]. On closer inspection however, the precise reason for the suppression of gaugino masses is a little more complicated. Moreover the ISS-like DSB sector itself may become phenomenologically important because, in direct mediation, it contains states charged under SM gauge groups that are light (typically of order 1 TeV).

This paper follows the story to its logical conclusion: we will catalogue the possible ways that such supersymmetry and R-symmetry breaking ends up in the visible sector, using various exemplary models of different types of breaking and gauge mediation (direct or indirect). We conclude that direct mediation generically yields phenomenology quite different from normal gauge mediation. This is due partly to the R-symmetry and partly to the fact that in direct mediation one of the fields to which the messengers couple is a pseudo-Goldstone mode. Generally the visible sector phenomenology ranges from a mildly split spectrum to a very heavy scalar (split-SUSY like) spectrum. In addition, in direct mediation the pseudo-Goldstone modes are expected to enter the visible spectrum, giving a rich source of new TeV mass particles associated with the SUSY breaking sector. This

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[1]Strictly speaking it is the mass scale governing the identification of the composite meson $Q\tilde{Q}$ of the electric ISS theory with the elementary meson Φ of the magnetic theory, $Q\tilde{Q} = \Lambda_{ISS}\Phi$

is similar to the effects of light pseudomoduli which have been found in [30] in the context of explicit R-symmetry breaking models.

We will also note that explicit mediation and spontaneously broken R-symmetry can be problematic in ISS-like models, due to the possibility that messengers become tachyonic. Thus the best prospect for indirect gauge mediation (i.e. with explicit messengers) is explicit R-symmetry breaking of the form discussed in ref. [3].

1.1 Overview

Our point of reference for the present paper is, the model of ref. [5], which introduced into the ISS superpotential a so-called "baryon deformation" that projected out some of the R-symmetry to satisfy the condition that some fields get R-charges different from 0 and 2 [4]. This baryon-deformed, or *ISSb* model, is a natural deformation of the ISS model which at tree-level has a runaway to *broken* supersymmetry. Upon adding the Coleman-Weinberg contributions to the potential, the runaway direction is stabilized at large field values where the R-symmetry is spontaneously broken. If part of the flavour symmetry of the ISS model is gauged and identified with the parent SU(5) of the Supersymmetric Standard Model (SSM), the magnetic quarks can then be enlisted to play the role of messengers, providing an extremely simple model of direct mediation. Moreover it was shown in ref. [31] that the Landau pole problem that usually plagues direct gauge mediation can be avoided: this is because the ISS model itself runs into a Landau pole above which a well-understood electric dual theory takes over. This results in a nett reduction in the effective number of messenger flavours coupling to the SSM *above* the scale Λ_{ISS}, and this in turn prevents the Standard Model coupling running to strong coupling — a scenario dubbed "deflected gauge unification".

In this paper we would like to generalize these observations to a much wider class of models. In order to do this we will begin in the following section by introducing an alternative way to break the R-symmetry of the ISS model spontaneously, by adding a meson term (with some singlet fields) to the superpotential. We call this the "meson-deformed" ISS model, or *ISSm* model. This bears some resemblance to the class of models considered previously in ref. [9], although now the R-symmetry is broken radiatively rather than at tree-level, thus allowing it to be somewhat simpler. We will show how supersymmetry and R-symmetry are broken, using both an analytic tree-level analysis and then a numerical minimization of the full Coleman-Weinberg potential.

We then, in section 3, go on to show how the supersymmetry breaking can subsequently be mediated, first in Subsection 3.1 with an explicit (indirect) mediation where we introduce an additional messenger sector, and then in Subsection 3.2 with direct mediation. In the former case the phenomenology is similar to the standard gauge mediation picture [32] – that is gauginos and scalars have similar masses governed by a single scale and related by functions of the gauge couplings and group theory indices. In particular the absence of tachyonic messenger states requires the explicit mediation model to lie in this regime, and we argue that this is likely to require additional *explicit* R-symmetry violating messenger mass terms. (In this case the spontaneous R-symmetry breaking that we have so carefully

arranged would become irrelevant.) Thus indirect gauge mediation in the ISS model works best with explicit R-symmetry breaking of the form discussed in ref. [3].

On the other hand in Subsection 3.2 we find that the directly mediated meson-deformed model *does* avoid tachyons without explicit R-symmetry breaking and gives phenomenology of a different sort, similar to that of the baryon-deformed model: the gaugino masses are suppressed. We then turn to one of the main goals of the paper which is to answer the question of *why* gauginos are so light compared to the scalar spectrum, and to see if this is a generic feature of spontaneously broken R-symmetry, or is more to do with how the mediation occurs. In fact we shall see that both aspects play a role: it occurs only with direct mediation, but is also related in a rather indirect way to the fact that the R-symmetry is broken spontaneously. The R-symmetry and the equations of motion enforce certain relations between the F-terms which makes the gaugino masses cancel at leading order in messenger mass-insertions. Once the one-loop Coleman-Weinberg contributions to the potential are included the F-terms violate these classical relations and generate non-trivial contributions to gaugino masses at the leading order in mass-insertions (the numerical analysis shows that the precise behaviour is rather complicated). It is this effect which gives the leading contribution to the gaugino masses.

In addition, in this gaugino suppressed regime, we shall find that the contribution from the adjoint pseudo-Goldstone modes , whose mass is lifted only at one-loop, can become important. In Subsection 3.2.2 we consider this second question in more detail. We shall see that the pseudo-Goldstone modes can have a significant impact on the SSM mass spectrum, and indeed their mediated contribution to the gaugino mass can be dominant in precisely the direct mediation models where the gauginos are light. This is because their one-loop suppressed mass makes them behave like a mediating sector with a correspondingly lower messenger scale. Moreover in order to give TeV scale SUSY breaking in the visible sector, the scale of hidden SUSY breaking is typically taken to be order $\frac{16\pi^2}{g^2}$ TeV or slightly higher. Thus the one-loop suppressed masses of the pseudo-Goldstone modes are typically around the scale of SUSY-breaking in the visible sector. This is a generic prediction: models of direct gauge mediation predict additional (with respect to the MSSM) scalar and fermion states in the visible sector, corresponding to pseudo-Goldstone modes, whose masses are close to the weak scale. We also note that the gaugino masses do not necessarily obey the usual relation where their mass ratios scale with the ratios of the coupling constants. Finally in section 4 we repeat the entire analysis for the baryon-deformed model. We find that the picture is similar.

Thus we conclude that indirect explicit mediation gives the standard picture of gauge mediation and that explicit R-symmetry breaking masses for messengers are most likely required. On the other hand direct mediation leads to a mass spectrum with heavy scalars and suppressed gaugino masses. Here the R-symmetry breaking *can* be spontaneous, in which case the pseudo-Goldstone modes can play a significant role in the mediation and in the visible sector phenomenology.

2 Meson-deformed ISS theory as the susy-breaking sector

As summarized in the Introduction, there are two simple types of deformation one might contemplate adding to the ISS model in order to make it spontaneously break R symmetry and generate Majorana gaugino masses in the visible sector. The first was presented in refs. [5, 8] and corresponds to adding a baryonic operator to the original model. That possibility will be examined and extended in section 4. Here we will discuss an alternative possibility which is to add appropriate mesonic deformations to the original model.

We will work entirely in the low-energy magnetic (i.e. relevant to collider phenomenology) description of the ISS model [1]; it contains N_f flavours of quarks and anti-quarks, φ and $\tilde{\varphi}$ respectively, charged under an $SU(N)$ gauge group, as well as an $N_f \times N_f$ meson Φ_{ij} which is a singlet under this gauge group. This is an $SU(N)$ gauge theory with $N = N_f - N_c$ which is weakly coupled in the IR. The ISS superpotential is given by

$$W_{ISS} = h(\Phi_{ij}\varphi_i.\tilde{\varphi}_j - \mu_{ij}^2\Phi_{ji}). \tag{2.1}$$

The coupling h is related to the different dynamical scales in the electric and magnetic theories (or equivalently the mapping between the two gauge couplings). The parameter μ_{ij}^2 is derived from a Dirac mass term $m_Q Q\tilde{Q}$ for the quarks of the electric theory: $\mu^2 \sim \Lambda_{ISS} m_Q$ where the meson field $\Phi_{ij} = \frac{1}{\Lambda_{ISS}}Q_i\tilde{Q}_j$ and where Λ_{ISS} is the Landau pole of the theory. Equation (2.1) gives the tree-level superpotential of the magnetic ISS SQCD theory; there is also the non-perturbatively generated

$$W_{\text{dyn}} = N\left(\frac{\det_{N_f} h\Phi}{\Lambda_{ISS}^{N_f - 3N}}\right)^{\frac{1}{N}}, \tag{2.2}$$

which gives negligible[2] contributions to physics around the SUSY-breaking vacuum.

The flavour symmetry of the magnetic model is initially $SU(N_f)$. When we do *direct* mediation, see section 3.2, an $SU(5)_f$ subgroup of this symmetry is gauged and identified with the parent $SU(5)$ of the Standard Model, so that $N_f \geq N+5$. On the other hand *indirect* mediation, considered in section 3.1, involves the introduction of *explicit* messengers and in that case N_f is a free parameter.

To visualise the the general set-up, let us first consider a simple example, which is appropriate for either case: we shall choose an $SU(2)$ gauge group for the magnetic dual theory and $N_f = 7$ flavours, with the flavour symmetry broken by μ_{ij} to $SU(2)_f \times SU(5)_f$. We will refer to this as the 2-5 model which was the also the prototype model[3] considered in refs. [5, 8]. The matter field decomposition under the $SU(2)_f \times SU(5)_f$ flavour subgroup and the charge assignments under $SU(2)_{\text{gauge}} \times SU(2)_f \times SU(5)_f \times U(1)_B \times U(1)_R$ are given in table 1. Note that we use an f-suffix to stand for "flavour" but one should remember

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[2]The only exception to this is the R-axion field. For this the explicit R-symmetry breaking contained in W_{dyn} gives a contribution to the mass [5] which importantly facilitates the evasion of astrophysical bounds [33–35]. For a recent discussion of the R-axion detection prospects at the LHC see [36].

[3]We will show momentarily that the meson-deformed ISS model actually requires a slightly more general flavour-breaking pattern which can be described by 1-1-5 and 2-2-3 models or their generalisations. For baryon-deformations all of these models, including the simplest 2-5 scenario will also work.

2-5 Model	$\mathrm{SU}(2)_{\mathrm{mg}}$	$\mathrm{SU}(2)_f$	$\mathrm{SU}(5)_f$	$\mathrm{U}(1)_B$	$\mathrm{U}(1)_R$
$\Phi_{ij} \equiv \begin{pmatrix} Y & Z \\ \tilde{Z} & X \end{pmatrix}$	$\mathbf{1}$	$\begin{pmatrix} Adj + \mathbf{1} & \bar{\square} \\ \square & \mathbf{1} \end{pmatrix}$	$\begin{pmatrix} \mathbf{1} & \square \\ \bar{\square} & Adj + \mathbf{1} \end{pmatrix}$	0	2
$\varphi \equiv \begin{pmatrix} \phi \\ \rho \end{pmatrix}$	$\square$	$\begin{pmatrix} \bar{\square} \\ \mathbf{1} \end{pmatrix}$	$\begin{pmatrix} \mathbf{1} \\ \bar{\square} \end{pmatrix}$	$\frac{1}{2}$	R
$\tilde{\varphi} \equiv \begin{pmatrix} \tilde{\phi} \\ \tilde{\rho} \end{pmatrix}$	$\bar{\square}$	$\begin{pmatrix} \square \\ \mathbf{1} \end{pmatrix}$	$\begin{pmatrix} \mathbf{1} \\ \square \end{pmatrix}$	$-\frac{1}{2}$	$-R$

Table 1. The 2-5 Model. We show the ISS matter field decomposition under the gauge SU(2), the flavour $\mathrm{SU}(2)_f \times \mathrm{SU}(5)_f$ symmetry, and their charges under the $\mathrm{U}(1)_B$ and R-symmetry. Both of the U(1) factors above are defined as tree-level symmetries of the magnetic ISS formulation in eq. (2.1). The (small) non-perturbative anomalous effects described by eq. (2.2) are not included. In the absence of baryon-deformations, the R-charges of magnetic quarks, $\pm R$, are arbitrary and can always be re-defined by considering instead a linear combination of $\mathrm{U}(1)_B$ and $\mathrm{U}(1)_R$ factors.

that in direct mediation $\mathrm{SU}(5)_f$ contains the gauge group of the Standard Model.

In the case of the 2-5 model, by a gauge and flavour rotation, the matrix μ_{ij}^2 can be brought to a diagonal 2-5 form:

$$2 - 5 \text{ Model}: \qquad \mu_{ij}^2 = \begin{pmatrix} \mu_Y^2 \mathbf{I}_2 & 0 \\ 0 & \mu_X^2 \mathbf{I}_5 \end{pmatrix}, \quad \mu_Y^2 > \mu_X^2 \,. \tag{2.3}$$

Now consider adding the following deformation[4] involving the meson plus some additional singlet fields A, B, C:

$$W_{\mathrm{meson-def}} = h(m_1 A^2 + m_2 BC + \lambda AB \operatorname{tr}(\Phi)) \,. \tag{2.4}$$

Here we chose to scale all the superpotential parameters with h. The meson deformation of the ISS model is characterised by the dimensionless coupling constant λ. In the electric-dual ISS formulation this deformation is $\sim \frac{1}{M_{Pl}} AB \operatorname{tr}(Q\tilde{Q})$ and thus

$$\lambda \sim \frac{\Lambda_{ISS}}{M_{Pl}} \ll 1 \,. \tag{2.5}$$

The new singlet fields are constrained to have R-charges given in table 2; these are different from 0 or 2, so spontaneous R-symmetry breaking is a possibility [4, 10].

The combined effect of $W_{ISS} + W_{\mathrm{meson-def}}$, gives a *generic* R-symmetry preserving superpotential which defines the low-energy magnetic formulation of our meson-deformed ISS theory. This is a self-consistent approach since, as pointed out in ref. [8], R-symmetry

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[4]A similar deformation involving a meson operator and two singlet fields was previously considered in ref. [6]. Their model, however, contained a runaway direction to a supersymmetric vacuum. For generic values of parameters, this makes the non-supersymmetric R-breaking vacuum of [6] short-lived and unstable to decay in the runaway direction. We will see below that our version of the meson-deformed model defined by eqs. (2.4), (2.1) with a 2-2-3 or 1-1-5 flavour patterns does not have a supersymmetric runaway, and the resulting susy-breaking vacuum is stabilised.

	$\mathrm{U}(1)_R$
A	1
B	-1
C	3

Table 2. R-charges of A, B, C singlet fields of the meson deformation in eq. (2.4).

breaking in the electric theory is controlled by a small parameter.[5] Terms quadratic in the meson Φ that could arise from lower dimensional irrelevant operators in the electric theory are forbidden by R-symmetry. Thus, our deformation is described by a *generic* superpotential and $W_{ISS} + W_{\mathrm{meson-def}}$ gives its leading-order terms.

Being an exact symmetry of the tree-level magnetic superpotential, the R-symmetry of this model is actually spontaneously-broken, as we have already alluded to above. We shall consider this R-symmetry breaking before we discuss the SUSY breaking and its mediation.

First note that for any non-zero $\langle AB \rangle$ we can define an effective μ^2 term

$$\mu_{\mathrm{eff}}^2 = \mu^2 - \lambda \langle A \rangle \langle B \rangle. \tag{2.6}$$

Thus the magnetic quarks acquire VEVs precisely as they do in the undeformed ISS but with μ^2 replaced by μ_{eff}^2;

$$\langle \rho \rangle = \langle \tilde{\rho} \rangle = 0 \tag{2.7}$$

$$\langle \phi \tilde{\phi} \rangle = \mu_{\mathrm{Y\,eff}}^2. \tag{2.8}$$

The VEVs of $\mathrm{tr}(\Phi)$ and C will simply set $\langle F_A \rangle = \langle F_B \rangle = 0$; that is

$$\langle \mathrm{tr}(\Phi) \rangle = -\frac{2m_1 \langle A \rangle}{\lambda \langle B \rangle} \tag{2.9}$$

$$\langle C \rangle = -\frac{\lambda \langle A \rangle \langle \mathrm{tr}(\Phi) \rangle}{m_2} = \frac{2m_1 \langle A \rangle^2}{m_2 \langle B \rangle}. \tag{2.10}$$

At this point the full potential is

$$V = \sum_{i=3}^{7} h^2 |(\mu_{\mathrm{eff}}^2)_{ii}|^2 + |F_C|^2 = 5h^2 |\mu_X^2 - \lambda \langle AB \rangle|^2 + h^2 m_2^2 |B|^2 \,, \tag{2.11}$$

so there is a runaway to unbroken SUSY in the direction $B \to 0$ and $A = \mu_X^2 / \lambda B \to \infty$ along which the R-symmetry is broken.

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[5]In principle, it is known that the apparent R-symmetry of the magnetic formulation of the ISS SQCD is an approximate symmetry of the underlying electric theory: it is broken by the anomaly as per eq. (2.2). (At the same time, the anomaly-free combination of $\mathrm{U}(1)_R$ and the axial symmetry $\mathrm{U}(1)_A$ is broken explicitly by the mass terms of electric quarks m_Q.) However, the R-symmetry is broken in the electric theory in a controlled way [8] by small parameter, $m_Q/\Lambda_{ISS} = \mu^2/\Lambda_{ISS}^2 \ll 1$. As such the R-symmetry is preserved to that order in the superpotential.

Now, in order to end up with broken SUSY we would like to stabilize this type of runaway with Coleman-Weinberg terms in the one-loop potential. (Note that alternatively one could stabilize the model at tree-level using a more complicated potential and R-symmetry as discussed in ref. [9].) We therefore need a runaway to *broken* SUSY since the Coleman-Weinberg contributions vanish where SUSY is unbroken. The classical runaway vacuum becomes non-supersymmetric if the components of the $\mu^2_{X\,ij}$ matrix on the right hand side of eq. (2.11) are no longer degenerate. This is easily achieved by breaking the flavour group into three rather than two factors.

For example, one can consider a 2-2-3 model. Here the original $SU(7)_f$ of the ISS $SU(2)_{mg}$ gauge theory is broken to $SU(2)_f \times SU(2)_f \times SU(3)_f$. This realisation can be thought of as the 2-5 model above where the $SU(5)_f$ flavour subgroup was further broken to $SU(2)_f \times SU(3)_f \times U(1)_{traceless}$ by splitting the eigenvalues of the μ^2_{ij} matrix. This does not cause problems for either explicit or direct mediation. Indeed in the case of direct gauge mediation the $SU(2)_L$ and $SU(3)_c$ components of μ^2_{ij} (or equivalently m_Q in the electric theory) renormalize differently below the GUT scale and so they are not expected to be the same.[6]

Alternatively, one can consider an even simpler example of a 1-1-5 model with $N_f = 7$ and $N_c = 6$ so that the magnetic 'number of colours', $N = 1$, and the magnetic group is trivial. By splitting the eigenvalues of the μ^2_{ij} matrix we choose the flavour breaking to have the 1-1-5 pattern, $SU(7)_f \to U(1)_f \times U(1)_f \times SU(5)_f$. For the case of direct mediation the SM gauge group is $SU(5)_f$.

To give a unified treatment of the 1-1-5 and the 2-2-3 models one can consider a general N-P-X model with $N + N_P + N_X = N_f$ and the μ^2_{ij} matrix given by:

$$\mu^2_{ij} = \begin{pmatrix} \mu^2_Y \mathbf{I}_N & 0 & 0 \\ 0 & \mu^2_P \mathbf{I}_{N_P} & 0 \\ 0 & 0 & \mu^2_X \mathbf{I}_{N_X} \end{pmatrix}, \quad \mu^2_Y > \mu^2_P, \mu^2_X, \quad \mu^2_P \neq \mu^2_X, \tag{2.12}$$

which corresponds to $SU(N_f) \to SU(N)_f \times SU(N_P)_f \times SU(N_X)_f$ as well as traceless $U(1)$ combinations which commute with the right hand side of eq. (2.12). For simplicity, the rank of top left Y-corner is identified with N, the number of magnetic colours, thus the original ISS rank condition which is responsible for the SUSY-breaking vacuum is arranged so that $F_\Phi = 0$ when $\Phi = Y$, see eq. (2.8), and $F_\Phi \neq 0$ when Φ is either P or X. The corresponding decomposition of ISS magnetic matter fields and their charges for this models are given in table 3.

The minimization with respect to C and $\text{tr}(\Phi)$ are as in eqs. (2.9)-(2.10) before, but minimization with respect to A, results in

$$\langle A \rangle = \frac{N_P \mu^2_P + N_X \mu^2_X}{N_P + N_X} \frac{1}{\lambda \langle B \rangle}, \tag{2.13}$$

---

[6]Note that renormalization of μ^2 above the scale Λ_{ISS} would be understood as renormalization of m_Q in the electric theory.

N-P-X Model	$\mathrm{SU}(N_P)_f$	$\mathrm{SU}(N_X)_f$	$\mathrm{SU}(N)_{\mathrm{mg}}$	$\mathrm{U}(1)_B$	$\mathrm{U}(1)_R$
$\Phi_{ij} \equiv \begin{pmatrix} Y & N & Z \\ \tilde{N} & P & M \\ \tilde{Z} & \tilde{M} & X \end{pmatrix}$	$\begin{pmatrix} 1 & \square & 1 \\ \bar{\square} & Adj+1 & \square \\ 1 & \bar{\square} & 1 \end{pmatrix}$	$\begin{pmatrix} 1 & 1 & \square \\ 1 & 1 & \square \\ \bar{\square} & \bar{\square} & Adj+1 \end{pmatrix}$	1	0	2
$\varphi \equiv \begin{pmatrix} \phi \\ \sigma \\ \rho \end{pmatrix}$	$\begin{pmatrix} 1 \\ \bar{\square} \\ 1 \end{pmatrix}$	$\begin{pmatrix} 1 \\ 1 \\ \bar{\square} \end{pmatrix}$	$\square$	$\frac{1}{N}$	R
$\tilde{\varphi} \equiv \begin{pmatrix} \tilde{\phi} \\ \tilde{\sigma} \\ \tilde{\rho} \end{pmatrix}$	$\begin{pmatrix} 1 \\ \square \\ 1 \end{pmatrix}$	$\begin{pmatrix} 1 \\ 1 \\ \square \end{pmatrix}$	$\bar{\square}$	$-\frac{1}{N}$	$-R$

Table 3. The N-P-X Model. We indicate ISS matter field decomposition under the flavour subgroup $\mathrm{SU}(N_P)_f \times \mathrm{SU}(N_X)_f$. In direct mediation we would gauge $\mathrm{SU}(N_P)_f \times \mathrm{SU}(N_X)_f \times \mathrm{U}(1)_{\mathrm{traceless}}$ or its subgroup, and identify it with the SM gauge group. We also show the gauge $\mathrm{SU}(N)$ and the charges under the $\mathrm{U}(1)_B$ and R-symmetry as in table 1.

and consequently the potential

$$
\begin{aligned}
V &= \sum_{i=N+1}^{N_f} h^2 |(\mu^2_{\mathrm{eff}})_{ii}|^2 + |F_C|^2 \\
&= h^2 N_P \left(\mu_P^2 - \frac{N_P \mu_P^2 + N_X \mu_X^2}{N_P + N_X} \right)^2 + h^2 N_X \left(\mu_X^2 - \frac{N_P \mu_P^2 + N_X \mu_X^2}{N_P + N_X} \right)^2 + h^2 m_2^2 |B|^2 \\
&= h^2 \frac{N_P N_X}{N_P + N_X} (\mu_X^2 - \mu_P^2)^2 + h^2 m_2^2 |B|^2 .
\end{aligned}
\tag{2.14}
$$

Again there is a runaway but now to broken supersymmetry as desired.

Note that in the case of explicit mediation the flavour symmetries in the ISS sector are divorced from the gauge symmetries of the Standard Model. In that case one can have a breaking of flavour symmetry that is more general than eq. (2.12), in terms of μ_{ii}^2. Defining the average μ_{ii} of the unbroken $\mathrm{SU}(N_f - N)$ factor as

$$
\overline{\mu^2} = \frac{1}{N_f - N} \sum_{i=N+1}^{N_f} \mu_{ii}^2 ,
\tag{2.15}
$$

we have

$$
\langle A \rangle = \frac{\overline{\mu^2}}{\lambda \langle B \rangle}
\tag{2.16}
$$

and then the generalisation of eq. 2.14 reads

$$
V = h^2 \sum_{i=N+1}^{N_f} (\mu_{ii}^2 - \overline{\mu^2})^2 + h^2 m_2^2 |B|^2 .
\tag{2.17}
$$

It is worth re-emphasizing that even in the limit $A, C \to \infty$ and $B \to 0$ the scalar potential V is non-zero, so we have a runaway to *broken* SUSY. Proceeding to one-loop, the

Coleman-Weinberg contribution to the potential is therefore expected to lift and stabilize this direction at the same time as lifting the pseudo-Goldstone modes.

The Coleman-Weinberg effective potential [37] sums up one-loop quantum corrections into the following form:

$$V_{\text{eff}}^{(1)} = \frac{1}{64\pi^2} \operatorname{STr} \mathcal{M}^4 \log \frac{\mathcal{M}^2}{\Lambda^2} \equiv \frac{1}{64\pi^2} \left(\operatorname{Tr} m_{\text{sc}}^4 \log \frac{m_{\text{sc}}^2}{\Lambda^2} - 2 \operatorname{Tr} m_{\text{f}}^4 \log \frac{m_{\text{f}}^2}{\Lambda^2} + 3 \operatorname{Tr} m_{\text{v}}^4 \log \frac{m_{\text{v}}^2}{\Lambda^2} \right) \tag{2.18}$$

where Λ is the UV cutoff,[7] and the scalar, fermion and vector mass matrices are given by [38]:

$$m_{\text{sc}}^2 = \begin{pmatrix} W^{ab}W_{bc} + D^{\alpha a}D^{\alpha}{}_c + D^{\alpha a}_c D^{\alpha} & W^{abc}W_b + D^{\alpha a}D^{\alpha c} \\ W_{abc}W^b + D^{\alpha}{}_a D^{\alpha}{}_c & W_{ab}W^{bc} + D^{\alpha}{}_a D^{\alpha c} + D^{\alpha c}_a D^{\alpha} \end{pmatrix} \tag{2.19}$$

$$m_{\text{f}}^2 = \begin{pmatrix} W^{ab}W_{bc} + 2D^{\alpha a}D^{\alpha}{}_c & -\sqrt{2}W^{ab}D^{\beta}{}_b \\ -\sqrt{2}D^{\alpha b}W_{bc} & 2D^{\alpha c}D^{\beta}{}_c \end{pmatrix} \qquad m_{\text{v}}^2 = D^{\alpha}{}_a D^{\beta a} + D^{\alpha a}D^{\beta}{}_a. \tag{2.20}$$

As usual, $W_c \equiv \partial W/\partial \Phi^c = F_{\Phi^c}^\dagger$ denotes a derivative of the superpotential with respect to the scalar component of the superfield Φ^c and the raised indices denote Hermitian conjugation, i.e. $W^{ab} = (W_{ab})^\dagger$. The D-terms are $D^\alpha = g z_a T^{\alpha a}{}_b z^b$ and they can be formally switched off by setting the gauge coupling $g = 0$, which we shall do for simplicity. All the above mass matrices will generally depend on field expectation values. The effective potential $V_{\text{eff}} = V + V_{\text{eff}}^{(1)}$ is the sum of the F-term (tree-level) potential and the Coleman-Weinberg contributions. To find the vacua of the theory we now have to minimize V_{eff}.

Now we can check the lifting of the classical runaway direction by quantum effects in the Coleman-Weinberg potential. We have done this numerically using *Mathematica* and have also checked it with *Vscape* program of ref. [39]. The non-supersymmetric vacuum is stabilised and in table 4 we give values of the VEVs for the 1-1-5 meson-deformed ISS model for a specific choice of external parameters. It is worth noting at this point that all the tree-level relations we have just derived get slightly shifted by the one-loop minization. As we shall see, these one-loop effects often give the leading contribution to the mediation of SUSY-breaking and so it is important to keep track of them. This is shown in table 4 where in the generic N-P-X model VEVs develop along the direction

$$\begin{aligned} \langle \tilde{\phi} \rangle &= \xi\, \mathbf{I}_N & \langle \phi \rangle &= \kappa\, \mathbf{I}_N \\ \langle Y \rangle &= \eta\, \mathbf{I}_N & \langle P \rangle &= p\, \mathbf{I}_{N_P} & \langle X \rangle &= \chi\, \mathbf{I}_{N_X}, \end{aligned} \tag{2.21}$$

accompanied by the A, B, C VEVs as before. These are the most general VEVs consistent with the tree-level minimization.

3 Models of mediation: from the meson-deformed ISS to the Standard Model

In the context of gauge mediation one can consider two distinct scenarios of how supersymmetry and R-symmetry breaking is transmitted to the visible Standard Model sector.

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[7]Which is traded for a renormalization scale at which the couplings are defined.

Vev	$\kappa/\mu_X = \xi/\mu_X$	η/μ_X	p/μ_X	χ/μ_X	A/μ_X	B/μ_X	C/μ_X
Tree-level constrained	4.7610	0	-0.1283	-4.8242	30.7086	7.5983	248.22
Unconstrained	4.7607	0.0026	-0.1129	-4.8283	30.7973	7.5617	248.96

Table 4. The 1-1-5 Model: Stabilized VEVs for a meson-deformed ISS theory with $N_f = 7$, $N_c = 6$, $h = 1$, $m_1/\mu_X = m_2/\mu_X = 0.03$, $\mu_Y/\mu_X = 5$, $\mu_P/\mu_X = 3$ and $\lambda = 0.01$. We show both the *constrained* VEVs (i.e. the VEVs obtained when the tree-level relations are enforced) and the true *unconstrained* VEVs resulting from complete minimization.

The first class is ordinary gauge mediation (i.e. mediation with explicit messengers), and the second class involves the models of direct gauge mediation. In this section we discuss how these two possibilities can be realized for the SUSY breaking models we have outlined in the previous section

3.1 Gauge mediation with explicit messengers

We begin in this subsection with explicit mediation. In this scenario one imagines that there is a third sector – the messenger fields — that is responsible for generating the SUSY breaking operators required in the visible sector. The approach in this paper is to try to have a preserved R-symmetry that is broken spontaneously. What we shall find is that we fall foul of the tachyonic messenger problem: ultimately we have to reintroduce explitcit R-symmetry breaking messenger masses to avoid this and we are forced back to the explicit mediation scenario of ref. [3].

To show this, let is first introduce an additional set of mediating fields f and $\tilde{f}$ transforming in the fundamental (and antifundamental respectively) of the Standard Model gauge groups. For concreteness we can take f and $\tilde{f}$ to be (anti)-fundamentals of the underlying GUT gauge group, e.g. $SU(5)_{\mathrm{GUT}}$. In explicit medation these messengers couple to the ISS sector via additional messenger coupling in the superpotential

$$W_{\mathrm{mess}} = \mathrm{Tr}(\tau\Phi)\, f \cdot \tilde{f}\,, \tag{3.1}$$

where τ_{ij} is an arbitrary coupling which from the electric theory perspective should scale as Λ_{ISS}/M_{Pl} as in ref. [3]. We remind the reader that there are no constraints on this coupling coming from the Standard Model, and that the ISS parameters, such as N, N_f are essentially unconstrained.

In order to see how the SUSY breaking enters the visible sector we need to exhibit the mass matrices for messenger fields explicitly. At tree-level the SUSY breaking enters into the scalar mass-squared matrices through the non-zero F_Φ-terms to which the messenger fields, f and $\tilde{f}$ couple. In general the matrices are given by (ignoring the D-terms)

$$m_{\mathrm{sc}}^2 = \begin{pmatrix} W^{ab}W_{bc} & W^{abc}W_b \\ W_{abc}W^b & W_{ab}W^{bc} \end{pmatrix},$$
$$m_{\mathrm{f}} = W_{ab}\,, \tag{3.2}$$

with the W_{ac} being the SUSY preserving mass of the fermions, and the off-diagonal terms $W^{abc}W_b$ containing the SUSY breaking. In this case $W_{f\tilde{f}} = \mathrm{Tr}(\langle\tau\Phi\rangle)$ is the Dirac mass of the fermionic superpartners, ψ_f and $\psi_{\tilde{f}}$, and the SUSY breaking contribution appears first in the tree-level mass-squared of the scalars, $S = (f, \tilde{f}^*)$. We have:

$$m_{\mathrm{sc}}^2 = \begin{pmatrix} |\mathrm{Tr}(\tau\Phi)|^2 & \mathrm{Tr}(\tau^\dagger F_\Phi^\dagger) \\ \mathrm{Tr}(\tau F_\Phi) & |\mathrm{Tr}(\tau\Phi)|^2 \end{pmatrix} . \tag{3.3}$$

Now, in order to avoid tachyonic messengers we must here impose the usual explicit mediation constraint that

$$|\mathrm{Tr}(\tau\langle\Phi\rangle)|^2 > |\mathrm{Tr}(\tau\langle F_\Phi\rangle)| \tag{3.4}$$

which is effectively a *lower* bound on the amount of spontaneous R-symmetry breaking (since $\langle\Phi\rangle$ is charged under R-symmetry). In particular this generally prevents us arranging a split scenario with gauginos much lighter than squarks and sleptons, since this would be a signature of approximate R-symmetry. (The situation is drastically different in models of direct mediation as we shall see in the following sections.)

As we have said dimensional arguments give

$$\tau \sim \lambda \sim \Lambda_{ISS}/M_{Pl} \ll 1$$

so the tachyonic inequality is delicate. If one assumes that $\Phi \sim \mu$ then it seems that the inequality is actually always violated when $\tau \ll 1$. But note that the same inequality can be equivalently written in terms of singlet VEVs,

$$\tau\Phi \sim \tau m_1 \frac{A}{\lambda B} \sim \frac{\tau}{\lambda^2} \frac{m_1 \mu^2}{B^2} , \tag{3.5}$$

which shows that the situation is quite complicated and can only be analyzed numerically. For the values in table 4 taking $\tau \sim \lambda$ violates the inequality which suggests that it may be problematic in general to avoid tachyonic messengers.

An explicit R-breaking mass term is a way to overcome this tachyon so that, as in ref. [3], eq. (3.1) becomes

$$W_{\mathrm{mess}} = \mathrm{Tr}(\tau\Phi)\, f \cdot \tilde{f} + M_f\, f \cdot \tilde{f} \tag{3.6}$$

Hence explicit gauge mediation and spontaneous R-symmetry breaking are inconsistent when the DSB is based on the ISS model. Note that we could have also added a term $\frac{A^2 f \cdot \tilde{f}}{M_{Pl}}$; however since we have $\langle A\rangle \sim \mu_P \ll \Lambda$ the effective mass that this induces for the messengers is even smaller than $\langle\mathrm{Tr}(\tau\Phi)\rangle$.

From here on the calculation of the SUSY spectrum is rather standard with values for gaugino masses being generated being of the same order as those for scalar masses; and so one expects a similar phenomenology to normal explicit gauge mediation [32], with the diagram that induces the gaugino mass in the present explicit mediation case as shown in figure 1.

However there is one feature of the present set-up that is rather interesting. The SUSY breaking effects in the visible sector, i.e. the gaugino and squark masses, are all proportional
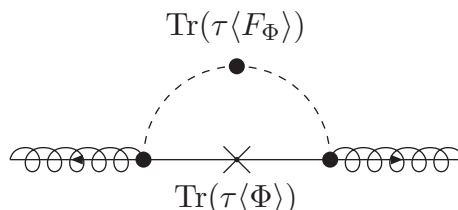
$$\text{Tr}(\tau \langle F_\Phi \rangle)$$

$$\text{Tr}(\tau \langle \Phi \rangle)$$

Figure 1. One-loop contribution to the gaugino masses from messengers f, $\tilde{f}$. The dashed (solid) line is a bosonic (fermionic) messenger. The blob on the scalar line indicates an insertion of the F-term VEV into the propagator of the scalar messengers and the cross denotes an insertion of the R-symmetry breaking VEV into the propagator of the fermionic messengers.

to the combination $W^{abc}W_b = Tr(\tau^\dagger F_\Phi^\dagger)$. But as we have seen in the previous section, the F-terms at the minimum (with VEV-less messengers, so that the SM gauge groups are not Higgsed) are given at tree-level by

$$F_{\Phi_{ij}}^\dagger = h\delta_{ij}(\mu_{ii}^2 - \overline{\mu^2})\,, \tag{3.7}$$

which clearly obeys

$$\text{Tr}(F_\Phi^\dagger) = 0\,. \tag{3.8}$$

This can be seen to result from the minimization of the tree-level potential with respect to A for a given B VEV:

$$\frac{\partial V}{\partial A} = \lambda B \text{Tr}(F_\Phi^\dagger) = 0\,. \tag{3.9}$$

Thus (at tree-level) the mediation of SUSY-breaking to the visible sector requires non-degenerate couplings τ_{ii}, and indeed we can write

$$\text{Tr}(\tau F_\Phi) = h(\overline{\tau\mu^2} - \bar{\tau}\overline{\mu^2})\,. \tag{3.10}$$

That is, only if *both* τ and μ have non-degeneracy can there be unsuppressed SUSY breaking mediation, even though SUSY breaking *per se* requires non-degeneracy only in the latter.

However, as we have said, when the full minimization is performed, tree-level relations such as $\text{Tr}(F_\Phi^\dagger) = 0$ are no longer expected to hold (for example, with the unconstrained values in the table we find $\text{Tr}(F_\Phi^\dagger) = -0.034\mu_X^2$): typically one finds $\text{Tr}(F_\Phi^\dagger) = \mu^2/(16\pi^2)$, since the effective F-term for mediation is one-loop suppressed. Thus when the τ are degenerate one can still get $m_\lambda \sim \frac{\mu^2}{16\pi^2 M_f} \frac{g^2}{(16\pi^2)} \sim 1\,\text{TeV}$ if $\mu^2/M_f \sim 10^7\,\text{GeV}$.

3.2 Direct gauge mediation

Now, let us compute gaugino masses for the direct gauge mediation scenario from the meson-deformed ISS sector. We first consider the effects of those direct messengers which obtain R-symmetry breaking masses at tree-level and which couple directly to the largest F-terms. These transform in the fundamental representation of the SM gauge groups,

and this constitutes a strictly one-loop and formally leading order effect. Then we will include additional, formally higher-loop, contributions from the pseudo-Goldstone modes transforming in both adjoint and (bi-)fundamental representations of the Standard Model gauge groups. It will turn out that the latter contributions can be of the same order.

3.2.1 Strict one-loop contributions to gaugino masses

To present a general discussion relevant for any deformation of the ISS model, by mesons, baryons or otherwise, we shall consider models of the form

$$W = h\Phi_{ij}\varphi_i.\tilde{\varphi}_j - h\mu_{ij}^2\Phi_{ji} + W_{\mathrm{meson-def}}(A_a, \Phi) + W_{\mathrm{baryon-def}}(A_a, \phi, \tilde{\phi}) \qquad (3.11)$$

where A_a denote generic singlets. The superpotential depends on Φ linearly, this is dictated by the R-symmetry of the model and is a central feature of direct mediation in the ISS context.

To keep the presentation simple in what follows we shall concentrate here on the 1-1-5 model, so that the parent gauge symmetry of the SM (in this case $\mathrm{SU}(5)_f$) is non-split. This discussion can also be straightforwardly generalised to the 2-2-3 and other N-P-X models by an appropriate reassembling of building blocks below.

The all important messenger/SUSY-breaking coupling in the superpotential is in this class of models is

$$\frac{1}{h}W \supset \Phi_{ij}\varphi_i.\tilde{\varphi}_j \supset \rho X\tilde{\rho} + \phi Z\tilde{\rho} + \rho\tilde{Z}\tilde{\phi} + \phi Y\tilde{\phi}. \qquad (3.12)$$

The field Φ is the pseudo-Goldstone mode, although note that F_ϕ and $F_{\tilde{\phi}}$ are non-zero as well as F_Φ – this will be important in what follows.

Gaugino masses are generated at one-loop order as indicated in figure 2. The fields propagating in the loop are fermion and scalar components of the direct mediation 'messengers'. Since gaugino masses are forbidden by R-symmetry one crucial ingredient in their generation is the presence of non-vanishing R-symmetry breaking VEVs. We are at this point interested in the contribution to the gaugino mass coming from those messenger fields transforming in the fundamental of SU(5), which formally give the leading-order contribution. (We shall consider the contribution from the X fields separately in section 3.2.2.)

First we exhibit the mass matrices of messenger fields. As before, they are given by (ignoring the D-terms)

$$m_{\mathrm{sc}}^2 = \begin{pmatrix} W^{ab}W_{bc} & W^{abc}W_b \\ W_{abc}W^b & W_{ab}W^{bc} \end{pmatrix}, \qquad m_{\mathrm{f}} = W_{ac}. \qquad (3.13)$$

The fundamental messengers are ρ, $\tilde{\rho}$ and Z, $\tilde{Z}$: we may define a messenger fermion multiplet,

$$\psi = (\rho_i, Z_i)_{\mathrm{ferm}},$$
$$\tilde{\psi} = (\tilde{\rho}_i, \tilde{Z}_i)_{\mathrm{ferm}}, \qquad (3.14)$$
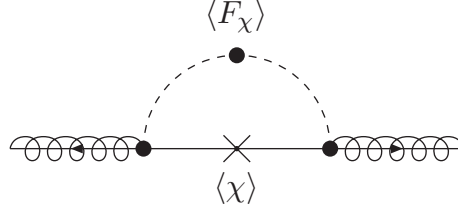
– 14 –

Figure 2. One-loop contribution to the gaugino masses. The dashed (solid) line is a bosonic (fermionic) messenger. The blob on the scalar line indicates an insertion of $\langle F_\chi \rangle$ into the propagator of the scalar messengers and the cross denotes an insertion of the R-symmetry breaking VEV into the propagator of the fermionic messengers.

where $i = 1 \ldots 5$. Then $\mathcal{L} \supset \psi \, m_{\rm f} \tilde{\psi}^T$ where the fermion messenger mass matrix is

$$m_{\rm f} = \mathbf{I}_5 \otimes \begin{pmatrix} \chi & \xi \\ \kappa & 0 \end{pmatrix}, \tag{3.15}$$

written in terms of the VEVs χ, κ and ξ (c.f. (2.21)):

$$\langle X \rangle = \chi I_5, \qquad \langle \phi \rangle = \kappa, \qquad \langle \tilde{\phi} \rangle = \xi. \tag{3.16}$$

For the scalar mass-squared matrix, we can define equivalent multiplets

$$S = (\rho_i, Z_i, \tilde{\rho}_i^*, \tilde{Z}_i^*)_{\rm sc}. \tag{3.17}$$

To proceed one can diagonalise the mass matrices and compute the full one-loop contribution to the gaugino mass. That is we define the diagonalisations:

$$\hat{m}_{\rm sc}^2 = Q^\dagger m_{\rm sc}^2 Q \tag{3.18}$$

$$\hat{m}_{\rm f} = U^\dagger m_{\rm f} V \tag{3.19}$$

with eigenvectors

$$\hat{S} = S.Q$$
$$\hat{\psi}_+ = \psi.U$$
$$\hat{\psi}_- = \tilde{\psi}.V^* \tag{3.20}$$

Here, the $m_{\rm f}$ diagonalisation is in general a biunitary transformation.

In order to calculate the gaugino mass, we need the gauge interaction terms given by

$$\mathcal{L} \supset i\sqrt{2} g_A \lambda_A (\psi_1 T^A S_1^* + \psi_2 T^A S_2^* + \tilde{\psi}_1 T^{*A} S_3 + \tilde{\psi}_2 T^{*A} S_4) + H.C. \tag{3.21}$$

$$= i\sqrt{2} g_A \lambda_A (\hat{\psi}_{+i} \hat{S}_k^* (U_{i1}^\dagger Q_{1k} + U_{i2}^\dagger Q_{2k}) + \hat{\psi}_{-i} \hat{S}_k (Q_{k3}^\dagger V_{1i} + Q_{k4}^\dagger V_{2i})) + H.C. \tag{3.22}$$

Then the diagram in figure 2 amounts to[8]

$$M_{\lambda_A}^{(\rho,Z)} = 4N g_A^2 \operatorname{tr}(T^A T^B) \sum_{ik} (U_{i1}^\dagger Q_{1k} + U_{i2}^\dagger Q_{2k})(Q_{k3}^\dagger V_{1i} + Q_{k4}^\dagger V_{2i}) \, I(\hat{m}_{\mathrm{f},i}, \hat{m}_{\mathrm{sc},k}) \quad (3.23)$$

where $I(\hat{m}_{\mathrm{f}}, \hat{m}_{\mathrm{sc}})$ is the appropriate one-loop integral with a fermion and a scalar. Here the "N" reinstates the possibility of an $SU(N)_{mg}$ gauge group. In the diagonal mass-basis

$$I(a,b) = \int \frac{d^4 k}{(2\pi)^4} \frac{a}{k^2 - a^2} \frac{1}{k^2 - b^2} = \frac{-a(\eta+1)}{16\pi^2} + \frac{1}{16\pi^2} \frac{a}{a^2 - b^2} \left[a^2 \log\left(\frac{a^2}{\Lambda^2}\right) - b^2 \log\left(\frac{a^2}{\Lambda^2}\right) \right] \quad (3.24)$$

and

$$\eta = \frac{2}{4-D} + \log(4\pi) - \gamma_E. \quad (3.25)$$

This integral is UV-divergent, but the divergences cancels in the sum over eigenstates as required.

Using (3.23) we can now evaluate gaugino masses in figure 2 generated by fundamental messengers ρ, $\tilde{\rho}$ and Z. Numerical values for the gaugino mass for a few different values of parameters of the model are given in the tables in section 3.3.

It is instructive to complement these numerical calculations by a simple analytic estimate, and in particular explain the smallness of these gaugino mass contributions. When the F-terms are small compared to μ^2 one can expand eqs. (3.23)-(3.24). We define a matrix of 'weighted' F-terms as:

$$\mathcal{F}^{ab} = W^{abc} W_c, \quad (3.26)$$

and to the leading order in $\mathcal{F}$ obtain,

$$M_{\lambda_A} = \frac{g_A^2}{8\pi^2} N \operatorname{tr}(T^A T^B) \operatorname{Tr}(\mathcal{F} \cdot m_{\mathrm{f}}^{-1}) + \mathcal{O}(\mathcal{F}^3). \quad (3.27)$$

This is a well-known leading order in $\mathcal{F}$ approximation which is basis-independent. In the appendix we give the derivation of eq. (3.27) in the general settings relevant to our model(s).

Clearly the matrix $\mathcal{F}$ is determined entirely by the contribution in eq. (3.12) to be

$$\mathcal{F} = W^{abc} W_c = h \begin{pmatrix} F_\chi & F_{\tilde{\phi}} \\ F_\phi & 0 \end{pmatrix} \quad (3.28)$$

and since $m_{\mathrm{f}}^{-1} = \begin{pmatrix} 0 & \frac{1}{\kappa} \\ \frac{1}{\xi} & -\frac{\chi}{\xi\kappa} \end{pmatrix}$ we find

$$M_{\lambda_A}^{(\rho,Z)} = \frac{g_A^2}{8\pi^2} N \operatorname{tr}(T^A T^B) \left(\frac{F_{\tilde{\phi}}}{\xi} + \frac{F_\phi}{\kappa} \right) + \mathcal{O}(\mathcal{F}^3) \quad (3.29)$$

Now consider the minimization condition for the tree-level potential, $V = \sum_c |F^c|^2$ with respect to Y^*.

$$\frac{1}{2} \frac{\partial V}{\partial Y^*} = 0 = \sum_c W^{Yc} F_c = \kappa F_{\tilde{\phi}} + \xi F_\phi + W_{\mathrm{meson-def}}^{Y A_a} F_{A_a} \quad (3.30)$$

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[8]More precisely, there are actually two diagrams of this type which are mirror images of each other.

(For the constrained 1-1-5 VEVs shown in table 4 this trivially sets $\eta = 0$.) This equation together with eq. (3.29) implies that the tree-level leading order gaugino mass is zero

$$M_{\lambda_A}^{(\rho,Z)} = 0 + \mathcal{O}(\mathcal{F}^3) \tag{3.31}$$

unless the additional singlet fields appearing in the meson deformation have non-zero F-terms as well. (This would require an additional source of SUSY breaking beyond the O'Raifeartaigh breaking of the ISS sector, and is therefore unattractive.) As we have stressed, these relations are perturbed when the potential is stabilized by one-loop effects (e.g. η is non-zero in the unconstrained model of table 4): then the estimate in eq. (3.29) is still reasonably good, with the F-terms being derived from the one-loop equations.

This leading order suppression for the gaugino mass explains the relative smallness of our numerical results in table 5 which shows the "reduced gaugino masses" $m_{1/2}$ defined by

$$M_{\lambda_A} = \frac{g_A^2}{16\pi^2} m_{1/2} \,. \tag{3.32}$$

In particular these values are much smaller than those derived for the scalars in table 6 where we show the "reduced scalar masses" m_0 defined by

$$m_{\text{sferm}}^2 = \sum_A \frac{g_A^4}{(16\pi^2)^2} C_A S_A m_0^2 \,, \tag{3.33}$$

where C_A and S_A are the standard Casimir/Dynkin indices as in ref. [40]. We note that this suppression is also related to that in ref. [41], which tells us that F_Φ does not contribute to the gaugino masses at leading order because of the structure of m_{f} (in particular the zero entry). Here we find that the argument extends to quite general models of direct mediation.

3.2.2 Additional contributions to gaugino masses

The effects considered above have so far generated rather small contributions to gaugino masses. Thus, we have to consider additional contributions, due to the adjoint X and P as well as the bifundamental M and $\tilde{M}$ messengers. These messengers are massless at tree-level and acquire masses only at loop-level. Thus their contributions to gaugino masses are formally a higher-loop effect. After a careful consideration we find that these indeed give a contribution to the gaugino masses which comparable to the strict one-loop effect described above. Scalar masses being unsuppressed at leading order are not significantly effected.

For 1-1-5 type models where the SM gauge group is $\text{SU}(5)_f$, the new contributions arise from the X_{ij} fields with $i, j = 1 \ldots 5$. They contribute through the diagram shown in figure 3. Note that the scalar vertex exists because the Coleman-Weinberg potential induces an R-symmetry violating mass term. The fermion mass-propagator is also absent at tree-level: since it is a Majorana term (and the X-fermions have R-charge 1) it also violates R-symmetry and by the non-renormalization theorem it vanishes in the absence of both R-symmetry and supersymmetry breaking. The naive expectation is therefore that this contribution will be three-loop suppressed. As we shall see, this is not the case, and in fact the contribution can be competitive with the previous contributions. This is because
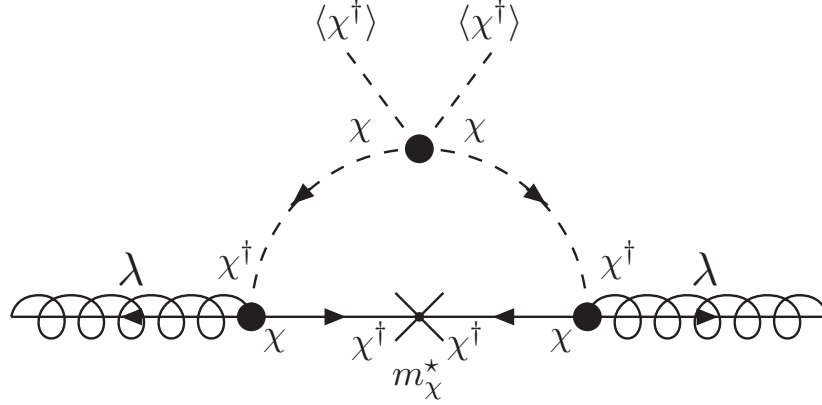
Figure 3. One-loop contribution to the gaugino masses from X-messengers. The dashed (solid) line is a bosonic (fermionic) component of X. The blob on the scalar line indicates an insertion of $\langle F_\chi \rangle$ into the propagator of the scalar messengers and the cross denotes an insertion of the R-symmetry breaking VEV into the propagator of the fermionic messengers.

the X modes are pseudo-Goldstone modes: *all* their masses arise at one-loop, and the lightness of these modes corresponds to a suppression of the effective messenger scale of the adjoints whose mass is in fact similar to $M_{\rm SUSY}$.

Let us estimate these effects in more detail. First the mass-insertions: the scalar mass-squareds come from the Coleman-Weinberg term

$$V_{\rm eff}^{(1)} \supset {\rm STr}\left(\frac{\mathcal{M}^4}{64\pi^2} \log \mathcal{M}^2 \right). \tag{3.34}$$

In particular there are terms involving $\bar{W}_{\rho Z} W^{Z\rho} \bar{W}_{\rho\tilde\rho} W^{\tilde\rho\rho} = h^4 \xi^2 |\delta X_{ij}|^2$ where $X = \langle X \rangle + \delta X$. Since typically $\xi \gg \mu \gg \kappa$ one expects R-symmetry conserving mass-squared for the adjoints of order

$$m_{XX^*}^2 \sim \frac{h^4 \xi^2}{64\pi^2} \tag{3.35}$$

at the minimum. R-symmetry violating masses are induced by terms such as $W_{\rho\tilde\rho}\bar{W}^{\tilde\rho\rho}W_{\rho\tilde\rho}\bar{W}^{\tilde\rho\rho} \supset h^4\langle X\delta X^\dagger\rangle^2 + h.c = h^4\chi^2(\delta X_{ij}^* \delta X_{ji}^*) + h.c.$ Hence we expect a neutral mass-squared matrix for $\mathbf{X} = (X^A, X_A^*)$ (where A is the adjoint index) of the form

$$m_{\mathbf{X}}^2 \sim \frac{\delta_{AB}}{64\pi^2} \begin{pmatrix} a & b \\ b^* & a \end{pmatrix},$$
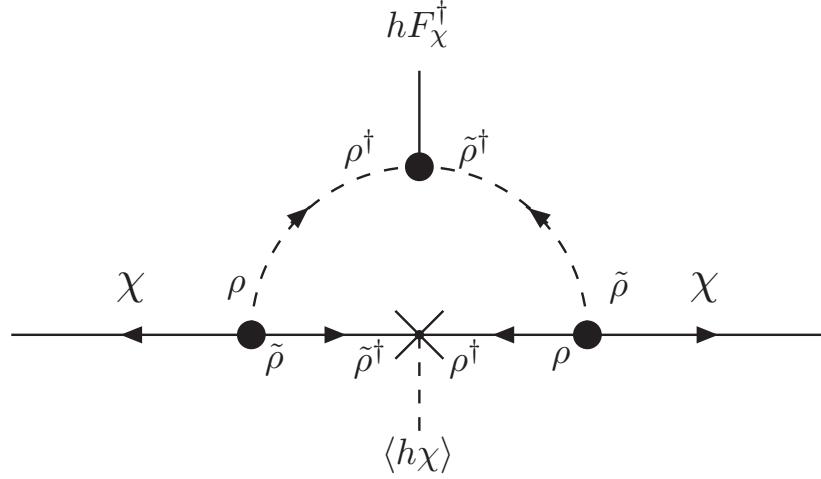$$a \sim \xi^2 \, ; \quad b \sim \chi^2 \, . \tag{3.36}$$

Figure 4. One-loop contribution to the Majorana masses of X-fermions. The dashed (solid) line is a bosonic (fermionic) messenger. The blob on the scalar line indicates an insertion of $\langle F_\chi \rangle$ into the propagator of the scalar messengers and the cross denotes an insertion of the R-symmetry breaking VEV into the propagator of the fermionic messengers.

Assuming b is real, the diagonalization of this matrix is $\hat{m}_{\mathbf{X}}^2 = (\mathcal{Q}^X)^T m_{\mathbf{X}}^2 \mathcal{Q}^X = \frac{h^4}{64\pi^2}\mathrm{diag}(a+b,\, a-b)$ where

$$\mathcal{Q}^X = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & -1 \\ 1 & 1 \end{pmatrix}. \tag{3.37}$$

We will call the two eigenvalues $\hat{m}_{X_\pm}^2$.

The R-breaking mass term for the adjoint fermion is generated from diagram shown in figure 4. The topology is identical to the one-loop gaugino diagram with internal states ψ, $\tilde{\psi}$ and S with the mass matrices and diagonalisations as in eqs. (3.18) and (3.19), although of course the vertices are different: they come from the $W \supset h\rho X\tilde{\rho}$ coupling and are given by

$$V \supset h\, (X\psi_1)S_3^* + h(X\tilde{\psi}_1)S_1 + h.c. \tag{3.38}$$

In terms of the previous mass eigenstates these become

$$V \supset hX\, (\hat{\psi}_{+i}\hat{S}_k^*(U_{i1}^\dagger Q_{3k}) + \hat{\psi}_{-i}\hat{S}_k(Q_{k1}^\dagger V_{1i}) + H.C.$$

where the diagonalisation matrices Q, U and V are exactly the same as in eqs. (3.18)–(3.19). Defining $X_{ij} = \sqrt{2}\, X_A T_{ij}^A$, and with standard Feynman parametization we find that the diagram in figure 4 generates

$$M_{\psi_X} = 4Nh^2\mathrm{tr}(T^A T^B) \sum_{ik} (U_{i1}^\dagger Q_{3k})(Q_{k1}^\dagger V_{1i})\, I(\hat{m}_{\mathrm{f},i}, \hat{m}_{\mathrm{sc},k})\,, \tag{3.39}$$

where $I(\hat{m}_{\mathrm{f},i}, \hat{m}_{\mathrm{sc},k})$ is the same integral (3.24) as in (3.23).

Note that although the diagram in figure 4 is similar to the fundamental contribution to the one-loop gaugino mass, there is less suppression. This is because the couplings of $\rho, \tilde{\rho}$ and $Z, \tilde{Z}$ to X are not degenerate as they are for the gaugino, indeed there is no equivalent of the $h\rho X\tilde{\rho}$ coupling for the $Z, \tilde{Z}$ fields at all; hence unitarity does not operate in the same way. Following the same steps as for the gaugino in the appendix we obtain a non-vanishing leading order result in $\mathcal{F}$,

$$M_{\psi_X^A} \approx 4\, N\, h^2 \text{tr}(T^A T^B) \sum_{ijk} \mathcal{A}_{jk}(U_{i1}^\dagger V_{1j})(U_{k1}^\dagger V_{1i})\,(\hat{m}_{\text{f}})_i\, J(\hat{m}_{\text{f}\,i}^2,\, \hat{m}_{\text{f}\,j}^2,\, \hat{m}_{\text{f}\,k}^2) \tag{3.40}$$

where the matrix $\mathcal{A}_{ij}$ was defined in eq. (A.3) and the function J is given by

$$J(a,b,c) = \frac{1}{8\pi^2} \frac{a^2 b^2 \log\left(\frac{a}{b}\right) + a^2 c^2 \log\left(\frac{c}{a}\right) + b^2 c^2 \log\left(\frac{b}{c}\right)}{(a^2 - b^2)(a^2 - c^2)(b^2 - c^2)}. \tag{3.41}$$

A very rough simple estimate is

$$M_{\psi_X^A} \sim \frac{h^2 \chi}{32\pi^2} \frac{F_X}{\xi^2}. \tag{3.42}$$

This should be compared to the equivalent contribution to the gaugino mass in section 3.2.1 which did vanish at this order (see eqs. (3.27),(3.31)).

Having determined the masses of X messengers we can now make an estimate for their contribution to the gaugino mass. The general expression is

$$M_{\lambda_A}^{(X)} = g_A^2\, N_X\, \left(I(M_{\psi_X}, \hat{m}_{X_+}) - I(M_{\psi_X}, \hat{m}_{X_-}) \right), \tag{3.43}$$

where N_X is the rank of the X lower-right corner in eq. (2.12), which in the case of 1-1-5 type models is $N_X = 5$.

Equation (3.43) allows us to evaluate gaugino masses generated by adjoint X-messengers. Numerical values for the full mass expressions (without relying on estimates and expansions in $\mathcal{F}$) for the model given in table 4 are presented in table 5 in section 3.3. In this table we give contributions from the ρ and Z messengers in the first column and from the X messengers in the second column. The third column gives the similar contribution from M messengers which we will comment on momentarily (see eq. (3.46)). The last column is the total result. Other tables in the same subsection follow the same structure and give results for other models.

To understand the order of magnitude can also be understood with the help of the following analytical estimates. As we have seen the masses are of the order

$$M_{\psi_X^A} \sim \frac{h^3 \chi}{32\pi^2} \frac{\mu^2}{\xi^2}$$

$$\hat{m}_{X_\pm}^2 \sim \frac{h^4}{64\pi^2}(\xi^2 \pm \chi^2).$$

Thus for $h \lesssim 1$ we expect $M_{\psi_X^A}^2 \ll \hat{m}_{X_\pm}^2$ and we find

$$M_{\lambda_A}^{(X)} = \frac{g_A^2\, N_X}{32\pi^2}\, M_{\psi_X^A} \log\left(\frac{\hat{m}_{X_+}^2}{\hat{m}_{X_-}^2}\right) \sim \frac{g_A^2 h^3\, N_X}{2(16\pi^2)^2} \frac{\chi^3 \mu^2}{\xi^4}, \tag{3.44}$$

where the last expression is valid for $\chi \lesssim \xi$. Note that, although in a "mass-insertion approximation" the leading order diagram is in principle three-loop, there is only a two-loop $1/(16\pi^2)^2$ suppression.

In addition to the contribution from the adjoint X fields we have a contribution from the M and $\tilde{M}$ fields. As can be seen from table 3 these are bifundamentals under the SU(N_X) and SU(N_P) groups.

$$M_{\psi_M} = Nh^2 \left[\sum_{ik} (U_{i1}^{P\dagger} Q_{3k}^X)(Q_{k1}^{X\dagger} V_{1i}^P) I(\hat{m}_{\mathrm{f},i}^P, \hat{m}_{\mathrm{sc},k}^X) + (X \leftrightarrow P) \right]. \tag{3.45}$$

Here, the labels P and X indicate the diagonalization matrizes for the SU(N_P) and SU(N_X) blocks, respectively (see table 3). In particular, the X is the diagonalization for the ρ and Z messengers whereas P corresponds to the σ and N.

The corresponding contribution to the gaugino mass is,

$$M_{\lambda^A}^M = 4\,\mathrm{tr}(T^A T^B) N_P \sum_{k=1}^{2} \mathcal{Q}_{1k}^M (\mathcal{Q}^M)_{k2}^T I(M_{\psi_M}, \hat{m}_{M,kk}), \tag{3.46}$$

where $\mathcal{Q}^M$ is the M-analog of $\mathcal{Q}^X$ matrix given in eq. (3.37). As mentioned earlier these contributions are shown in the third column of table 5 and similar ones in the subsection 3.3.

3.2.3 Scalar masses

Having determined the gaugino masses in the preceding subsections, we now outline the procedure for the generation of sfermion masses of the supersymmetric standard model. As in ref. [8] we follow the calculation of Martin in ref. [40] adapted to our direct mediation models.

Sfermion masses are generated by the two-loop diagrams shown in figure 5. In [40] the contribution of these diagrams to the sfermion masses was determined to be,

$$m_{\tilde{f}}^2 = \sum_{mess.} \sum_{a} g_a^4 C_a S_a(mess.)[\text{sum of graphs}], \tag{3.47}$$

where we sum over all gauge groups under which the sfermion is charged, g_a is the corresponding gauge coupling, $C_a = (N_a^2 - 1)/(2N_a)$ is the quadratic Casimir and $S_a(mess.)$ is the Dynkin index of the messenger fields (normalized to $1/2$ for fundamentals).

As in the calculation of the gaugino mass we use the propagators in the diagonal form and insert the diagonalisation matrices directly at the vertices. For the diagrams 5a to 5f we have closed loops of purely bosonic or purely fermionic mass eigenstates of our messenger fields. It is straightforward to check that in this case the unitary matrices from the diagonalisation drop out. We then simply have to sum over all mass eigenstates the results for these diagrams computed in ref. [40].

The next diagram 5g is slightly more involved. This diagram arises from the D-term interactions. D-terms distinguish between chiral and antichiral fields, in our case ρ, Z and $\tilde{\rho}, \tilde{Z}$, respectively. We have defined our scalar field S in (3.17) such that all component fields have equal charges. Accordingly, the ordinary gauge vertex is proportional
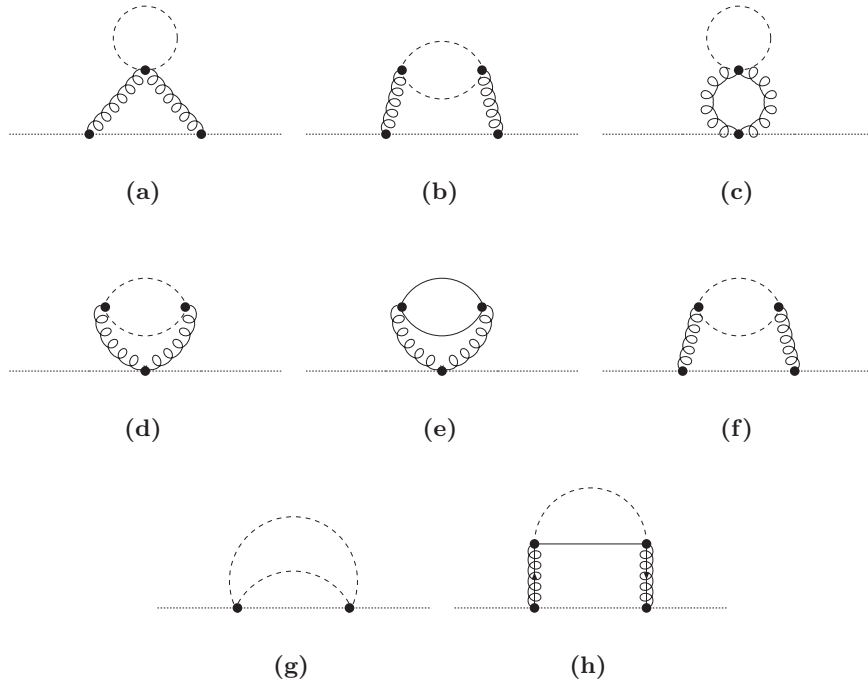
Figure 5. Two-loop diagrams contributing to the sfermion masses. The long dashed (solid) line is a bosonic (fermionic) messenger. Standard model sfermions are depicted by short dashed lines.

to a unit matrix in the component space (cf. eq. (3.21)). This vertex is then 'dressed' with our diagonalisation matrices when we switch to the $\hat{S}$ basis, (3.22). This is different for diagram 5g. Here we have an additional minus-sign between chiral and antichiral fields. In field space this corresponds to a vertex that is proportional to a matrix $V_D = \mathrm{diag}(1, 1, -1, -1)$. We therefore obtain,

$$\text{Figure } 5\text{g} = \sum_{i,m}(Q^T V_D Q)_{i,m} J(\hat{m}_{0,m}, \hat{m}_{0,i})(Q^T V_D Q)_{m,i}, \tag{3.48}$$

where J is the appropriate two-loop integral for figure 5g which can be found in [40].

Finally, in 5h we have a mixed boson/fermion loop. The subdiagram containing the messengers is similar to the diagram for the gaugino mass. The only difference is the direction of the arrows on the gaugino lines. Indeed the one-loop sub-diagram corresponds to a contribution to the kinetic term rather than a mass term for the gauginos. (The mass term will of course contribute as well but will be suppressed by quark masses.) Using eq. (3.22) we find,

$$\text{Figure } 5\text{h} = \sum_{ik}(|U^\dagger_{i1}Q_{1k} + U^\dagger_{i2}Q_{2k}|^2 + |Q^\dagger_{k3}V_{1i} + Q^\dagger_{k4}V_{2i}|^2)L(\hat{m}_{1/2,i}, \hat{m}^2_{0,k}), \tag{3.49}$$

where L is again the appropriate loop integral from [40].

Summing over all diagrams we find the sfermion masses which are typically significantly larger than the gaugino masses calculated earlier. Indeed, the scalar masses roughly follow

Contribution (in units of μ_X)	$\rho,\ \tilde{\rho},\ Z,\ \tilde{Z}$	X	$\tilde{M}\ M$	total
Tree-level constrained	8.22×10^{-5}	0	0	8.22×10^{-5}
Unconstrained (tree scalar mass matrix)	5.34×10^{-3}	0	0	5.34×10^{-3}
Unconstrained (mass matrix with CW)	2.81×10^{-3}	4.49×10^{-3}	8.3×10^{-5}	7.38×10^{-3}

Table 5. Contributions to the reduced gaugino mass $m_{1/2}$ for the meson-deformed 1-1-5 model of table 4.

the estimate

$$m_{\tilde{f}}^2 \sim \frac{g^4}{(16\pi^2)^2}\mu^2. \tag{3.50}$$

This is precisely the leading order effect which in our direct mediation scenario is absent for the gaugino masses.

So far we have taken into account the ρ, Z (or similarly the σ, M) contributions which as we just explained give a non-vanishing leading order effect. In distinction to our earlier calculation of the gaugino masses we do not need to include the sub-dominant contributions from other messengers (which were massless at tree-level).[9]

3.3 Summary of signatures in the directly mediated meson-deformed model

Here we present and summarize our result for gaugino and sfermion masses for a variety of our meson-deformed models. These results are most conveniently expressed in terms of the reduced gaugino ($m_{1/2}$)

$$M_{\lambda_A} = \frac{g_A^2}{16\pi^2}m_{1/2}\,, \tag{3.51}$$

and scalar masses (m_0^2)

$$m_{\text{sferm}}^2 = \sum_A \frac{g_A^4}{(16\pi^2)^2}C_A S_A m_0^2. \tag{3.52}$$

We similarly define reduced masses for the pseudo-Goldstone components of the direct messengers (appearing in tables 7, 11, 17) by including a factor of $16\pi^2$,

$$m_{\text{reduced}} = 16\pi^2 m_{\text{phys}} \tag{3.53}$$

The first three tables 5, 6 and 7 summarize our results for the mass spectrum at the high scale for meson-deformed 1-1-5 model specified in table 4.

The following four tables 8, 9, 10 and 11 give results for the same 1-1-5 model but with a different choice of parameters. Comparing the last lines in table 5 and table 9 we see that the contribution from the X messengers can be of the same order but the relative sizes of the different contributions can vary quite significantly.

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[9]Inclusion of such effects would be actually not completely straightforward because our mass-insertion technique breaks down when used in the two-loop diagrams for the scalars. The reason for this can be traced to the non-cancelation of the UV cutoff dependent terms. This problem would disappear if one performs a complete higher-loop calculation. In any case since the leading order result for scalars was non-vanishing we do not expect any significant changes from this.

Contribution (in units of μ_X)	ρ, $\tilde{\rho}$, Z, $\tilde{Z}$
Tree-level constrained	0.48
Unconstrained (tree scalar mass matrix)	0.48
Unconstrained (mass matrix with CW)	not consistent

Table 6. Contributions to the reduced sfermion masses m_0 (only ρ, $\tilde{\rho}$, Z, $\tilde{Z}$ contribution) for the meson-deformed 1-1-5 model of table 4. The third line in the table indicates that the use of the full CW corrected masses is inappropriate in this case (see text).

Particle	Reduced Mass$/\mu_X$
sfermions	0.48
gauginos	7.4×10^{-3}
χ_f	0.13
χ_s	1.33, 2.35
M_f, $\tilde{M}_f$	0.42
M_s, $\tilde{M}_s$	9.58, 9.73

Table 7. Reduced masses for the various particles charged under the SM gauge group for the meson-deformed 1-1-5 model of table 4, with $M_{\mathrm{SUSY}}/\mu_X = 2.7$.

Vev	$\kappa/\mu_X = \xi/\mu_X$	η/μ_X	p/μ_X	χ/μ_X	A/μ_X	B/μ_X	C/μ_X
Tree-level constrained	4.7610	0	-0.0881	-2.5014	17.1430	13.6110	215.92
Unconstrained	4.7603	0.0017	-0.0783	-2.5127	17.1978	13.5634	217.38

Table 8. Stabilized vevs for a meson model with $N_f = 7$, $N_c = 6$, $h = 1$, $m_1/\mu_X = 0.05$, $m_2/\mu_X = 0.01$, $\mu_Y/\mu_X = 5$, $\mu_P/\mu_X = 3$ and $\lambda = 0.01$.

In total both models give rather similar predictions. with scalars being two orders of magnitude heavier than the gauginos. This is a "slightly" split-SUSY scenario which is expected in all of our direct mediation ISS-SSM models.

In addition, as can be seen from tables 7, 11, some of the messengers which are charged under the Standard Model gauge group are relatively light with masses somewhere in between the scalars and the gauginos.

The remaining six tables in this subsection give an example for a 2-2-3 model — a model with a non-trivial magnetic group. This model has very similar features with the only exception being that the reduced gaugino and sfermion masses differ for the different gauge groups. This shows that one can achieve a deviation from the simple scaling of the full physical masses with the gauge couplings eqs. (3.51) and (3.52) because $m_{1/2}$ and m_0 now actually depend on the index A specifying the gauge group.

We have generated the soft SUSY breaking terms of the SSM at the high (messenger)

Contribution (in units of μ_X)	$\rho,\ \tilde{\rho},\ Z,\ \tilde{Z}$	X	$\tilde{M}\ M$	total
Tree-level constrained	5.91×10^{-5}	0	0	5.91×10^{-5}
Unconstrained (tree scalar mass matrix)	3.45×10^{-3}	0	0	3.45×10^{-3}
Unconstrained (mass matrix with CW)	1.78×10^{-3}	7.06×10^{-4}	1.34×10^{-5}	2.50×10^{-3}

Table 9. Contributions to the reduced gaugino mass $m_{1/2}$ for the meson-deformed 1-1-5 model of table 8.

Contribution (in units of μ_X)	$\rho,\ \tilde{\rho},\ Z,\ \tilde{Z}$
Tree-level constrained	0.53
Unconstrained (tree scalar mass matrix)	0.54
Unconstrained (mass matrix with CW)	not consistent

Table 10. Contributions to the reduced sfermion masses m_0 (only ρ, $\tilde{\rho}$, Z, $\tilde{Z}$ contribution) for the meson-deformed 1-1-5 model of table 8.

Particle	Reduced Mass/μ_X
sfermion	0.54
gauginos	2.5×10^{-3}
χ_f	8.83×10^{-2}
χ_s	2.39, 2.71
$M_f,\ \tilde{M}_f$	0.24
$M_s,\ \tilde{M}_s$	10.20, 10.16

Table 11. Reduced masses for the various particles charged under the SM gauge group for the meson-deformed 1-1-5 model of table 8, with $M_{\text{SUSY}}/\mu_X = 2.7$.

Vev	$\kappa/\mu_X = \xi/\mu_X$	η/μ_X	p/μ_X	χ/μ_X	A/μ_X	B/μ_X	C/μ_X
Tree-level constrained	4.5607	0	-1.3999	-4.3327	33.2544	12.6299	105.071
Unconstrained	4.5613	0.0021	-1.3433	-4.2233	33.5579	12.4704	103.038

Table 12. Stabilized vevs for a meson model with $N_f = 7$, $N_c = 5$, $h = 1$, $m_1/\mu_X = 0.03$, $m_2/\mu_X = 0.05$, $\mu_Y/\mu_X = 5$, $\mu_P/\mu_X = 3$ and $\lambda = 0.01$.

scale. In order to determine the mass spectrum at the electroweak scale the soft SUSY breaking parameters given in the tables should be renormalization group evolved. But we expect that the overall pattern remains the same.

In summary, we see that all our direct models have the following features: 1) A heavy scalar spectrum; 2) The pseudo-Goldstone direct messengers are relatively light and the effective low energy theory is always extended away from the MSSM; 3) We can have

Contribution (in units of μ_X)	$\sigma, \tilde{\sigma}, N, \tilde{N}$	P	$\tilde{M}\, M$	total
Tree-level constrained	-4.5×10^{-3}	0	0	-4.5×10^{-3}
Unconstrained (tree scalar mass matrix)	4.1×10^{-3}	0	0	-4.4×10^{-3}
Unconstrained (mass matrix with CW)	-4.52×10^{-2}	-4.17×10^{-4}	-3.6×10^{-4}	-4.60×10^{-2}

Table 13. Contributions to the reduced mass $m_{1/2}^{(2)}$ of the SU(2) gaugino for the meson-deformed 2-2-3 model of table 12.

Contribution (in units of μ_X)	$\rho, \tilde{\rho}, Z, \tilde{Z}$	χ	$\tilde{M}\, M$	total
Tree-level constrained	2.8×10^{-3}	0	0	2.8×10^{-3}
Unconstrained (tree scalar mass matrix)	1.1×10^{-2}	0	0	1.1×10^{-2}
Unconstrained (mass matrix with CW)	1.05×10^{-2}	1.05×10^{-2}	-2.4×10^{-4}	2.1×10^{-2}

Table 14. Contributions to the reduced gluino mass $m_{1/2}^{(3)}$ for the meson-deformed 2-2-3 model of table 12.

Contribution (in units of μ_X)	$\sigma, \tilde{\sigma}, N, \tilde{N}$
Tree-level constrained	2.93
Unconstrained (tree scalar mass matrix)	2.94
Unconstrained (mass matrix with CW)	not consistent

Table 15. Contributions to the reduced masses $m_0^{(2)}$ of the SU(2) sfermions (only $\rho, \tilde{\rho}, Z, \tilde{Z}$ contribution) for the meson-deformed 2-2-3 model of table 12.

Contribution (in units of μ_X)	$\rho, \tilde{\rho}, Z, \tilde{Z}$
Tree-level constrained	1.74
Unconstrained (tree scalar mass matrix)	1.74
Unconstrained (mass matrix with CW)	not consistent

Table 16. Contributions to the SU(3) sfermion masses $m_0^{(3)}$ (only $\sigma, \tilde{\sigma}, N, \tilde{N}$ contribution) for the meson-deformed 2-2-3 model of table 12.

deviations from the standard gaugino/sfermion mass pattern dictated by the Standard Model gauge couplings.

4 The baryon-deformed ISS theory and its mediation patterns

In this section we revisit models with the hidden sector given by baryon-deformed ISS theory introduced in [5, 8]. These models form extensions/deformations of the ISS which are complimentary to the meson deformations discussed above. We will extend the analysis to include the effects of the X and M messengers.

Particle	Reduced Mass/μ_X
sfermions SU(2)	2.95
sfermions SU(3)	1.74
gauginos SU(2)	4.6×10^{-2}
gauginos SU(3)	2.1×10^{-2}
χ_f	0.41
χ_s	14.46, 15.06
P_f	0.62
P_s	5.40, 8.56
M_f, $\tilde{M}_f$	0.47
M_s, $\tilde{M}_s$	11.79, 11.56

Table 17. Reduced masses for the various particles charged under the SM gauge group for the meson-deformed 2-2-3 model of table 12 with $M_{\mathrm{SUSY}}/\mu_X = 2.96$.

4.1 The baryon-deformed model

We start with an ISS model with $N_c = 5$ colours and $N_f = 7$ flavours, which has a magnetic dual description as an SU(2) theory, also with $N_f = 7$ flavours and following [5, 8] we deform this theory by the addition of a baryonic operator. The resulting superpotential is given by

$$W = \Phi_{ij}\varphi_i.\tilde{\varphi}_j - \mu_{ij}^2\Phi_{ji} + m\varepsilon_{ab}\varepsilon_{rs}\varphi_r^a\varphi_s^b \tag{4.1}$$

where $i, j = 1\ldots7$ are flavour indices, $r, s = 1, 2$ run over the first two flavours only, and a, b are SU(2) indices. This is the superpotential of ISS with the exception of the last term which is a baryon of the magnetic SU(2) gauge group. Note that the 1,2 flavour indices and the 3...7 indices have a different status and the full flavour symmetry $SU(7)_f$ is broken explicitly to $SU(2)_f \times SU(5)_f$. As before, the *direct* gauge mediation is implemented by gauging the $SU(5)_f$ factor and identifying it with the parent SU(5) gauge group of the Standard Model. The matter field decomposition under the magnetic $SU(2)_{\mathrm{gauge}} \times SU(5)_f \times SU(2)_f$ and their $U(1)_R$ charges are given in table 1 with $R = 1$.

Using the notation established in the previous sections for the meson model the baryon-deformed model defined by eq. (4.1) is a 2-5 model. It is straightforward to consider alternatives such as a 1-5 model where the magnetic gauge group is empty and the baryon deformation is a linear operator,

$$W_{1-5} = \Phi_{ij}\varphi_i.\tilde{\varphi}_j - \mu_{ij}^2\Phi_{ji} + k\varphi_1, \tag{4.2}$$

or, for example, a 2-2-3 model as before. In all of those models Landau poles inherent in the direct mediation can be avoided by using the deflected unification mechanism of [31]. This works most effectively in the 1-5 model due to its minimal matter content. The discussion of these models is virtually identical to that which we will now present for the 2-5 model.

At the Lagrangian level this baryon-deformed model respects R-symmetry. Thanks to the baryon deformation, the structure of R-charges allows for spontaneous R symmetry

breaking and it was shown in [5] that this does indeed happen. We also stress that our baryon deformation is the leading order deformation of the ISS model that is allowed by R-symmetry of the full theory imposed at the Lagrangian level. As explained in [8] this is a self-consistent approach. For example, terms quadratic in the meson Φ that could arise from lower dimensional irrelevant operators in the electric theory are forbidden by R-symmetry. Thus, our deformation is described by a *generic* superpotential and (4.1) gives its leading-order terms.

Using the $\mathrm{SU}(2)_f \times \mathrm{SU}(5)_f$ symmetry, the matrix μ_{ij}^2 can be brought to the form (2.3). The baryon operator can be identified with a corresponding operator in the electric theory. Indeed the mapping from baryons B_E in the electric theory to baryons B_M of the magnetic theory, is $B_M \Lambda_{ISS}^{-N} \leftrightarrow B_E \Lambda_{ISS}^{-N_c}$ (we neglect factors of order one). Thus one expects

$$m \sim M_{Pl} \left(\frac{\Lambda_{ISS}}{M_{Pl}} \right)^{N_f - 2N} = \frac{\Lambda_{ISS}^3}{M_{Pl}^2}, \tag{4.3}$$

where M_{Pl} represents the scale of new physics in the electric theory at which the irrelevant operator B_M is generated.

The F-term contribution to the potential at tree-level is

$$
\begin{aligned}
V = {} & \sum_{ar} |Y_{rs}\tilde{\phi}_s^a + Z_{r\hat{i}}\tilde{\rho}_{\hat{i}}^a + 2m\varepsilon_{ab}\varepsilon_{rs}\phi_s^b|^2 \\
& + \sum_{a\hat{i}} |\tilde{Z}_{\hat{i}r}\tilde{\phi}_r^a + X_{\hat{i}\hat{j}}\tilde{\rho}_{\hat{j}}^a|^2 + \sum_{as} |\phi_r^a Y_{rs} + \rho_{\hat{i}}^a \tilde{Z}_{\hat{i}s}|^2 + \sum_{a\hat{j}} |\phi_r^a Z_{r\hat{j}} + \rho_{\hat{i}}^a X_{\hat{i}\hat{j}}|^2 \\
& + \sum_{rs} |(\phi_r.\tilde{\phi}_s - \mu_Y^2 \delta_{rs})|^2 + \sum_{r\hat{i}} |\phi_r.\tilde{\rho}_{\hat{i}}|^2 + \sum_{r\hat{i}} |\rho_{\hat{i}}.\tilde{\phi}_s|^2 + \sum_{\hat{i}\hat{j}} |(\rho_{\hat{i}}.\tilde{\rho}_{\hat{j}} - \mu_X^2 \delta_{\hat{i}\hat{j}})|^2
\end{aligned}
\tag{4.4}
$$

where a, b are $\mathrm{SU}(2)_{mg}$ indices. The flavor indices r, s and $\hat{i}$, $\hat{j}$ correspond to the $\mathrm{SU}(2)_f$ and $\mathrm{SU}(5)_f$, respectively. It is straightforward to see that the rank condition works as in ISS; that is the minimum for a given value of X, Y, Z and $\tilde{Z}$ is along $\rho = \tilde{\rho} = 0$ and

$$\langle \phi \rangle = \frac{\mu_Y^2}{\xi} \, \mathbf{I}_2 \,, \qquad \langle \tilde{\phi} \rangle = \xi \, \mathbf{I}_2, \tag{4.5}$$

where ξ parameterizes a runaway direction that will be stabilized by the Coleman-Weinberg potential eq. (2.18). This then gives $Z = \tilde{Z} = 0$. In addition Y becomes diagonal and real (assuming m is real). Defining $\langle Y_{rs} \rangle = \eta \, \mathbf{I}_2$, the full potential is

$$V = 2 \left| \eta\, \xi + 2m\frac{\mu_Y^2}{\xi} \right|^2 + 2 \left| \eta \frac{\mu_Y^2}{\xi} \right|^2 + 5\mu_X^4. \tag{4.6}$$

Using R symmetry we can choose ξ to be real.[10] Minimizing in η we find

$$\eta = -2m \left(\frac{\xi^2}{\mu_Y^2} + \frac{\mu_Y^2}{\xi^2} \right)^{-1}. \tag{4.7}$$

---

[10]The phase of ξ corresponds to the R-axion.

Vev	κ/μ_X	ξ/μ_X	η/μ_X	χ/μ_X
Tree-level constrained	1.1005	8.1781	-0.0793	-0.3493
Unconstrained	1.1004	8.1766	-0.0792	-0.3470

Table 18. Stabilized VEVs for a 2-5 baryon-deformed model with $N_f = 7$, $N_c = 5$, $h = 1$, $m/\mu_X = 0.3$ and $\mu_Y/\mu_X = 3$.

Vev	κ/μ_X	ξ/μ_X	η/μ_X	χ/μ_X
Tree-level constrained	1.76214	5.1074	-0.05248	-0.20720
Unconstrained	1.7620	5.1067	-0.05227	-0.2037

Table 19. Stabilized VEVs for a 1-5 baryon-deformed model with $N_f = 6$, $N_c = 5$, $h = 1$, $k/\mu_X^2 = 0.3$ and $\mu_Y/\mu_X = 3$.

Substituting $\eta(\xi)$ into eq. (4.6) we see that $\xi \to \infty$ is a runaway direction along which

$$V(\xi) = 8m^2\mu_Y^2 \left(\frac{\xi^6}{\mu_Y^6} + \frac{\xi^2}{\mu_Y^2} \right)^{-1} + 5\mu_X^4. \tag{4.8}$$

Since in the limit $\xi \to \infty$, the scalar potential V is non-zero, we have a runaway to broken supersymmetry, hence the Coleman-Weinberg potential again lifts and stabilizes this direction, which is indeed the case [5]. As in eqs. (2.21) we parameterise the pseudo-Goldstone and runaway VEVs by

$$\langle \tilde{\phi} \rangle = \xi \, \mathbf{I}_2 \qquad\qquad \langle \phi \rangle = \kappa \, \mathbf{I}_2 \tag{4.9}$$

$$\langle Y \rangle = \eta \, \mathbf{I}_2 \qquad\qquad \langle X \rangle = \chi \, \mathbf{I}_5. \tag{4.10}$$

Stabilized VEVs for a 2-5 and a 1-5 model are shown in tables 18 and 19, respectively. Constrained VEVs in these tables arise from using the tree-level equations of motion eqs. (4.5) and (4.7). Again, the difference between constrained and unconstrained VEVs is rather small but the general discussion of subsection 3.2 indicates that this difference has crucial effects on the generation of gaugino masses in direct mediation.

Explicit mediation has been studied in [8] and leads to the usual standard GMSB pattern (as also discussed for the meson-deformed model in subsection 3.1).

4.2 Summary of signatures in the directly mediated baryon-deformed model

The basic equations for calculating gaugino and scalar masses are the same as in subsection 3.2. Only the VEV configurations and the structure of the messenger mass matrices know about the difference in the deformation.

Our results for the soft SUSY breaking parameters at the messenger scale are presented below following the same structure as before. The first three tables correspond to the 2-5 model given in table 18. The next three correspond to the 1-5 model specified in table 19.

Contribution	ρ, $\tilde{\rho}$, Z, $\tilde{Z}$	χ	total
Tree-level constrained	4.17×10^{-5}	0	4.17×10^{-5}
Unconstrained (tree scalar mass matrix)	1.74×10^{-3}	0	1.74×10^{-3}
Unconstrained (mass matrix with CW)	-1.57×10^{-3}	9.61×10^{-7}	-1.57×10^{-3}

Table 20. Contributions to the reduced gaugino mass for the baryon-deformed 2-5 model of table 18.

Contribution	ρ, $\tilde{\rho}$, Z, $\tilde{Z}$
Tree-level constrained	0.70
Unconstrained (tree scalar mass matrix)	0.70
Unconstrained (mass matrix with CW)	not consistent

Table 21. Contributions to the reduced sfermion masses (only ρ, $\tilde{\rho}$, Z, $\tilde{Z}$ contribution) for the baryon-deformed 2-5 model of table 18.

Particle	Mass/μ_P
sfermion	0.70
gauginos	1.57×10^{-3}
χ_f	1.92×10^{-2}
χ_s	2.923, 2.925

Table 22. Reduced masses for the various particles charged under the SM gauge group for the baryon-deformed 2-5 model of table 18.

Evidently, the dominant contribution to the gaugino mass comes from unconstraining the VEVs and putting in the full one-loop mass matrices. Overall this leads again to models with heavy scalars and, in distinction to our earlier paper [8] (where the constrained VEVs were used), we do not need to fine tune the different μ^2 parameters to achieve a moderately split spectrum. It is remarkable that in all of the directly mediated ISS models gaugino masses are this sensitive to quantum corrections (due to the inevitable cancellation at tree-level).

5 Conclusions

We have investigated different scenarios of gauge mediation which incorporate a dynamical SUSY breaking (DSB) sector coupled to a supersymmetric Standard Model. The DSB sector was realized in terms of two different types of deformations of the ISS model. These models generate all SUSY breaking parameters at the messenger scale in a calculable way from relatively simple supersymmetric Lagrangians. In all of the models investigated we find rather model independent signatures for the direct gauge mediation which include:

Contribution	ρ, $\tilde{\rho}$, Z, $\tilde{Z}$	χ	total
Tree-level constrained	2.67×10^{-5}	0	2.67×10^{-5}
Unconstrained (tree scalar mass matrix)	7.49×10^{-4}	0	7.49×10^{-4}
Unconstrained (mass matrix with CW)	-5.97×10^{-4}	3.60×10^{-7}	-5.96×10^{-4}

Table 23. Contributions to the reduced gaugino mass for the baryon-deformed 1-5 model of table 19.

Contribution	ρ, $\tilde{\rho}$, Z, $\tilde{Z}$
Tree-level constrained	0.61
Unconstrained (tree scalar mass matrix)	0.61
Unconstrained (mass matrix with CW)	not consistent

Table 24. Contributions to the reduced sfermion masses (only ρ, $\tilde{\rho}$, Z, $\tilde{Z}$ contribution) for the baryon-deformed 1-5 model of table 19.

Particle	Mass/μ_P
sfermions	0.61
gauginos	5.96×10^{-4}
χ_f	1.1×10^{-2}
χ_s	2.921, 2.919

Table 25. Reduced masses for the various particles charged under the SM gauge group for the baryon-deformed 1-5 model of table 19.

- Scalars are typically two orders of magnitude or more heavier than gauginos.

- The low energy effective theory of the visible sector i.e. particles charged under the Standard Model gauge groups is necessarily extended by light pseudo-Goldstone messenger fields.

- Direct mediation models easily allow for deviations from the mass patterns dictated by the gauge couplings, familiar from standard gauge mediation.

It is also possible to implement indirect gauge mediation, by adding an explicit messenger sector. In this case we find a rather standard pattern of gauge mediated supersymmetry breaking.

Finally we would like to briefly comment on how the usual little hierarchy problem of the supersymmetric Standard Model manifests itself. First of all, the non-observation of the Higgs at LEP requires that the mass of the lightest Higgs, $m_{h^0} > 115$ GeV. On the other hand, supersymmetric models predict an upper bound so that

$$(115 \, \text{GeV})^2 \, < \, m_{h^0}^2 \, < \, \cos^2(2\beta) m_Z^2 \, + \, \text{rad. corr.} , \tag{5.1}$$

where the radiative corrections $\sim m_t^2 \log(m_{\tilde{t}}/m_t)$. To fulfill this one needs a rather large stop mass, which our models deliver. On the other hand, the conditions for electroweak symmetry breaking require that at the electroweak scale

$$m_Z^2 = -2(m_{H_u}^2 + |\mu_{\text{MSSM}}|^2) + \mathcal{O}(1/\tan^2(\beta)). \tag{5.2}$$

The scalar masses, including m_{H_u}, and their loop corrections are of the order of $m_{\tilde{t}}$ and are (as just argued) much bigger than the electroweak scale. This requires a fine-tuning of μ_{MSSM} of the order of 10^{-2}. In the direct mediation scenarios with a mildly split SUSY spectrum, $m_{\tilde{t}}$ is bigger than the minimal required value from eq. (5.1) resulting in a somewhat higher degree of fine-tuning of the order of $10^{-4} - 10^{-5}$. In this paper we are treating μ_{MSSM} as a free parameter and do not attempt to solve this problem.

Acknowledgments

A   Leading order contribution to the gaugino mass

To develop a perturbative approximation of eqs. (3.23)–(3.24) we note that when the F-terms are small compared to μ^2, we may first go to the "fermion-diagonal basis", by making a rotation on the scalars given by

$$Q_0 = \begin{pmatrix} U & \mathbf{0} \\ \mathbf{0} & V \end{pmatrix} \tag{A.1}$$

where the U and V matrices are the fermion-diagonalisation matrices defined in (3.19). In this basis the scalar mass-squareds are

$$\tilde{m}_{\text{sc}}^2 = Q_0^\dagger m_{\text{sc}}^2 Q_0 \approx \begin{pmatrix} \hat{m}_{\text{f}}^2 & \mathcal{A} \\ \mathcal{A}^\dagger & \hat{m}_{\text{f}}^2 \end{pmatrix} \tag{A.2}$$

where

$$\mathcal{A}_{ij} = U_{ia}^\dagger W^{abc} W_c V_{bj} = (U^\dagger \mathcal{F} V)_{ij} \tag{A.3}$$

in terms of the F-term matrix $\mathcal{F}^{ab} \equiv W^{abc} W_c$. Evaluating the diagram for the gaugino mass in this basis (cf. eqs. (3.23)-(3.24)) and suppressing the overall factor $2g_A^2 \text{tr}(T^A T^B)$, yields,

$$\int \frac{d^4k}{(2\pi)^4} \sum_{k,l=1}^{4} \sum_{i,j=1}^{2} (U_{i1}^\dagger Q_{0,1k} + U_{i2}^\dagger Q_{0,2k}) \left(\frac{1}{k^2 - \tilde{m}_{\text{sc}}^2} \right)_{kl} \left(\frac{\hat{m}_f}{k^2 - \hat{m}_f^2} \right)_{ij} (Q_{0,l3}^\dagger V_{1j} + Q_{0,l4}^\dagger V_{2j})$$

$$= \int \frac{d^4k}{(2\pi)^4} \sum_{i,j,k,l=1}^{2} (U_{i1}^\dagger U_{1k} + U_{i2}^\dagger U_{2k}) \left(\frac{1}{k^2 - \tilde{m}_{\text{sc}}^2} \right)_{k,(l+2)} \left(\frac{\hat{m}_f}{k^2 - \hat{m}_f^2} \right)_{ij} (V_{l1}^\dagger V_{1j} + V_{l2}^\dagger V_{2j})$$

$$= \int \frac{d^4k}{(2\pi)^4} \sum_{i,j,k,l=1}^{2} \delta_{ik} \left(\frac{1}{k^2 - \tilde{m}_{\text{sc}}^2} \right)_{k,(l+2)} \left(\frac{\hat{m}_f}{k^2 - \hat{m}_f^2} \right)_{ij} \delta_{jl}, \tag{A.4}$$

where, in the last step, we have made use of the unitarity of the U and V matrices.

The fermion propagator is already diagonal, but the boson propagator has off diagonal terms $\sim \mathcal{A}$. Expanding in powers of $\mathcal{A}$ we have,

$$\left(\frac{1}{k^2-\tilde{m}_{\rm sc}^2}\right)_{k,(l+2)} = \left(\frac{1}{k^2-\hat{m}_{\rm f}^2}\mathcal{A}\frac{1}{k^2-\hat{m}_{\rm f}^2} + \frac{1}{k^2-\hat{m}_{\rm f}^2}\mathcal{A}\frac{1}{k^2-\hat{m}_{\rm f}^2}\mathcal{A}^\dagger\frac{1}{k^2-\hat{m}_{\rm f}^2}\mathcal{A}\frac{1}{k^2-\hat{m}_{\rm f}^2}+\cdots\right)_{kl} \tag{A.5}$$

Using that $\hat{m}_{\rm f}$ is a diagonal matrix we find to lowest order in $\mathcal{A}$,

$$M_{\lambda^A} = 2g_A^2 \,{\rm tr}(T^A T^B){\rm Tr}(\mathcal{A}I^{(1)}(\hat{m}_{\rm f})) \tag{A.6}$$

where

$$I_{ij}^{(1)} = {\rm diag}(I(\hat{m}_{ii})) \tag{A.7}$$

and

$$I^{(1)}(m) = \int \frac{d^4k}{(2\pi)^4}\frac{m}{(k^2-m^2)^3} = \frac{1}{32\pi^2}\frac{1}{m}. \tag{A.8}$$

Using the explicit form of $I^{(1)}$ we have the leading order contribution to the gaugino masses:

$$M_{\lambda^A} = \frac{g_A^2}{16\pi^2}\,{\rm tr}(T^A T^B){\rm Tr}(\mathcal{A}\hat{m}_{\rm f}^{-1}) = \frac{g_A^2}{16\pi^2}\,{\rm tr}(T^A T^B){\rm Tr}(\mathcal{F}m_{\rm f}^{-1})\,. \tag{A.9}$$

This reproduces eq. (3.27).

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