Composite Intermediary and Mediator Models of Gauge-Mediated Supersymmetry Breaking

Csaba Csáki, Lisa Randall and Witold Skiba

Center for Theoretical Physics
Laboratory of Nuclear Science and Department of Physics
Massachusetts Institute of Technology
Cambridge, MA 02139, USA
csaki@mit.edu, lisa@ctptop.mit.edu, skiba@mit.edu

Abstract

We discuss gauge-mediated models which employ a tree-level mass term in the superpotential. We give explicit composite realizations in which the mass terms are not fundamental. Instead, they arise as effective terms in the superpotential from confining gauge dynamics.
1 Introduction

Gauge-mediated supersymmetry breaking is in principle an attractive alternative to gravity-mediated supersymmetry breaking in that the flavor changing neutral current (FCNC) problem is automatically solved. However many of the existing models are complicated or contrived so it is worthwhile to explore alternative model-building ideas.

In Ref. [1], two classes of gauge-mediated models were introduced. In one class, called “Intermediary” models, there are two massive singlets, one of which couples to the dynamical supersymmetry breaking sector and the other one to the messenger quarks, but for which there is a Dirac mass term coupling the two. Upon integrating out the singlet there is automatically a coupling between the dynamical supersymmetry breaking (DSB) and visible sectors which can communicate supersymmetry breaking to the messenger quarks and hence to the visible sector. The phenomenology of these models is similar in many respects to other messenger models [2, 3, 4].

In the other class of models, called “Mediator” models, there are massive mediator fields which carry the gauge charge of a gauged messenger group of the DSB sector and also standard model gauge charge. These fields therefore permit communication of supersymmetry breaking, but at high loop order. The scalar partners obtain a two-loop supersymmetry breaking mass, whereas the gauginos obtain a three-loop mass. The phenomenology of these models is therefore distinctive (and probably somewhat more fine-tuned) in that the scalars are expected to be about an order of magnitude heavier than the gauginos.

The advantage of both these models is that there is no complicated superpotential required in order to communicate supersymmetry breaking to a fundamental singlet which is coupled to the messenger quarks. In the intermediary models, the communication is automatic upon integrating out the singlet. In the mediator models, a dynamical messenger sector [5, 6, 7] which automatically communicates supersymmetry breaking to messenger gauginos and ultimately to the visible sector is assumed. The relative simplicity of these scenarios occurs because we have not made the very restrictive assumption about the absence of mass terms in the superpotential. When this assumption is relaxed, more direct communication of supersymmetry breaking in a stable or sufficiently stable vacuum occurs reasonably simply.

However, fundamental mass terms with a scale other than that determined by the Planck
scale or non-perturbative gauge dynamics would be a strong assumption. Our philosophy in generating these models was to first see what works, with the assumption that it would be relatively straightforward to realize the required masses and couplings in a composite scenario.

In this paper, we give explicit realizations of composite dynamics which produce successful Intermediary and Mediator models. These models serve as existence proofs and allow a more accurate determination of how complicated these models are compared to the microscopic realization of alternative gauge-mediated scenarios. Of course it is conceivable that simpler composite implementations exist; the models presented here set an upper bound for complexity.

In the second section we discuss composite intermediary models. We give explicit realizations of the underlying gauge dynamics. We also explain the motivation for the various elements assumed in these models. In the third section we discuss composite mediator scenarios.

Alternative models based on a dynamical messenger sector generated by an inverted hierarchy have been recently presented in Ref. [7].

2 Composite Intermediary Models

The intermediary models employ two massive gauge singlet fields $S$ and $\overline{S}$ to communicate supersymmetry breaking to the visible sector. The role of these singlets is to generate an effective non-renormalizable operator connecting the fields in the supersymmetry breaking sector to the messenger quarks $Q, \overline{Q}$. Such an operator is suppressed by the mass of the singlet fields $M_S$. The effects of this operator will be identical to the coupling of the gauge singlet to the messenger quarks of the original model of DNNS [2] without requiring a messenger gauge group or complicated superpotential.

In the simplest version of the intermediary models, one has a pair of vector-like fields $V$ and $\overline{V}$ in the supersymmetry breaking sector which have a non-vanishing F-term. One of the singlets, $S$, couples to the vector-like flavor of the supersymmetry breaking sector $V, \overline{V}$, while the other singlet, $\overline{S}$, couples to the messenger quarks $Q, \overline{Q}$. The two sectors communicate only via the mass term for the two singlets, $M_S S \overline{S}$. To see the effect of this coupling we
examine the superpotential

\[ W = SVM + SQQ + MS\overline{S}. \]  

(1)

Integrating out the massive singlets \( \overline{S} \) produces an effective superpotential term

\[ W_{\text{eff}} = -\frac{(VM)(\overline{Q})}{MS}, \]  

(2)

below the mass scale \( MS \) which can mimic the effects of the coupling of the singlet to the messenger quarks in the original DNNS models [2], provided that the composite field \( V\overline{V} \) has the correct F-type expectation value. In order to obtain the correct gluino and squark masses one might also require an explicit mass term \( MQ\overline{Q} \) for the messenger quarks and a mass term \( MV\overline{V} \). We will see examples where the term \( MQ\overline{Q} \) is not necessary; the requirement on \( MV\overline{V} \) depends on whether it is necessary to lift dangerous flat directions. It is most convenient to assume it is present. Thus the full superpotential required for the intermediary models is of the form

\[ W = SVM + SQQ + MS\overline{S} + MQ\overline{Q} + MV\overline{V}. \]  

(3)

Once we have the superpotential of Eq. 3 with the right values of mass terms and the correct F-terms for the fields \( V\overline{V} \), the mechanism described above generates the effective \( -(VM)(\overline{Q})/MS \) coupling and supersymmetry breaking is mediated to the visible sector in the usual way. This mechanism is clearly rather simple and generic.

The question one has to address however is how the superpotential of Eq. 3 can be obtained without introducing artificially small mass terms (compared to the Planck-scale) while simultaneously including unsuppressed Yukawa couplings. We will present several models in which the effects of confining dynamics are exactly to produce the superpotential of Eq. 3. In these models the singlets \( S, \overline{S} \), the messenger quarks \( Q, \overline{Q} \) and the vector-like flavor \( V, \overline{V} \) are composites of an underlying strongly interacting gauge group. The effect of the confining dynamics will be to bind the “preon” fields into the composites \( S, \overline{S}, Q, \overline{Q}, V, \overline{V} \) and to generate an \( O(1) \) Yukawa coupling for these fields required for the superpotential of Eq. 3. The mass terms will be obtained from non-renormalizable preon operators that turn into mass terms for the composite fields after confinement [3]. Throughout this paper we will assume that the only scales present in the models are the dynamical scales of various confining groups and and the Planck-scale suppressing possible non-renormalizable operators.
2.1 The confining sector

We will first present a model which generates the $M_Q Q$ and $\lambda S Q Q$ terms, where $\lambda$ is of order one. These terms are a part of the superpotential in Eq. $\mathbb{3}$. It will become clear that the method is general and can be used with other gauge groups.

The idea is to use a confining theory whose global symmetries contain an $SU(5)$ subgroup and then gauge this subgroup which we identify with the ordinary $SU(5)$. The best known example of a confining supersymmetric theory is SUSY QCD with the number of flavors equal to the number of colors plus one. Here we want to gauge an $SU(5)$ symmetry, so we can take an $SU(4)$ theory with five flavors of fundamentals $q$ and antifundamentals $\bar{q}$. These fields transform as follows under the strong $SU(4)$ and global symmetry $SU(5) \times SU(5)$.

\[
\begin{array}{c|ccc|c}
SU(4) & SU(5) & SU(5) & SU(5)_D \\
q & 4 & 5 & 1 & 5 \\
\bar{q} & 5 & 5 & \end{array}
\]

The $SU(5)_D$ in the above table is not an additional symmetry, but the diagonal subgroup of $SU(5) \times SU(5)$ which we gauge.

This theory confines at low energies $\mathbb{3}$. The composite spectrum consists of the mesons $M \sim q\bar{q}$, the baryons $B \sim q^4$ and the antibaryons $\bar{B} \sim \bar{q}^4$, whose transformation properties are

\[
\begin{array}{c|ccc|c}
M \equiv \Sigma + S & SU(5) & SU(5) & SU(5)_D \\
& 5 & 5 & 24 + 1 \\
B \equiv Q & 5 & 1 & 5 \\
\bar{B} \equiv \bar{Q} & 1 & 5 & 5 \\
\end{array}
\]

Again, we denoted explicitly the gauged subgroup of $SU(5) \times SU(5)$ in the last column. The strong dynamics generates a superpotential

\[
W_{\text{conf}} = \frac{1}{\Lambda'} (\det M - BM\bar{B}),
\]

which needs to be expressed in terms of fields with definite $SU(5)_D$ quantum numbers. The composite meson field contains an adjoint and a singlet $S$ of the gauged $SU(5)$, while the baryon and the antibaryon are identified with the messenger quarks. Among other terms, the above superpotential contains the term $SQQ$, with a coefficient of order one for canonically normalized fields.
We also want to generate masses for the fields $Q$ and $\overline{Q}$. This can be achieved by adding a tree-level term in the microscopic theory of the form $W_{\text{tree}} = \frac{\lambda^6_{SU(4)}}{M_P}$. For normalized fields this translates to giving mass $M_Q = \frac{\Lambda^6_{SU(4)}}{M_P}$ for the composite fields $Q$, $\overline{Q}$. This way we have achieved generating a mass term for $Q$ and $\overline{Q}$ and a Yukawa term with the coupling of order one for the composite fields. This mechanism will be useful in a large class of our examples.

It is now obvious that the same mechanism can be used for fields transforming as fundamentals under an $SU(N)$ group. The confining interactions are then based on an $SU(N-1)$ gauge group. As before, the Yukawa coupling between the composite fundamentals and the composite singlet has a coefficient of order one. The mass term, on the other hand, is obtained from a tree-level superpotential term of the form $\frac{(q'q)(q'q)}{M_P^{2N-5}}$. Therefore the mass of the fundamental fields is $\frac{\Lambda_{SU(N)}^{2N-4}}{M_P^{2N-5}}$. In this scenario, the adjoint of the gauged subgroup has to be given a mass by adding a tree-level term which becomes a mass term after confinement.

We can now generate the remaining terms in the superpotential of Eq. 3 by the same method. Since in the first model presented below $V$ and $\overline{V}$ transform under $SU(3)$, we take a confining $SU(2)$ theory and obtain appropriate Yukawa coupling $SV\overline{V}$ and a mass $\frac{\lambda_{SU(2)}^2}{M_P}V\overline{V}$. We also can obtain a mass term for $S$ and $\overline{S}$. This requires a tree-level term $\frac{\lambda_{SU(2)}^2\lambda_{SU(3)}^2}{M_P}$, which gives $M_S = \frac{\Lambda_{SU(2)}^2\Lambda_{SU(3)}^2}{M_P}$. To summarize, assuming that the fields $V, \overline{V}, Q, \overline{Q}, S, \overline{S}$ are composites results in the superpotential of Eq. 3, where the trilinear Yukawas are order one due to the confining dynamics, while the masses are given by

$$M_Q = \frac{\Lambda^6_{SU(4)}}{M_P^6}, \quad M_V = \frac{\Lambda^2_{SU(2)}}{M_P^2}, \quad M_S = \frac{\Lambda_{SU(2)}^2\Lambda_{SU(4)}^2}{M_P^2}. \quad (4)$$

2.2 The 3-2 model as the supersymmetry breaking sector

In the first explicit example we use the “3-2 model” of Affleck, Dine and Seiberg [10] for the dynamical supersymmetry breaking sector, but add an additional flavor $(V + \overline{V})$ which transforms as $3 + \overline{3}$ under the $SU(3)$. This modified model clearly breaks supersymmetry when the extra flavor is massive. The superpotential terms $SV\overline{V} + M_VV\overline{V}$ are generated by a confining $SU(2)$ group with three flavors (six doublets) as described above, which will result in a mass term $M_V = \frac{\lambda_{SU(2)}^2}{M_P}$ and the order one Yukawa coupling $SV\overline{V}$. The superpotential terms for $\overline{S}$ and the messenger quarks are generated by a confining $SU(4)$ group as described above, resulting in $M_Q = \frac{\Lambda^6_{SU(4)}}{M_P^6}$. Finally, the mass for the singlets is
given by \( M_S = \frac{\Lambda_{SU(2)}\Lambda_{SU(3)}}{M_p} \). The fields of the supersymmetry breaking \( SU(3) \times SU(2) \) sector are

|       | \( SU(3) \) | \( SU(2) \) |
|-------|-------------|-------------|
| \( P \) | 3           | 2           |
| \( \bar{U} \) | 3           | 1           |
| \( \bar{D} \) | 3           | 1           |
| \( V \) | 3           | 1           |
| \( \bar{V} \) | 3           | 1           |
| \( L \) | 1           | 2           |

together with the tree-level superpotential for the 3-2 model \( W_{3-2} = \lambda P\bar{U}L \).

With this information at hand we can find out what the effective F-term multiplying the messenger quarks \( Q\bar{Q} \) is. For this we consider the effective theory below the mass scales \( M_S \) and \( M_V \). In order for this approximation to make sense we need to assume that \( M_V > \Lambda_{SU(3)} \), where \( \Lambda_{SU(3)} \) is the scale of the \( SU(3) \) group of the 3-2 model responsible for the breaking of supersymmetry. Below the scale \( M_S \), the singlets \( S, \bar{S} \) are integrated out resulting in the superpotential

\[
W = -(\frac{V\bar{V})(Q\bar{Q})}{M_s} + M_VV\bar{V} + M_QQ\bar{Q}.
\]

It is useful to analyze this theory by treating the coefficient of \( V\bar{V} \) as an effective (field dependent) mass for \( V\bar{V} \). This effective mass is thus \( M_V - \frac{Q\bar{Q}}{M_S} \). The effective scale of the \( SU(3) \) is (after integrating out \( V, \bar{V} \))

\[
\Lambda_{eff}^3 = (M_V - \frac{Q\bar{Q}}{M_S})\Lambda_{SU(3)}^3.
\]

Thus the effective superpotential is

\[
W = \frac{M_V(1 - \frac{Q\bar{Q}}{M_VM_S})\Lambda_{SU(3)}^6}{(P\bar{U})(P\bar{D})} + M_QQ\bar{Q} + \lambda P\bar{U}L. \tag{5}
\]

Provided that the VEV of the messenger quarks \( Q, \bar{Q} \) vanishes, the dynamical superpotential breaks supersymmetry, and an effective F-term is generated for the messenger quarks. One can estimate the magnitude of this effective F-term to be

\[
F_{eff} = \frac{F\Phi^2}{M_SM_V}, \tag{6}
\]
where $F$ is the magnitude of the F-terms in the supersymmetry breaking sector, while $\Phi$ is the value of the VEV’s in the supersymmetry breaking sector. In the case of the 3-2 model $\Phi \sim \Lambda_{SU(3)}$, $F \sim \Lambda_{SU(3)}^2$. 

The constraints on the parameters of all models of this sort are:

1. $\frac{F_{\text{eff}}}{M_Q} \simeq 10^{4.5}$ GeV to ensure the generation of the correct values of the gaugino and squark masses [3].

2. $F_{\text{max}} \leq 10^{18}$ GeV$^2$, in order to suppress gravity mediated soft supersymmetry breaking terms which would reintroduce the problem of flavor changing neutral currents.

3. $M_Q \geq \frac{F_{\text{eff}}}{M_Q} \simeq 10^{4.5}$ GeV to avoid the appearance of negative mass squared messenger quarks.

4. $M_V, M_S \geq \Lambda_{SU(3)}$ so that the presented effective theory approach is trustworthy.

These four constraints will be the same in all composite intermediary models; the only difference is in the expressions for $F_{\text{eff}}, M_Q, M_V$ and $M_S$ in terms of the dynamical scales of the different strongly interacting gauge groups.

In the case of the 3-2 model discussed above, one has to use Eq. 6 together with the expressions for $M_S, M_V$ and $M_Q$ of Eq. 4 obtained from confinement to get the bound on the “parameters” $\Lambda_{SU(2)}, \Lambda_{SU(4)}$ and $\Lambda_{SU(3)}$, where $\Lambda_{SU(2)}$ and $\Lambda_{SU(4)}$ are the scales of the confining groups and $\Lambda_{SU(3)}$ is the scale of the $SU(3)$ group of the 3-2 model. Putting these constraints together we obtain the bounds

\[
\Lambda_{SU(4)} \geq 10^{15.75} \text{ GeV } \\
10^{15.75} \text{ GeV } \geq \Lambda_{SU(2)} \geq 10^{12.15} \text{ GeV } \\
\Lambda_{SU(3)} \leq 10^{9} \text{ GeV},
\]

such that the inequality $\Lambda_{SU(4)}^7 \Lambda_{SU(2)}^3 \leq 10^{157.5}$ GeV$^{10}$ is satisfied. The bound on $\Lambda_{SU(2)}$ is relaxed to $10^{17}$ GeV $\geq \Lambda_{SU(2)} \geq 10^{12.15}$ GeV if we allow for a weaker $F_{\text{max}} \leq 10^{20}$ GeV$^2$ constraint instead of $F_{\text{max}} \leq 10^{18}$ GeV$^2$.

In order to improve the allowed range of parameters of the previous model one could consider a different dynamical supersymmetry breaking sector for the same kind of composite intermediary models. Clearly, if we use the same model of compositeness we will not change the bound $\Lambda_{SU(4)} \geq 10^{15.75}$ GeV since this is purely a consequence of the fact that the

\footnote{We have suppressed the dependence on the Yukawa coupling $\lambda$. We assume throughout this paper that Yukawa couplings are $O(1)$ and do not affect the estimates of mass scales much.}
messenger quarks transform as $5 + \overline{5}$ under the ordinary $SU(5)$ group. A different model of dynamical supersymmetry breaking could lead to a better model, provided that one can take a smaller representation than the $3 + \overline{3}$ of the 3-2 model, since in this case one would need to add a smaller confining group which would result in weaker bounds on the confining scale. However, the 3-2 model is one of the smallest models of dynamical supersymmetry breaking; thus much improvement cannot be achieved on the above bounds either.

Let us summarize the field content and the interactions of the model we introduced in this section. As a summary we repeat the complete structure of the model at high energies, above the compositeness scales at which the composite messenger quarks $Q$ and $\overline{Q}$ and the extra composite flavor $V$, $\overline{V}$ are generated. The field content is given in the table below.

|       | $SU(2)$ | $SU(3)$ | $SU(3)$ | $SU(2)$ | $SU(4)$ | $SU(5)$ | $SU(5)$ |
|-------|---------|---------|---------|---------|---------|---------|---------|
| $P$   | 1       | 3       | 1       | 2       | 1       | 1       | 1       |
| $\overline{U}$ | 1       | 3       | 1       | 1       | 1       | 1       | 1       |
| $D$   | 1       | 3       | 1       | 1       | 1       | 1       | 1       |
| $q$   | 2       | 3       | 1       | 1       | 1       | 1       | 1       |
| $\overline{q}$ | 2       | 1       | 3       | 1       | 1       | 1       | 1       |
| $L$   | 1       | 1       | 1       | 2       | 1       | 1       | 1       |
| $p$   | 1       | 1       | 1       | 1       | 4       | 5       | 1       |
| $\overline{p}$ | 1       | 1       | 1       | 1       | 4       | 1       | 5       |

where the first $SU(2)$ group is the confining $SU(2)$ producing the composite $V$ and $\overline{V}$ fields. The two $SU(3)$ factors are the global symmetries of the confining $SU(2)$ group, where the diagonal $SU(3)$ is gauged and is identified with the $SU(3)$ group of the supersymmetry breaking 3-2 model. The second $SU(2)$ factor is the $SU(2)$ of the 3-2 model, and the $SU(4)$ is the confining group producing the composite messengers. The two $SU(5)$’s are the global symmetries of the confining $SU(4)$ group, with the diagonal $SU(5)$ being identified with the ordinary $SU(5)$. The superpotential of this model is given at high energies by

$$W = P \overline{U} L + \frac{1}{M_P} p^\dagger \overline{p} + \frac{1}{M_P} q^2 \overline{q} + \frac{1}{M_P} q \overline{q} \overline{p} \overline{p},$$

which then results in gauge mediation of supersymmetry breaking as described above. We emphasize again that this model does not have gauge singlet chiral superfields. However, the model is not completely chiral. Therefore certain mass terms are forbidden by discrete or global symmetries.
Finally, we comment on the relation of the composite intermediary models and the models presented in Ref. [11], in which a nonrenormalizable operator coupling the messengers and the DSB sectors was assumed. In those models, it was found that one generally required a new mass scale to suppress these operators; \( M_P \) was too big. So an operator suppressed with an explicit new mass scale was necessary. It seemed difficult to realize this composite dynamics, in which both visible sector and DSB sector fields participated in common gauge dynamics without introducing problematic new operators. We therefore chose to realize this mass scale explicitly through the mass of the \( S, \overline{S} \) fields. The dynamics of the two sectors could then be separated and the operators which result are precisely those which are listed. However, the net result is similar in spirit; a composite operator which links the two sectors and obviates the need for a singlet F-term.

### 2.3 Models without explicit mass for the messengers

In this section we consider the possibility that there is no explicit mass term \( M_Q \) present in the superpotential for the messengers. To constrain the type of models we assume first that there is a non-renormalizable tree-level superpotential of the form \( W = \frac{\Phi^k}{M_P^k} \) where \( \Phi \)'s are fields from the supersymmetry breaking sector. We also assume that supersymmetry is broken after the generation of a dynamical superpotential of the form \( W_{dyn} = \frac{\Lambda^{p+3}}{\Phi^p} \). If we assume that there is no explicit messenger mass \( M_Q \) present, but the messenger mass comes purely from the expectation value of the effective operator of Eq. 2 then conditions 1-4 of the previous section imply \( k \leq 5 \), independent of the details of the supersymmetry breaking sector, \( i.e. \) independent of \( p \). To allow a larger \( \Lambda_{DSB} \) than \( 10^9 \) GeV one needs to consider theories with non-renormalizable operators. Since \( k \leq 5 \) we consider a theory which has \( k = 4 \), namely the \( SU(6) \times U(1) \) model of DNNS [2]. In this theory the dynamical

\[ \text{A similar analysis shows that the constraints 1-4 are independent of the supersymmetry breaking sector (} i.e. \text{ independent of } p \text{) even when an explicit mass term } M_Q \text{ for the messengers is included. In this case the conditions 1-4 yield a constraint } \left( \frac{M}{M_P} \right)^{k-3} M^2 M_Q < 10^{31.5} \text{ GeV}^3 \text{ where the coupling between the messengers } Q, \overline{Q} \text{ and the fields of the supersymmetry breaking sector } \phi \text{ has been assumed to be of the form } \frac{\phi^k \overline{Q}}{M_P}. \]
supersymmetry breaking sector is given by

\begin{align*}
| & S(6) & U(1) \\
A & 15 & 1 \\
\overline{F}^+ & 6 & -2 \\
\overline{F}^- & 6 & -2 \\
S^0 & 1 & 3 \\
S^+ & 1 & 3 \\
S^- & 1 & 3 \\
V & 6 & 0 \\
\overline{V} & 6 & 0 \\
\end{align*}

where we again added the extra vector-like flavor \( V, \overline{V} \) to the field content. The non-renormalizable superpotential required for dynamical supersymmetry breaking is

\[ W_{\text{tree}} = \frac{1}{M_P} A \overline{F}^+ \overline{F}^- S^0. \]

In order to generate the composite \( V, \overline{V} \) fields we need to introduce a confining \( SU(5) \) group with six flavors, and the resulting mass term for the \( V \) is given by

\[ M_V = \frac{\Lambda_{SU(5)}^8}{M_P^2}. \]

Similarly, the mass term for the singlets is given by

\[ M_S = \frac{\Lambda_{SU(4)} \Lambda_{SU(5)}}{M_P^2}. \]

The effective coupling between the messenger quarks and the fields \( \Phi \) of the supersymmetry breaking sector obtained after integrating out the fields \( S, \overline{S}, V, \overline{V} \) is

\[ \Phi^4 \frac{1}{M_P M_S M_V} Q \overline{Q}. \]

Since we have no explicit \( M_Q \) mass term for the messenger quarks present, the effective \( F/\Phi \) value just coincides with the original value of \( F/\Phi \) in the supersymmetry breaking sector, which fixes the scale of the supersymmetry breaking \( SU(6) \) group to be \( \Lambda_{SU(6)} = 10^{9.9} \) GeV. Since supersymmetry breaking is achieved through a non-renormalizable operator, the F-term in the supersymmetry breaking sector is \( F = \frac{\Lambda_{SU(6)}^{5/2}}{M_P^2} \). The condition \( F \leq 10^{18} \) GeV²
results in the requirement $\Lambda_{SU(6)} \leq 10^{10.8}$ GeV, which is clearly obeyed by the required value of $\Lambda_{SU(6)}$.

The constraint of having only positive mass-squared eigenvalues in the messenger sector is different here than in Section 2.2, since there is no explicit mass term $M_Q$ present. Now the condition $M_Q^2 > F_{\text{eff}}$ should be written as

$$\left( \frac{\Phi^4}{M_P M_S M_V} \right)^2 > \frac{F_\Phi \Phi^3}{M_P M_S M_V},$$

which results in the constraint

$$M_S M_V \leq 10^{22.5} \text{ GeV}^2.$$

Thus the allowed range for the mass parameters $M_S, M_V$ is

$$10^{10} \text{ GeV} \leq M_S, M_V \leq 10^{12} \text{ GeV},$$

such that $M_S M_V < 10^{22.5} \text{ GeV}^2$. The $10^{10} \text{ GeV} \leq M_V$ constraint together with Eq. 8 results in $\Lambda > 10^{17}$ GeV.

Another possible model with a non-renormalizable term necessary for supersymmetry breaking is the $Sp(4) \times U(1)$ model of Ref. [13]. The field content of the supersymmetry breaking sector is given by

|       | $Sp(4)$ | $U(1)$ |
|-------|---------|---------|
| $A$   | 5       | 2       |
| $q_1$ | 4       | $-3$    |
| $q_2$ | 4       | $-1$    |
| $s_1$ | 1       | 2       |
| $s_2$ | 1       | 4       |
| $V$   | 4       | 0       |
| $\nabla$ | 4   | 0       |

where the field $A$ is in the traceless antisymmetric representation of $Sp(4)$ and we have again introduced the vector-like flavor $V, \nabla$ into the supersymmetry breaking sector. The superpotential needed for supersymmetry breaking is

$$W = q_1 q_2 s_2 + \frac{1}{M_P} q_1 A q_2 s_1.$$

The vacuum structure of this theory can be analyzed by first assuming that the coefficient of the non-renormalizable operator in the tree-level superpotential is turned off. Then there is
a runaway direction, along which the $Sp(4)$ gauge group is completely broken. Turning the coefficient of the operator $\frac{1}{M_P} q_1 A q_2 s_1$ on yields a finite vacuum in which $Sp(4)$ is completely broken and the remaining uneaten singlets are massive. Carrying out this analysis with the requirement that the supersymmetry breaking F-term is less than $10^{18}$ GeV$^2$ gives the bound $\Lambda_{Sp(4)} < 10^{11.4}$ GeV. We again assume the same kind of compositeness as presented in Section 2.1, i.e. there is a confining $SU(3)$ group producing the composite fields $V, \nabla$ and $S$, while the confining $SU(4)$ group produces the messenger quarks $Q, \overline{Q}$ and the singlet $\overline{S}$. Since we assume no explicit $M_Q$ mass for the messengers the condition $(F/\phi)_{eff} = 10^{4.5}$ GeV fixes the $Sp(4)$ scale to be $\Lambda_{Sp(4)} = 10^{10.3}$ GeV, which is below the bound $10^{11.4}$ GeV obtained above. The remaining constraints yield the following requirements for the scales

$$\Lambda_{SU(3)} > 10^{16} \text{ GeV, } 10^{14.5} \text{ GeV} > \Lambda_{SU(4)} > 10^{10.3} \text{ GeV.}$$

### 2.4 General considerations on variations of the composite intermediary models

First, one should ask the question of why we need the fields $V, \nabla$ at all. In this scenario one of the singlets $S$ would be elementary and directly coupled to the supersymmetry breaking sector through a non-renormalizable operator, while the other singlet $\overline{S}$ and the messenger quarks $Q, \overline{Q}$ are still composite. One model would be for example to consider the 3-2 model without an extra $SU(3)$ flavor, but with the additional coupling

$$W = \frac{1}{M_P} SPUL.$$

However the condition $F < 10^{18}$ GeV$^2$ results in $\Lambda_{SU(4)} \leq 10^{14}$ GeV, which is incompatible with the condition $\Lambda \geq 10^{15.75}$ GeV which comes from $M_S > \Lambda_{SU(3)}$ and so this model is excluded. Since the most restrictive condition $\Lambda \geq 10^{15.75}$ GeV is a consequence of our choice of model of compositeness one could ask whether it is possible to make this model work by choosing a different scheme for compositeness. For example one could instead of gauging the diagonal $SU(5)$ subgroup of the $SU(5) \times SU(5)$ global symmetry of a confining $SU(4)$ theory with five flavors we just gauge one of the global $SU(5)$’s and identify that with the
ordinary $SU(5)$. In this case the field content of the confining module is

\[
\begin{array}{c|ccc}
   & SU(4) & SU(5) & SU(5) \\
\hline
  q & 4 & 5 & 1 \\
  \overline{q} & 4 & 1 & 5 \\
p_i & 1 & 5 & 1 \\
\end{array}
\]

, \ i = 1, \ldots, 4.

The fields $p_i$ are included to cancel the $SU(5)$ anomalies. The messenger quark could be identified with one of the meson fields $q \overline{q}$ and the baryon $q^4$, while the singlet could be identified with a component of the antibaryon $\overline{q}^4$. In this case the expressions for the composite masses for the singlets and the messenger quarks are modified to be

\[
M_S = \frac{\Lambda_{SU(4)}^3}{M_P^2}, \quad M_Q = \frac{\Lambda_{SU(4)}^4}{M_P^3}.
\]

Using these expressions together with the constraints 1-4 of Section 2.2 we get that the allowed range for $\Lambda_{SU(4)}$ is

\[
10^{14.63} \text{ GeV} \leq \Lambda_{SU(4)} \leq 10^{14.78} \text{ GeV}.
\]

The upper bound is extended to $10^{15.3}$ GeV if we allow for F-terms up to $10^{20}$ GeV$^2$. Thus we can see that there is no fundamental need for the extra vector-like flavor $V + \overline{V}$ in the dynamical supersymmetry breaking sector, but models with vector fields have a larger allowed parameter range.

Next, we address the question of why we needed to rely in our composite models on the dynamically generated confining superpotential terms. As an alternative approach one could just add the appropriate tree-level superpotential term for the preon fields by hand. In this case however, the Yukawa coupling terms would not be of order one, but would be suppressed by powers of the ratios of the confinement scales to the Planck scale. For example in the model with mass terms present for $V, \overline{V}, S, \overline{S}, Q, \overline{Q}$ the effective $F/\phi$ is just

\[
\lambda \frac{F^2}{M_S M_V M_Q} \sim 10^{4.5} \text{ GeV}, \tag{9}
\]

where $\lambda = \frac{M_Q M_S \Lambda_S^{p-1}}{M_P^2 M_P^3 M_P^{p-1}}$, where $\Lambda_S$ is the compositeness scale of the group which binds $p$ preons into $S$. Thus the constraint of Eq. (9) is just

\[
\frac{F^2 M_P^{p-3}}{M_V \Lambda_S^{p-1}} = 10^{4.5} \text{ GeV},
\]

13
which together with $F \leq 10^{18}$ GeV$^2$ results in unacceptably small values of the product $M_V\Lambda_S^{p-1}$. Thus we conclude that one really needs the order one Yukawa couplings resulting from confinement in order to obtain a viable spectrum using the intermediary model.

Finally, we consider the possibility of using the dynamical supersymmetry breaking model of Ref. [12] for the intermediary model. Recall that this model has an $SU(2)$ gauge symmetry with four doublets $q_i$ and six singlets $S^{ij}$ with a tree-level superpotential $S^{ij}q_iq_j$. Since this model has naturally singlets and vector-like flavors $q_i$ in it one would think that this model could be used as the dynamical supersymmetry breaking sector in an intermediary model. This is however not true. The reason is that the fields $q_i$ don’t have an F-term. This can be seen by writing down the scalar potential

$$V = |\text{Pf} M - \Lambda_{SU(2)}^4|^2 + |M_{ij}|^2 + |S^{ij}|^2 + \lambda \epsilon^{ijkl} M_{kl}|^2,$$

where $\lambda$ is the Lagrange multiplier enforcing the quantum modified constraint and $M_{ij}$ is the meson field $q_iq_j$. The point is that the F-terms corresponding to $M_{ij}$ can always be satisfied by an appropriate choice of the singlet VEV’s and thus the mesons will not obtain F-term in this model of supersymmetry breaking, therefore one can not use it directly in an intermediary model.

To summarize this section, we have shown two explicit composite intermediary models which provide a viable sparticle spectrum similar to the original DNNS spectrum. In one case, we used the 3-2 model as the dynamical supersymmetry breaking sector, and in the other case, we used the non-renormalizable $SU(6) \times U(1)$ model of DNNS for supersymmetry breaking. Since all fields of these models are composites, there are no fundamental singlets present in these theories. The Yukawa couplings are of order one due to confinement and the mass terms arise from tree-level superpotential terms turning into masses after confinement. We have calculated the viable parameter range for the confinement scales of these models.

### 3 Mediator Models

We begin by reviewing the mediator models. There is a weakly gauged messenger group $G_m$ acting on fields in the dynamical supersymmetry breaking sector. This gauge group may or may not be broken, but a supersymmetry breaking gaugino mass, $M_m$, for the gauge bosons of this group is essential. We assume this is achieved through a dynamical
messenger sector, such as those discussed in Refs. \cite{5,6}. However, in the models of Ref. \cite{5,6} the weakly gauged group was $SU(5)$, so that supersymmetry breaking was communicated directly, whereas we instead gauge a messenger gauge group $G_m$. For definiteness, we will take the messenger gauge group $G_m$ to be $SU(3)$. This additional step in communicating supersymmetry breaking solves the negative mass squared problem for squarks and sleptons which was encountered in those models \cite{5,6}.

Second, there are “mediator fields”, which we call $T$ and $\tilde{T}$, which transform both under the messenger gauge group, $G_m$, and the standard model gauge group, $G_{SM}$ (or an extension thereof) and therefore “talk” to both the dynamical supersymmetry breaking and visible sectors. Third, there is a supersymmetric mass term $M_T T \tilde{T}$. As discussed in Ref. \cite{1}, $M_T$ is constrained to lie between $M_m$, the messenger gaugino mass, and the mass of the heaviest of the hidden sector scalars, $M_{DSB}$.

In these models, there is no need to generate an $F$ term for a singlet in order to generate the gaugino mass. Instead, it occurs from the three-loop diagram given in Figure 1.

The scalar squared mass occurs at four-loops. The leading contribution can readily be obtained when there is a large separation of mass scales as follows. As computed in Ref. \cite{5,6}, there is a negative logarithmically enhanced contribution to $ST r M^2$ arising at two-loops. This feeds into the squark and slepton mass squared with an additional two-loops. The nice
feature is the net contribution to the squark and slepton mass squared is positive if the initial $STrM^2$ in the messenger sector was positive, as it often is. This means that the model is phenomenologically viable, although it can be more fine-tuned than the more conventional gauge-mediated mass predictions. In particular, there is a relatively large prediction for the scalar to gaugino mass ratio (of like charges) which is

$$\frac{m_i}{m_{i,1/2}} \approx \alpha_m \frac{50}{N_f} \left( \frac{k_i}{k_m} \right)^{1/2} \left[ \beta \log \left( \frac{M}{k_i} \right)^2 - N_f \right]^{1/2} > 1 \quad (10)$$

where $N_f$ is the number of flavors charged under $G_m$ and can be order 10, $\beta$ is a number of $O(1)$, $M$ is the biggest mass scale in the supersymmetry breaking sector, $\alpha_m = g_m^2/4\pi$, where $g_m$ is the messenger gauge coupling. Therefore the ratio, although large, might be viable.

One interesting feature of these models is that they have the correct global minimum. That is, the minimum is the color and charge-preserving minimum in which supersymmetry is broken. This model is the only known example we know with this property (in which additional fields are not added ad hoc in order to guarantee this property).

We now address the question of how to realize these models in such a way that the mass term is not introduced by hand. As we will see, this is fairly straightforward.

### 3.1 A Composite Mediator Model

In the explicit model we present the mediator fields $T, \overline{T}$ are composites of a confining $Sp(4) \times Sp(4)'$ gauge group. Each of the $Sp(4)$ groups has eight fields transforming as a fundamental representation 4. A gauged $SU(5)$ subgroup of the global $SU(8)$ symmetry is identified with the ordinary $SU(5)$ group and a gauged $SU(3)$ group is identified with the messenger $SU(3)_m$ gauge group. The fields transforming under the confining $Sp(4)$ groups

\footnote{We thank Asad Naqvi for sharing his result.}
are summarized in the table below.

|        | $SU(5)_{SM}$ | $SU(3)_m$ | $Sp(4)$ | $Sp(4)'$ |
|--------|--------------|-----------|---------|----------|
| $T_5$  | 5            | 1         | 4       | 1        |
| $T_3$  | 1            | 3         | 4       | 1        |
| $\bar{T}_5$ | 5       | 1         | 1       | 4        |
| $\bar{T}_3$ | 1       | 3         | 1       | 4        |

The bound state spectrum of this theory is given by

$$W_{tree} = \lambda_1 T_5^2 \bar{T}_5^2 + \lambda_2 T_5 T_3 \bar{T}_5 T_3 + \lambda_3 T_3^2 \bar{T}_3^2$$

is present in the theory for the preon fields. After confinement, the tree-level Yukawa coupling are turned into mass terms and a confining superpotential $PfM$ is generated resulting in the low energy superpotential

$$W = PfM + PfM' + \lambda_1 (T_5^2)(\bar{T}_5^2) + \lambda_2 T \bar{T} + \lambda_3 (T_3^2)(\bar{T}_3^2).$$

Thus, the $T \bar{T}$ mass term is given by $M_T = \Lambda_{Sp(4)} \Lambda_{Sp(4)'}/M_P$. Assuming the two $Sp(4)$ groups have scales of the same order, and assuming that the masses in the dynamical supersymmetry breaking sector lie between $10^4$ GeV and $10^9$ GeV, we get a bound on the confining scale of $10^{11}$ GeV < $\Lambda_{Sp(4)} < 10^{13.5}$ GeV. This model, besides being quite simple also has the advantage that the F-term in the supersymmetry breaking sector is not required to be close to $10^9$ GeV. Therefore this model can also satisfy the constraints of Ref. [7] coming from nucleosynthesis.

\[8\]Note that similar product group structures involving symplectic groups arise naturally in the orbifold construction of Ref. [14].
4 Conclusions

We have presented explicit realizations of the “Intermediary” and “Mediator” models of gauge mediated supersymmetry breaking presented in Ref. [1]. Both of these models have non-vanishing tree-level mass terms as essential ingredients. In the models we presented these mass terms arise due to confining dynamics. The mass scales are not put in by hand but rather determined as a function of the dynamical scales of the confining groups and the Planck scale. These masses always arise after tree-level Yukawa couplings (sometimes non-renormalizable operators suppressed by the Planck scale) turn into mass terms for the confined low-energy degrees of freedom. The intermediary models also require order one Yukawa couplings of the confined fields. In our models, these are generated dynamically via the confining superpotentials of Ref. [8].

The intermediary models have two singlets but no messenger gauge groups. The interactions of these singlets with the messenger quarks and with the fields in the dynamical supersymmetry breaking sector generate an effective operator that has similar effects as the singlet coupling to the messenger quarks in the DNNS models. In the composite versions of this model both singlets and the messenger quarks are composite and thus there is no elementary singlet present at high energies. We have presented several models of this sort and given the allowed range of the parameters of the theory. The mediator models employ two fields $T, \overline{T}$ which carry both the charges of the messenger gauge group and of ordinary $SU(5)$. However, an explicit mass term for these fields together with masses for the messenger gluinos are crucial for this model to work. We presented an example in which the mass term of the mediators arises via confinement where tree-level Yukawa terms turn into masses for the composite mediator fields.

The models presented here, while being explicit realizations of the scenarios presented in Ref. [1], are new complete examples of gauge mediation of dynamical supersymmetry breaking. The intermediary models don’t have a messenger gauge group and the composite versions of these models do not contain elementary gauge singlet fields either, thus presenting a simplification compared to the original model of Ref. [2]. The mediator models do have a messenger gauge group, but the additional structure required for gauge mediation is simpler than in the intermediary models. The explicit models presented in this paper represent viable alternatives to the conventional models of gauge mediated supersymmetry breaking.
References

[1] L. Randall, *Nucl. Phys.* **B495**, 37 (1997), hep-ph/9612426.

[2] M. Dine, A.E. Nelson, and Y. Shirman, *Phys. Rev.* **D51**, 1362 (1995), hep-ph/9408384; M. Dine, A. Nelson, Y. Nir, and Y. Shirman, *Phys. Rev.* **D53**, 2658 (1996), hep-ph/9507378.

[3] K. R. Dienes, C. Kolda, J. March-Russell, *Nucl. Phys.* **B492**, 104 (1997), hep-ph/9610479.

[4] I. Dasgupta, B.A. Dobrescu, and L. Randall, *Nucl. Phys.* **B483**, 95 (1997), hep-ph/9507487; N. Arkani-Hamed, C. D. Carone, L. J. Hall, H. Murayama, *Phys. Rev.* **D54**, 7032 (1996), hep-ph/9607298.

[5] E. Poppitz and S. Trivedi, *Phys. Rev.* **D55**, 5508 (1997), hep-ph/9609529; *Phys. Lett.* **401B**, 38 (1997), hep-ph/9703246.

[6] N. Arkani-Hamed, J. March-Russell, and H. Murayama, hep-ph/9701286.

[7] H. Murayama, hep-ph/9705271; S. Dimopoulos, G. Dvali, R. Rattazzi, and G. Giudice, hep-ph/9705307.

[8] N. Seiberg, *Phys. Rev.* **D49**, 6857 (1994), hep-th/9402044; C. Csáki, M. Schmaltz and W. Skiba, *Phys. Rev. Lett.* **78**, 799 (1997), hep-th/9610139; *Phys. Rev.* **D55**, 7840 (1997), hep-th/9612207.

[9] E. Poppitz, Y. Shadmi, and S. Trivedi, *Phys. Lett.* **388B**, 561 (1996), hep-th/9606184; C. Csáki, L. Randall, and W. Skiba, *Nucl. Phys.* **B479**, 65 (1996), hep-th/9605108; C. Csáki, L. Randall, W. Skiba, and R. Leigh, *Phys. Lett.* **387B**, 791 (1996), hep-th/9607021; R. Leigh, L. Randall, and R. Rattazzi, hep-ph/9704246.

[10] I. Affleck, M. Dine, and N. Seiberg, *Nucl. Phys.* **B256**, 557 (1985).

[11] Y. Shadmi, hep-ph/9703312.

[12] K.-I. Izawa and T. Yanagida, *Prog. Theor. Phys.* **95**, 829 (1996), hep-ph/9602180; K. Intriligator and S. Thomas, *Nucl. Phys.* **B473**, 121 (1996), hep-th/9603158.
[13] C. Csáki, W. Skiba and M. Schmaltz, *Nucl. Phys.* **B487**, 128 (1997), hep-th/9607210

[14] K. Intriligator, *Nucl. Phys.* **B496**, 177 (1997), hep-th/9702038