A Model of Direct Gauge Mediation

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Abstract

We present a simple model of gauge mediation (GM) which does not have a messenger sector or gauge singlet fields. The standard model gauge groups couple directly to the sector which breaks supersymmetry dynamically. This is the first phenomenologically viable example of this type in the literature. Despite the direct coupling, the model can preserve perturbative gauge unification. This is achieved by the inverted hierarchy mechanism which generates a large scalar expectation value compared to the size of supersymmetry breaking. There is no dangerous negative contribution to the squark, slepton masses due to two-loop renormalization group equation. The potentially non-universal supergravity contribution to the scalar masses can be suppressed enough to maintain the virtue of the gauge mediation. The model is completely chiral, and one does not need to forbid mass terms for the messenger fields by hand. Beyond the simplicity of the model, it possesses cosmologically desirable features compared to the original models of GM: an improved gravitino and string moduli cosmology. The Polonyi problem is back unlike in the original GM models, but is still much less serious than in hidden sector models.

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Recently, gauge mediation of supersymmetry breaking [1] has attracted interests. The primary motivation of this scheme is to guarantee degeneracy between the masses of sfermions belonging to different generations, thereby solving the supersymmetric flavor-changing problem. The pioneering works of Dine, Nelson and collaborators [2, 3, 4] demonstrated that such a scheme is phenomenologically viable, and they presented explicit models which realize the gauge mediation mechanism. Their models, however, require three relatively decoupled sectors “insulated” from each other, namely the standard model sector, the sector of dynamical supersymmetry breaking (DSB), and the so-called messenger sector. It is an important question to ask whether this structure is inevitable for successful gauge mediation or whether it simply provides an existence proof, calling for further simplification.

Some progress has been made in this direction [5, 6]. Poppitz and Trivedi [5] demonstrated that one can couple the DSB sector and the standard model gauge groups directly without spoiling perturbative gauge unification. However, their model suffered from two problems. One is that the supersymmetry breaking scale is so high that the supergravity contribution to squark and slepton masses dominate over the contributions from the gauge mediation. Therefore, the degeneracy among sfermions is not an automatic consequence of the model. Secondly, there are fields which are charged under the standard model gauge groups below $10^5$ GeV, whose scalar components have supersymmetry breaking soft scalar masses of (a few $\times 10^4$ GeV)$^2$. They contribute to the two-loop renormalization group equations (RGE) of squark and slepton masses, driving them negative at low energies [6, 8]. The problem with supergravity contributions was surmounted in the model constructed by Arkani-Hamed, March-Russell and the author [6], but the problem with light charged fields with large soft scalar masses afflicts both classes of models. Despite the problem, these works left hope that the DSB sector may not need to be as “insulated” as in the original models. The basic ingredient here is that the mass of the messengers $M_Q^2$ can be much larger than their supersymmetry breaking bilinear mass $B_Q M_Q$ due to the dynamics of the models; then the ratio $(B_Q M_Q)/M_Q$ can be kept around $10^4$ GeV to generate desired magnitude of gaugino masses while $M_Q$ is close to the unification scale, which makes it easy to maintain the perturbative gauge unification.

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See [7] for other attempts to simplify the structure of gauge mediation models along different lines.
In this letter, we present a simple model of gauge mediation without a messenger sector or gauge singlet fields. The standard model gauge groups are an important part of the dynamics. The model overcomes the difficulties mentioned above: it preserves the perturbative gauge unification while suppressing the supergravity contributions enough to avoid flavor-changing processes. There are light multiplets in the DSB sector which are charged under the SM gauge group but they do not have large soft masses. Therefore there is no dangerous negative contribution to the squark, slepton masses due to two-loop RGE, and the model is phenomenologically viable. An aesthetically appealing feature of the model is that it is completely chiral, and one does not need to forbid mass terms for the messenger fields by hand unlike in the earlier models. Since there is no messenger sector in this model, the vector-like messengers are produced as a consequence of dynamics (this is a feature also shared by the models in [5, 6]). Furthermore, it possesses cosmologically desirable features compared to the original models of gauge mediation.

Our model is based on the vector-like $SP(N)$ models by Izawa, Yanagida [9] and by Intriligator, Thomas [10]. We take $SP(4)$ model with 5 flavors, i.e. 10 fundamentals, which we denote by $Q^i$ ($i = 1, \cdots, 10$). In addition, there are $SP(4)$ singlet fields $S_{ij}$ which couple to the $Q^i$’s in the superpotential,

$$W = \lambda S_{ij} Q^i Q^j.$$  \hfill (1)

This coupling lifts all flat directions: $Q^i Q^j = 0$, while the quantum modified constraint requires $\text{Pf}(Q^i Q^j) = \Lambda^{10}$. The contradiction between two conditions implies that supersymmetry is broken.

Into the SU(10) global symmetry of the model, we embed SU(5)$_L \times$SU(5)$_R$ gauge group.‡ The particle content of the model is shown in Table 1. The only difference from the original model is that we need to add $\phi^a$ and $\bar{\phi}^a$ ($a = 1, 2$) to cancel the SU(5) anomalies. We assume that both SU(5) are weaker than SP(4) and treat them perturbatively.§ The superpotential is nothing but

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‡This is the gauge group above the GUT-scale. Below the GUT scale, one of the SU(5) groups must be broken down to the SU(3)$\times$SU(2)$\times$U(1) gauge group to preserve the observed gauge unification. We use SU(5) language to both SU(5) factors to keep track of the quantum numbers easily.

§It would be an interesting exercise to see how the description of the dynamics changes when one or both of the SU(5) groups are stronger than the SP(4) group.
Table 1: The particle content of our model under $\text{SP}(4) \times \text{SU}(5)_L \times \text{SU}(5)_R$ gauge group and a global $\text{SU}(2)$ symmetry. The symbol $\cdot$ refers to singlets.

|        | $\text{SP}(4)$ | $\text{SU}(5)_L$ | $\text{SU}(5)_R$ | $\text{SU}(2)$ |
|--------|-----------------|------------------|------------------|---------------|
| $Q$    | $\Box$          | $\Box$           | $\cdot$          | $\cdot$       |
| $\bar{Q}$ | $\Box$         | $\cdot$          | $\Box$           | $\cdot$       |
| $\Sigma$ | $\cdot$     | $\Box$           | $\cdot$          | $\cdot$       |
| $S$    | $\cdot$         | $\Box$           | $\cdot$          | $\cdot$       |
| $\bar{S}$ | $\cdot$        | $\cdot$          | $\Box$           | $\cdot$       |
| $\phi^a$ | $\cdot$   | $\Box$           | $\cdot$          | $\Box$       |
| $\bar{\phi}^a$ | $\cdot$ | $\cdot$          | $\Box$           | $\Box$       |

what is obtained by decomposing the original one (1) into $\text{SU}(5)_L \times \text{SU}(5)_R$ multiplets,

$$W = \frac{1}{2} g S^{kl} Q_k Q_l + \frac{1}{2} g \bar{S}_{\kappa \lambda} \bar{Q}^\kappa \bar{Q}^\lambda + \lambda \Sigma^k_{\kappa} \bar{Q}^\kappa Q_k. \quad (2)$$

Latin and Greek letters refer to different $\text{SU}(5)$ indices. This is the most general renormalizable superpotential consistent with gauge invariance and global $\text{SU}(2)$ symmetry. Just as in the original model, the quantum modified moduli space requires (at least some of) the meson operators $QQ$, $\bar{Q}Q$ or $\bar{QQ}$ not to vanish, while the superpotential requires all meson operators to vanish. Therefore supersymmetry is broken. Note that this model is completely chiral, i.e. none of the fields in the model can have mass terms in the superpotential. In particular, the absence of terms in the superpotential such as $S$, $S^2$ or $S^3$ etc. had to be imposed by hand in the vector-like $\text{SP}(N)$ models in order to break supersymmetry, while they are automatically forbidden by the gauge invariance of the model here. There are no gauge singlet fields in this model which is also aesthetically appealing.

We are primarily interested in the situation where $\Sigma$ acquires a large expectation value. Existence of a potential minimum with large expectation

\footnote{This SU(2) can also be promoted to a local one.}
value will be shown later. Consider the following configuration,

\[
\Sigma^k = \frac{v}{\sqrt{5}} \delta^k_\kappa = \frac{1}{\sqrt{5}} \begin{pmatrix} v & v \\ v & v \\ v & v \end{pmatrix}
\]

(3)

with \( v \gg \Lambda \). One can view this direction as a gauge invariant polynomial \( \det \Sigma \). This configuration breaks SU(5)_L \times SU(5)_R gauge group down to their diagonal SU(5) subgroup. We identify the unbroken diagonal SU(5) as “our” gauge group into which the standard model gauge groups are embedded. Along this direction (3), the \( \Sigma \) field has an exactly flat potential at the tree-level. This can be seen by the following analysis. First of all, all \( Q \)'s become massive with mass \( \lambda v / \sqrt{5} \) and can be integrated out from the theory. All components of \( \Sigma \) except the \( v \) direction are eaten by the broken SU(5) generators. Therefore the low energy theory is a pure SP(4) gauge theory with a singlet chiral superfield \( v \) together with fields charged under the unbroken diagonal SU(5) \( S, \bar{S}, \phi, \bar{\phi} \). The low-energy SP(4) gauge group develops a gaugino condensate, which results in an effective superpotential

\[
W_{\text{eff}} = \Lambda^2 (\text{Pf} M_{ij})^{1/5},
\]

(4)

where the (anti-symmetric) mass matrix \( M_{ij} \) of the quarks \( Q^i \) are given by the expectation values of \( \lambda \Sigma, gS \) and \( \bar{g} \bar{S} \),

\[
M_{ij} = \begin{pmatrix} gS & \lambda \Sigma \\ -\lambda' \Sigma & \bar{g} \bar{S} \end{pmatrix}.
\]

(5)

Along the direction of our interest Eq. (3), the effective superpotential (4) is

\[
W_{\text{eff}} = \frac{1}{\sqrt{5}} \lambda \Lambda^2 v
\]

(6)

and the \( v \) derivative does not vanish; hence supersymmetry is broken. With a canonical (tree-level) Kähler potential for the \( v \) direction, the potential is

\[\text{Here, } v \text{ is a chiral superfield, but we refer to its scalar expectation value by the same symbol.}\]
exactly flat. On the other hand, the $v$ direction acquires an $F$-component, $F_v = \lambda \Lambda^2 / \sqrt{5}$. Therefore, the $Q$ fields have a large mass $M_Q = \lambda v / \sqrt{5}$ with a bilinear supersymmetry breaking mass term $B_Q M_Q = \lambda F_v / \sqrt{5}$, and hence act as messengers of supersymmetry breaking. The ratio

$$B_Q = \frac{F_v}{v} = \frac{\lambda \Lambda^2}{\sqrt{5} v}$$

(7)

determines the size of sfermion, gaugino masses induced by standard model gauge interactions from the loops of $Q$'s (gauge mediation). The generated gaugino and scalar masses is $N_m = 8$ times larger than the case of the minimal messenger sector $5 + 5^*$ in [2, 3]. We would like this ratio to be around $B_Q \sim 10^4$ GeV to generate supersymmetry breaking terms of the desired magnitude. Recall that $Q, \bar{Q}$ are chiral under the original gauge group; they become vector-like messengers only after the symmetry breaking $SU(5)_L \times SU(5)_R \rightarrow SU(5)$. Since the vector-like messengers originate dynamically in this model, there is no need for a separate and “insulated” messenger sector.

The rest of the mass spectrum of the model is the following at the renormalizable level (we will discuss additional mass terms generated by Planck-scale operators later). The $S$ and $\bar{S}$ fields become vector-like $(\bf{10} + 10^*)$ under the unbroken diagonal $SU(5)$ and acquire an invariant mass. By expanding the effective superpotential Eq. (4.5) up to the second order in $S, \bar{S}$, one finds that they acquire a mass of $g\bar{g}\Lambda^2 / (\sqrt{5} \lambda v) = (g\bar{g} / \lambda^2) B_Q \sim 10^4$ GeV. $\phi^a$ and $\bar{\phi}^a$ remain massless at this point. The heavy gauge multiplet has a mass of order $v$, and contributes to the gaugino masses as well [11]. Two-loop contribution of heavy gauge multiplet to the scalar masses has not been calculated to the best of our knowledge.

So far $v$ acquired the $F$-component, but its $A$-component is undetermined because its potential is exactly flat at the tree-level. In fact, its potential is lifted by the inverted hierarchy mechanism [4] as follows. The Kähler potential for the $v$ direction can be calculated by perturbation theory since it does not participate in strong $SP(4)$ dynamics, as long as $v \gg \Lambda$. Its wave function receives renormalization due to the superpotential coupling $\lambda$ and the gauge couplings of $SU(5)_L \times SU(5)_R$. From the effective Lagrangian for

**If one follows the argument in [12], one would conclude that there is also a contribution from $SP(4)$ gauge coupling to the effective potential $V_{eff}$. This is not true because the wave
the $v$ field
\[ \mathcal{L} = \int d^4 \theta Z(v) v^* v + \int d^2 \theta \lambda \Lambda^2 v / \sqrt{5}, \]  
the effective potential for $v$ is obtained as
\[ V_{\text{eff}} = \frac{1}{5 Z(v)} |\lambda \Lambda^2|^2. \]  
The point is that the Yukawa contribution and the gauge contribution in $Z(v)$ have the opposite sign:
\[ Z(v) = 1 + \frac{12}{5} g_L^2 + g_R^2 \ln \frac{v^2}{M^2} - 40 \frac{(\lambda/\sqrt{5})^2}{16 \pi^2} \ln \frac{v^2}{M^2}. \]  
Here, $g_L$ ($g_R$) is the gauge coupling constant of $SU(5)_L$ ($SU(5)_R$) group and $M$ the ultraviolet cutoff. This is the situation which makes the inverted hierarchy possible. The gauge couplings make $Z(v)$ larger at higher energies which makes the potential decrease, while the Yukawa coupling makes the potential increase. It can well happen that the gauge piece dominates at lower energies which makes the potential to decrease, but at some point the Yukawa piece wins over the gauge piece due to the renormalization group running. The potential is minimized at the energy scale where this turnover occurs. Because the running is logarithmic in scales, the potential can

function of $\Sigma$ receives renormalization from $S\!P(4)$ interaction only at the two-loop level. Since the effective superpotential Eq. (6) is exact, there is no other effect which modifies the effective potential other than the wave function renormalization $Z(v)$. The fact that there is no dependence on $S\!P(4)$ gauge coupling or on the wave function renormalization of $Q, \bar{Q}$ is also consistent with the correct treatment of Wilsonian renormalization group running.\[\dag\]

Consider, for instance, the case where all standard model fields couple to $SU(5)_L$ which is broken to $SU(3) \times SU(2) \times U(1)$ at the GUT-scale and $SU(5)_R$ remains unbroken between $v$ and the GUT-scale. Then all $g_L$'s grow slowly as a function of energy but remain perturbative, while $g_R$ is asymptotically free. From the low-energy data, $\alpha^{-1}_3(v) \sim 8$–16 depending on the value of $v$ as will be determined later. Therefore the combination $(1/g_L^2 + 1/g_R^2)(v)$ is constrained by the low-energy data, while we are free to choose $g_R^2 \gtrsim 2g_L^2$. Then $g_L^2 + g_R^2$ decreases for higher energies. On the other hand, a Yukawa coupling can easily increase as a function of energy as known to be the case in the top Yukawa coupling. It is quite natural that a turnover occurs at some energy scale with $\lambda \sim O(1)$. For different embeddings of the standard model generations, the details can be somewhat different. We do not commit ourselves to one particular embedding in this letter and do not go into more quantitative analysis of the running coupling constants.
develop an exponentially large expectation value for \( v \) compared to the size of supersymmetry breaking given by \( \Lambda \).

Even though the minimum of the effective potential discussed above is a well-defined consistent local minimum, we cannot exclude a possibility that there is another minimum close to the origin. For instance, the point \( \Sigma = S = \bar{S} = 0 \) has an enhanced symmetry (global U(1), which is an accidental symmetry of our model) and is an extremum of the potential. Whether it is a minimum or a maximum cannot be answered because it is in a strongly coupled regime and the theory is not calculable. We note, however, that the minimum close to the origin (if any) is so far away from the other minimum \( v \gg \Lambda \) such that the vacuum tunneling is presumably highly suppressed, and hence we can live on the minimum discussed.

We have shown that our model which is completely chiral generates vector-like messenger fields due to the dynamical breakdown of the gauge symmetries. Because of the inverted hierarchy, \( v \) can be naturally much larger than \( \Lambda \), which makes the messenger fields very heavy. Now we study the values of \( v \) which give us a viable phenomenology. First requirement is that the potentially non-universal supergravity contribution to the squark, slepton masses are suppressed relative to the gauge-mediated contribution. The typical size of the supergravity contribution is characterized by the gravitino mass:

\[
m_{3/2} = |F_v|/(\sqrt{3}M_*) = B_Qv/\sqrt{3}M_*.
\] (11)

Here, \( M_* = M_{\text{Planck}}/\sqrt{8\pi} \approx 2 \times 10^{18} \) GeV is the reduced Planck mass. We require that the gravitino mass is less than 10% of the gluino mass \( M_3 = N_m(\alpha_s(m_Z)/4\pi)B_Q \) such that the RGE-induced squark mass squareds have degeneracy at the 1\% level. We find

\[
v \lesssim 0.1N_m(\alpha_s(m_Z)/4\pi)(\sqrt{3}M_*) \sim 3 \times 10^{16} \text{ GeV}.
\] (12)

It can be as large as the unification scale\( ^\dagger \)

So far the \( \phi^a, \bar{\phi}^a \) fields are massless, which is phenomenologically unacceptable. However, they, together with \( S, \bar{S} \) fields, can easily made much heavier by non-renormalizable interactions suppressed by the reduced Planck mass.

\( ^\dagger \)It is an interesting question whether such a large \( \Sigma \) is related to the breaking of grand unified SU(5) group. A mechanism for triplet-doublet splitting discussed \([16]\) based on \( \text{SU(5)}_L \times \text{SU(5)}_R \) may be possible within this framework.
mass $M_\ast$\footnote{Unfortunately, one possible non-renormalizable term in the superpotential $\det \Sigma/M_\ast^2$ not forbidden by the gauge invariance could screw up the inverted hierarchy mechanism. We simply assume that it is absent. Absence of an operator consistent with all symmetries of the model is not uncommon in string derived models. Yet higher order term $(\det \Sigma)^2/M_\ast^4$ is not harmful.} Adding the following terms to the superpotential,
\[ \Delta W = \frac{1}{M_\ast^2 (2!3!)^2} \epsilon_{klmn} \epsilon^{\kappa\lambda\mu\nu\rho} S^k \tilde{S}_\kappa \Sigma^l \Sigma^m \Sigma^n \Sigma^r + \frac{1}{M_\ast^3 (4!)^2} \epsilon_{klmn} \epsilon^{\kappa\lambda\mu\nu\rho} \phi^a \bar{\phi}_\kappa \Sigma^l \Sigma^m \Sigma^n \Sigma^r, \quad (13) \]
one generates masses for $S$, $\bar{S}$, $\phi$, $\bar{\phi}$ of $m_S = (v/\sqrt{5})^3/M_\ast^2$ and $m_\phi = (v/\sqrt{5})^4/M_\ast^3$. For $\phi$, $\bar{\phi}$ to be heavier than the experimental constraints of order 100 GeV\footnote{\cite{2}} we find
\[ v \gtrsim 3 \times 10^{14} \text{ GeV}. \quad (14) \]
If they happen to be this light, they leave charged tracks in the detector with anomalous $dE/dx$ because they do not have renormalizable interactions to the standard model particles due to the global SU(2) invariance. On the other hand for the largest possible $v \sim 3 \times 10^{16}$ GeV \cite{12}, $m_S \sim 10^{11}$ GeV and $m_\phi \sim 10^{9}$ GeV and then they are well beyond the experimental reach. In either case, the gauge coupling constants remain well perturbative up to the unification scale with these additional mass terms. Note that these fields also contribute to the gaugino and scalar masses when they decouple. Because of the high power in $\Sigma$ field, their soft supersymmetry breaking bilinear masses are enhanced. $S$, $\bar{S}$ fields give effectively $N_m = 9$, and $\phi$, $\bar{\phi}$ fields $N_m = 8$. This makes the degeneracy of squarks, sleptons even better.

It is an interesting question to which SU(5) each of the standard model fields are coupled above the scale $v$. The easiest option, of course, is the case where all standard model fields are coupled to one of the SU(5), while

\footnote{\cite{17} excluded a color-triplet quark bound in unit-charge hadron up to 139 GeV. Assuming 50-50 probability for $\phi^a$, $\bar{\phi}^a$ to form neutral or unit-charge meson-type states and adding $a = 1, 2$, this can be regarded as the lower bound on $m_\phi$. LEP-172 has excluded stable heavy unit-charged particles close to their kinematic reach; for instance the cross section limit from DELPHI \cite{18} can be interpreted as a lower bound of 84 GeV for a stable lepton doublet. Due to the RGE scaling, the bound on $m_\phi$ at a high scale determined from these experimental bounds is substantially lower, around 50 GeV for both quarks and leptons in $\phi$, $\bar{\phi}$.}
the other SU(5) has only $Q$, $S$, $\Sigma$ and $\phi$ as its matter content. More exotic and interesting possibility is that the Higgs field and the third generation couple to one SU(5) while the first and second generation to the other SU(5). For such a choice, the Yukawa couplings of first and second generation are naturally suppressed by a ratio $\langle \Sigma \rangle / M$, where $M$ is a scale which generates non-renormalizable interactions. It would be amusing to construct a realistic model of fermion masses along this line, but it is beyond the scope of this letter.

We now briefly discuss various distinctive features in the phenomenology of our model compared to the original models of gauge mediation (see [6] for a more detailed discussion of these issues). Since the supersymmetry breaking scale is much higher than in the original models, the logarithmic running of the supersymmetry breaking parameters is quite sizable, and the superparticle mass spectrum is significantly different from the original models because of these large logarithms [11, 12], as well as the large effective number of messengers $N_m = 8 + 9 + 8 = 25$. The experiments will be able to differentiate these predictions. The cosmological problem of a light stable gravitino overclosing the Universe [21, 22] is much less severe in our model simply because of the larger gravitino decay constant. The problem with the coherent oscillation of string moduli fields [22] is also less severe because their masses are heavier and can be diluted away by a late inflation such as thermal inflation [23]. The cosmological problem of a Polonyi-like field [24] is worse than in the original GM models because of the large scale $v$, but is much less serious than in the hidden sector models. The $v$ field acts effectively as a Polonyi field in hidden sector supersymmetry breaking, with a mass of order $10^3$ GeV. But unlike the Polonyi fields in the hidden sector case, they decay before nucleosynthesis, because of less suppressed coupling to the light fields $1/v$ rather than $1/M_*$, which provides a crucial improvement. This can dilute a pre-existing baryon asymmetry by a factor of $\sim 10^{15}(\langle v \rangle / M_*)^3$, which is rather mild for $v \lesssim 10^{15}$ GeV. Even for a larger $\langle v \rangle \sim 10^{16}$ GeV, the entropy production is still much smaller than in the hidden sector case $10^{15}$. Affleck–Dine baryogenesis [25] may be efficient enough for this purpose [24].

In summary, we have presented a simple model of gauge mediation using

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‡The $v$ direction is exactly flat at the tree-level, while is lifted due to the inverted hierarchy mechanism at the one-loop level. Therefore, the scalar mass along this direction is expected to be $m_v^2 \sim (10^4 \text{ GeV})^2/16\pi^2$. 
the $\text{SP}(4) \times \text{SU}(5)_L \times \text{SU}(5)_R$ gauge group, where the standard model gauge groups is embedded into the diagonal subgroup of two SU(5)'s. The DSB sector is coupled directly to the standard model and there is no separate messenger sector. The model is completely chiral, but yet provides vector-like messengers as a result of the dynamics. Despite the direct coupling, perturbative unification can be maintained thanks to the inverted hierarchy mechanism. The potentially non-universal supergravity contribution to the squark and slepton masses is under control. We briefly discussed several cosmologically desirable features of our model: not-too-serious Polonyi problem, and less severe cosmological problems associated with the light stable gravitino and string moduli fields than in the original gauge mediation models.

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