Strongly Coupled Gauge Mediation

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Abstract

A model of gauge mediation with $m_{3/2} \simeq \mathcal{O}(1)\text{eV}$ is considered to motivate 100 TeV colliders. Massive mediators with standard-model and supersymmetry-breaking gauge quantum numbers let low-scale dynamics induce sizable soft masses of the standard-model superpartners. We circumvent potential phenomenological difficulties that such low-scale models tend to cause.
1 Introduction

Light gravitino of mass $\lesssim 10$ eV is very interesting, since it does not have any astrophysical and cosmological problems [1]. Furthermore, it brings an intriguing possibility of the axion dark matter [2] back into our attention. This leads us to construct an explicit model of gauge mediation [3] with the gravitino mass $m_{3/2} \simeq \mathcal{O}(1)$ eV.\footnote{Several attempts have been made to construct such models of low-scale gauge mediation. See Ref.[4,5].}

2 SUSY Breaking

We adopt an $Sp(2)$ gauge theory with 6 chiral superfields $Q_i$ in the fundamental 4-dimensional representation, where $i$ is a flavor index ($i = 1, \cdots, 6$) and the gauge index is omitted. Without a superpotential, this theory has a flavor $SU(6)_F$ symmetry. This $SU(6)_F$ symmetry is explicitly broken down to a flavor $Sp(3)_F$ by a superpotential in our model. That is, we add gauge singlets $Z^a$ ($a = 1, \cdots, 14$) to obtain the tree-level superpotential

$$W_0 = Z^a(QQ)_a,$$

where $(QQ)_a$ denotes a flavor 14-plet of $Sp(3)_F$ given by a suitable combination of gauge invariants $Q_iQ_j$.

The effective superpotential which may describe the dynamics of the $Sp(2)$ gauge interaction is given by

$$W_{\text{eff}} = X(PfV_{ij} - \Lambda^6) + Z^aV_a$$

in terms of low-energy degrees of freedom

$$V_{ij} \sim Q_iQ_j, \quad V_a \sim (QQ)_a,$$

where $X$ is an additional chiral superfield and $\Lambda$ denotes a dynamical scale of the gauge interaction. This effective superpotential implies that, among the gauge invariants $Q_iQ_j$, the $Sp(3)_F$ singlet $(QQ)$ condenses as $\langle (QQ) \rangle = \Lambda^2$.\footnote{Several attempts have been made to construct such models of low-scale gauge mediation. See Ref.[4,5].}
Let us consider
\[ W = W_0 + \lambda Z(QQ). \] (4)

When the coupling \( \lambda \) is small, the effective superpotential \( W_{\text{eff}} \) implies that we obtain the following vacuum expectation values:
\[ \langle (QQ) \rangle \simeq \Lambda^2, \quad \langle V_a \rangle \simeq 0. \] (5)

Then the low-energy effective superpotential may be approximated by
\[ W_{\text{eff}} \simeq \lambda \Lambda^2 Z, \] (6)

which finally yields dynamical supersymmetry (SUSY) breaking \[6\]
\[ F_Z \simeq \lambda \Lambda^2. \] (7)

There seem two alternative possibilities in this SUSY-breaking dynamics: \( \langle \lambda Z \rangle = 0 \) or \( \langle \lambda Z \rangle \neq 0 \) \[7, 8\]. We first bear in mind the latter case \( \langle \lambda Z \rangle \neq 0 \) and later in section 4, we proceed to the other case \( \langle \lambda Z \rangle = 0 \).

3 Massive Mediators

We introduce mediator chiral superfields \( \Phi, \bar{\Phi} \), which are (anti-)fundamentals under the standard-model (SM) gauge group\(^2\) and fundamentals under the SUSY-breaking \( Sp(2) \) gauge group. The superpotential of our model for the mediators is given by \[4\]
\[ W_{\text{med}} = m \Phi \bar{\Phi}, \] (8)

where \( m \sim 4\pi \Lambda \) is a mass parameter. This is a supersymmetric mass term similar to that of the Higgs doublets (so-called \( \mu \)-term) in the supersymmetric SM. Notice that the \( Sp(2) \) theory is asymptotically free for a pair of the chiral multiplets \( \Phi, \bar{\Phi} \).

\(^2\)Namely, for definiteness, we assume that \( \Phi \) transforms as the anti-down quark and the lepton doublet. Since these mediators have the SM quantum numbers, they can decay into the SM particles (through higher-dimensional interactions) without yielding dangerously long-lived cosmological relics. The singlets \( Z^a \) and the bound states \( (QQ)_a \) can also decay into the SM particles if one allows flavor \( Sp(3)_F \) violating interactions such as \( Z(QQ)_a \).
The mediators $\Phi$ and $\bar{\Phi}$, which are massive, may be integrated out.\textsuperscript{3} Owing to their SM and SUSY-breaking gauge interactions, they induce SUSY-breaking soft masses of the SM superpartners. When the $Sp(2)$ gauge interaction becomes strong around the mass scale $m$ with the mediators decoupling \cite{4}, we obtain

\[ m_0^2 \simeq 4 \left( \frac{\alpha}{4\pi} \right)^2 \frac{1}{m^2} |\lambda F_Z|^2 \simeq 4 \left( \frac{\alpha}{4\pi} \right)^2 \frac{|\lambda^2 \Lambda^2|^2}{m^2}, \]

and

\[ m_{1/2} \simeq 4 \frac{\alpha}{4\pi m^2} \lambda F_Z \langle \lambda Z \rangle^* \simeq 4 \frac{\alpha}{4\pi m^2} \lambda^2 \Lambda^2 \langle \lambda Z \rangle^*, \]

where the overall factor 4 takes into account the mediator number of SM flavors and naive dimensional counting \cite{9} is applied. Here, $\alpha$ corresponds to the relevant SM gauge couplings to the sfermions and the gauginos. On the other hand, the gravitino mass is obtained as

\[ m_{3/2} \simeq \frac{F_Z}{\sqrt{3} M_{Pl}} \simeq \frac{\lambda \Lambda^2}{\sqrt{3} M_{Pl}}, \]

where $M_{Pl}$ denotes the reduced Planck scale: $M_{Pl} \simeq 2.4 \times 10^{18}$ GeV.

For $m \simeq 4\pi \Lambda$,

\[ m_0^2 \simeq \left( \frac{\alpha}{4\pi} \right)^2 \frac{1}{4\pi^2} |\lambda^2 \Lambda|^2, \]

and

\[ m_{1/2} \simeq \frac{\alpha}{4\pi} \frac{\lambda^2 \Lambda}{\pi}, \]

where we have used $\langle \lambda Z \rangle \simeq 4\pi \Lambda$. We see that $m_{3/2} \simeq \mathcal{O}(1)$ eV is achieved for $\sqrt{\lambda} \Lambda \simeq 100$ TeV with $m_0 \simeq m_{1/2} \simeq \mathcal{O}(100)$ GeV.

\textsuperscript{3}Through an effective identification $X \sim m \Phi \bar{\Phi}/\Lambda^6$, we see

\[ \langle m \Phi \Phi \rangle \simeq \frac{5}{3} \langle \lambda Z \rangle \Lambda^2, \]

where the factor 5 comes from the mediator number of the $Sp(2)$ quartet pairs. We suppose that this condensation is a SM singlet contained in $\Phi \otimes \bar{\Phi}$ and respects the SM gauge symmetry.
4 R-axion and \(\mu\)-term

The above setup respects R-symmetry and \(\langle \lambda Z \rangle \simeq 4\pi \Lambda\) results in an R-axion [10], which is marginally consistent to astrophysical and cosmological constraints for \(\sqrt{\lambda} \Lambda\) as low as 100 TeV. We note here that we can safely satisfy these constraints as follows. We further introduce additional chiral superfields \(\psi, \bar{\psi}\), which are (anti-)fundamentals under the SM gauge group and singlets under the SUSY-breaking \(Sp(2)\) gauge group. The superpotential for these fields is given by

\[
W_{\text{add}} = (M + kZ)\psi \bar{\psi},
\]

where \(M(\gg |\langle kZ \rangle|)\) is an additional mass parameter and \(k\) a coupling constant. This interaction makes the SUSY-breaking vacuum \(|\langle kZ \rangle| \ll M\) a local one with the supersymmetric minimum far away at \(kZ = -M\).\(^4\)

The above terms explicitly break the R symmetry and make the would-be R-axion completely harmless for \(M \ll M_{Pl}\).\(^5\) Note that we also obtain additional contributions

\[
\Delta m_0^2 \simeq \left(\frac{\alpha}{4\pi}\right)^2 \frac{|kF_{Z}|^2}{M^2},
\]

and

\[
\Delta m_{1/2} \simeq \frac{\alpha}{4\pi} \frac{kF_{Z}}{M}.
\]

This contribution can make gauginos massive even when \(\langle \lambda Z \rangle = 0\).

For the case \(\langle \lambda Z \rangle = 0\), the strongly coupled mediators provide a straightforward origin of the supersymmetric Higgsino mass term. Let us consider a superpotential of the form

\[
W = hQ\bar{\Phi}H + \bar{h}Q\Phi \bar{H},
\]

where \(h\) and \(\bar{h}\) denote coupling constants. Then, the integration of heavy modes induces the \(\mu\)-term

\[
W_{\text{eff}} \simeq -h\bar{h} \frac{\langle (QQ) \rangle}{m} H \bar{H},
\]

\(^4\)Another way to avoid an unwanted R-axion (and a supersymmetric minimum) is to introduce a singlet \(Y\) with superpotential terms such as \(\langle QQ \rangle Y + Y^3\) [7].

\(^5\)They also break an accidental \(Z_3\) symmetry in the SUSY-breaking dynamics and eliminate the corresponding domain wall problem. The scale \(M\) can be as large as an intermediate scale \(\mathcal{O}(10^{14})\) GeV for \(\sqrt{\lambda} \Lambda \simeq 100\) TeV to make the would-be R-axion sufficiently massive.
which yields

\[ \mu \simeq -h \bar{h} \frac{\Lambda^2}{m}, \]  

(19)

without a large B-term.\(^6\) It is interesting that \(\mu\) and \(m_0\) obey the same scaling law as \(\Lambda^2/m\).

5 Conclusion

We confirm the presence of an explicit model for the light gravitino of mass \(m_{3/2} \simeq \mathcal{O}(1)\) eV without any phenomenological difficulties. That constitutes one of the motivations for experiments at barely imagined colliders of, say, 100 TeV. We also note that the present setting is indeed applicable for a wide range of the gravitino mass \(m_{3/2} \simeq 1\) eV \(- 100\) TeV, among which only experiments tell what is realized in Nature.

\(^6\)The case \(\langle \lambda Z \rangle \neq 0\) suffers from a large B-term in the present setup.
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