Gravitational Supersymmetry Breaking

Izawa K.-I., T. Kugo, and T.T. Yanagida

\[1\] Yukawa Institute for Theoretical Physics, Kyoto University, Kyoto 606-8502, Japan
\[2\] Institute for the Physics and Mathematics of the Universe, University of Tokyo, Chiba 277-8568, Japan

Abstract

We consider supersymmetry breaking models with a purely constant superpotential in supergravity. The supersymmetry breaking is induced for the vanishing cosmological constant. As a hidden mediation sector of supersymmetry breaking, it naturally leads to a split spectrum in supersymmetric standard model. We also point out possible utility of our setup to construct nonlinear sigma model and/or Fayet-Iliopoulos-like term in broken supergravity.
It naively seems that perturbative gravity as an effective theory provides only small corrections to low-energy particle physics on the flat background and thus it would not essentially alter the physical contents such as phase structure from those of the corresponding particle physics without gravity. However, as for supergravity in superspace, such a naive guess is not necessarily true. Simple-minded decoupling of the supergravity multiplet sometimes results in qualitatively different physics on the flat background.

In this paper, we consider supersymmetry (SUSY) breaking models with a purely constant superpotential in supergravity. Although the constant superpotential has no physical meaning in the case of rigid SUSY without gravity, the SUSY breaking is induced for the vanishing cosmological constant in the case of supergravity. As a hidden mediation sector of SUSY breaking, it naturally leads to a split spectrum in supersymmetric standard model. We also point out possible utility of our setup to construct nonlinear sigma model and/or Fayet-Iliopoulos-like $D$-term in broken supergravity.

Let us adopt a purely constant superpotential\(^1\)

\[
W = c \neq 0,
\]

and a Kähler potential $K$ of chiral superfields $\Phi_i$, where the subscript $i$ denotes a flavor index. Our crucial premise is that the Kähler potential $K$ is tuned to make the cosmological constant vanishing so that SUSY is spontaneously broken with its breaking scale solely determined by the constant $W = c$ or the gravitino mass\(^2\). We now proceed to investigate physical contents of this setup.

With the graviton field and all the fermionic fields vanishing, the supergravity action takes the form

\[
\int d^4 x \, d^2 \theta d^2 \theta^* \varphi \varphi^* \Omega + \left( \int d^4 x \, d^2 \theta \, \varphi^3 W + \text{h.c.} \right) + \cdots,
\]

where $\Omega = -3 \exp(-K/3)$ with the reduced Planck scale unity and the chiral compensator $\varphi = 1 + \theta^2 F$. Here the compensator lowest component is gauge-fixed to be one for

---

1. Classically the Kähler-Weyl transformation can achieve the constant superpotential without loss of generality, while quantum-mechanically it is generically anomalous.

2. Concrete structure of the Kähler potential was investigated by Hebecker \(^\text{[1]}\) in the case of a single chiral superfield with the constant superpotential.
simplicity. The ellipses stand for extra derivative terms for (the lowest components of) \( \Phi_i \) which come from the vector auxiliary field \([2]\).

The algebraic equations of motion for the auxiliary fields \( F (\equiv F_\varphi) \) and \( F_{\Phi_i} \) determine them in terms of the dynamical fields \( \Phi_i \), which yield the (non-canonically normalized) scalar potential

\[
V = -3cF,
\]

as will be shown just below. Hence the vanishing of the cosmological constant results in the vanishing of the compensator auxiliary field \( F \) in the vacuum. In particular, it means an important implication that the so-called anomaly mediation \([3]\) does not occur.

More generally, we can show that the vanishing \( F \) is realized in the case of vanishing \( F_W \) for a generic superpotential \( W \). The action

\[
\int d^4x d^2\theta d^2\theta^* \varphi \varphi^* \Omega + \left( \int d^4x d^2\theta \varphi^3W + \text{h.c.} \right),
\]

with the chiral compensator \( \varphi = 1 + \theta^2F \), yields the scalar potential

\[
-V = |F|^2\Omega + F\Omega^iF^*_i + F^*\Omega^iF_i + \Omega^{ij}F^*_iF_j + 3FW + F_W + 3F^*W^* + F^*_W
\]

\[
= F^*(F\Omega + \Omega^iF_i + 3W^*) + F^*_i(F\Omega^i + \Omega^{ij}F_j + W^i) + 3FW + F_W,
\]

with \( F_i = F_{\Phi_i} \) and \( F_W = W^iF_i \), where the superscripts \( i \) to \( \Omega \) and \( W \) denote partial derivatives with respect to \( \Phi_i \) and the repeated indices are summed over.

The algebraic equations of motion for the auxiliary fields \( F^* \) and \( F^*_i \) read, respectively,

\[
F\Omega + \Omega^iF_i + 3W^* = 0,
\]

\[
F\Omega^i + \Omega^{ij}F_j + W^i = 0,
\]

which lead to \( V = -3FW - F_W \).

With this, we see that the demand \( \langle V \rangle = \langle F_W \rangle = 0 \) generally implies \( \langle F \rangle = 0 \) under \( \langle W \rangle \neq 0 \). That is, the vacuum value of the compensator auxiliary field vanishes. When \( W \) is a constant \( (W = c) \), in particular, then \( F_W = 0 \) since \( W^i = 0 \), which results in \( V = -3cF \) as advocated above. Eq.(6) also implies that, when \( \langle F \rangle = 0 \), SUSY is spontaneously broken since \( \langle W \rangle \neq 0 \) requires \( \langle \Omega^iF_i \rangle \neq 0 \) so that \( \langle F_i \rangle \neq 0 \) for some \( i \).
Let us turn to phenomenological implications of our setup with the constant super-potential ($W = c$). Since we assume the unique input scale $c$ as the origin of SUSY breaking, we can make order estimation of soft masses based on possible forms of effective operators. Under the vanishing compensator auxiliary field ($F = 0$), we have three possible contributions to the gaugino mass term or the lowest component of $W^\alpha W_\alpha$ with the reduced Planck scale as the cutoff.

i) gravitational mediation:

The contribution to the gaugino masses is estimated to be of order $c^3$ due to effective operators

$$\int d^2\theta d^2\theta^* |\Phi_i|^2 W^\alpha W^\alpha W_\alpha,$$

of the neutral R charge [4]. In other words, the maximum $\langle F_\Phi \rangle$ is naturally of order the gravitino mass $m_{3/2}$, so that the gaugino masses are expected to be given by $O(m_{3/2}^3)$. To obtain the gaugino masses at the order of 100GeV, we need to choose $m_{3/2} \simeq 10^{13}\text{GeV}$.

ii) moduli mediation:

The contribution to the gaugino masses is estimated to be of order $c$ or less due to effective operators

$$\int d^2\theta \Phi_i W^\alpha W_\alpha.$$

The above two contributions constitute gravity mediation of SUSY breaking suppressed by the Planck scale, while gauge mediation [5] can incorporate a lower scale of SUSY breaking mediation as follows.

iii) gauge mediation:

The contribution to the gaugino masses is estimated to be of order $c/M$ or less due to effective operators

$$\int d^2\theta \frac{\Phi_i}{M} W^\alpha W_\alpha,$$

where $M$ denotes the scale of gauge mediation.

In the three contributions, the first one is ubiquitous in contrast to the remaining two. Without the moduli and gauge mediation, the gravitational mediation naturally
leads to a split spectrum in supersymmetric standard model due to the vanishing anomaly mediation, where gaugino masses are many orders of magnitude lighter than scalar masses.

One peculiar realization may appear as a tiny bino mass\(^3\) less than 10eV, which avoids cosmological problems due to thermal bino abundance. Heavier bino mass would necessitate tuning for co-annihilation with slepton in order to achieve efficient bino annihilation in the early universe.

Even in the presence of other dominant contributions to the bino mass, the absence of the anomaly mediation may have some virtues. For instance, sizable anomaly mediated contribution might induce CP-violating phase difference among different origins of the gaugino masses\(^6\), which seems undesirable in view of experimentally limited CP-violating effects.

Finally a few comments are in order. As a by-product, the present Kähler potential can be utilized to construct nonlinear sigma model and/or Fayet-Iliopoulos-like D-term in broken supergravity.

For simplicity, let us restrict ourselves to the single chiral superfield \(\Phi\) with \(\Omega(\phi)\) and \(W = c \neq 0\), where \(\phi = \Phi + \Phi^*\). Then, in view of Eq(6), the vanishing cosmological constant requires

\[
\langle \Omega'(\phi) F_\phi \rangle + 3c^* = 0,
\]

which implies \(\langle \Omega'(\phi) \rangle \neq 0\).

When we have the Kähler potential \(G\) of a nonlinear sigma model with \(\langle G\rangle = 0\) in rigid SUSY\(^7\), we can obtain a corresponding nonlinear sigma model in broken supergravity in terms of the above setup with \(\Omega(\phi + G)\), since its expansion

\[
\Omega(\phi + G) = \Omega(\phi) + \Omega'(\phi) G + \frac{1}{2} \Omega''(\phi) G^2 + \cdots
\]

contains the leading term \(\Omega'(\phi) G\) under \(\langle \Omega'(\phi) \rangle \neq 0\), provided SUSY is almost intact in the sigma model sector.

Even if a symmetry transformation in rigid SUSY varies \(G\) by a holomorphic function plus its complex conjugate, the variation can be absorbed in \(\phi\), and the symmetry is maintained in the supergravity.

\(^3\)The other gauginos are assumed to obtain larger masses from other contributions.
We may also realize a Fayet-Iliopoulos-like $D$-term in supergravity \cite{8} by replacing the above $G$ with a $U(1)$ real superfield \cite{4} whose gauge transformation has a similar property of holomorphy.

Acknowledgements

This work is supported by the Grant-in-Aid for Yukawa International Program for Quark-Hadron Sciences, the Grant-in-Aid for the Global COE Program ”The Next Generation of Physics, Spun from Universality and Emergence”, and World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan. TK is also partially supported by a Grant-in-Aid for Scientific Research (B) (No. 20340053) from the Japan Society for the Promotion of Science.

References

\begin{itemize}
  \item [1] A. Hebecker, \url{arXiv:hep-th/0407196}.
  \item [2] T. Kugo and S. Uehara, Nucl. Phys. B 226 (1983) 49, and references therein.
  \item [3] L. Randall and R. Sundrum, \url{arXiv:hep-th/9810155}.
  \hspace{1em} G.F. Giudice, M.A. Luty, H. Murayama, and R. Rattazzi, \url{arXiv:hep-ph/9810442}.
  \item [4] N. Arkani-Hamed and S. Dimopoulos, \url{arXiv:hep-th/0405159}.
  \item [5] For a review, G.F. Giudice and R. Rattazzi, \url{arXiv:hep-ph/9801271}.
  \item [6] See M. Endo, M. Yamaguchi, and K. Yoshioka, \url{arXiv:hep-ph/0311206}.
  \item [7] For a review, M. Bando, T. Kugo, and K. Yamawaki, Phys. Rept. 164 (1988) 217.
  \item [8] See Z. Komargodski and N. Seiberg, \url{arXiv:0904.1159} [hep-th].
\end{itemize}

\footnote{If instead we replace the above $G$ with $G$ plus a $U(1)$ real superfield, the chiral superfield $\Phi$, which contains a massless scalar, becomes a gauge degree of freedom so that the physical massless modes are exclusively contained in those of the original sigma model.}