

Spontaneous SUSY breaking without R symmetry in supergravity

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We discuss spontaneous supersymmetry (SUSY) breaking in a model with an anomalous $U(1)_A$ symmetry. In this model, the size of the each term in the superpotential is controlled by the $U(1)_A$ charge assignment and SUSY is spontaneously broken via the Fayet-Iliopoulos of $U(1)_A$ at the metastable vacuum. In the global SUSY analysis, the gaugino masses become much smaller than the sfermion masses, because an approximate R symmetry appears at the SUSY breaking vacuum. In this paper, we show that gaugino masses can be as large as gravitino mass, taking the supergravity effect into consideration. This is because the R symmetry is not imposed so that the constant term in the superpotential, which is irrelevant to the global SUSY analysis, largely contributes to the soft SUSY breaking terms in the supergravity. As the mediation mechanism, we introduce the contributions of the field not charged under $U(1)_A$ and the moduli field to cancel the anomaly of $U(1)_A$. We comment on the application of our SUSY breaking scenario to the grand unified theory.

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I. INTRODUCTION

The conditions for spontaneous supersymmetry (SUSY) breaking have been pointed out in the literatures [1,2]. Nelson and Seiberg [2] speculated that without R symmetry SUSY cannot be broken spontaneously in global minimum of the scalar potential with generic interactions, and no counter example for this speculation has been known. On the other hand, since R symmetry forbids gaugino and Higgsino masses, the R symmetry must be broken to obtain realistic models. However, spontaneous R symmetry breaking results in massless R axion which is potentially suffering from astrophysical problems. In the supergravity (SUGRA), R symmetry can be broken by constant term in the superpotential without changing the arguments in global SUSY (and in many cases, it needs to obtain Minkowski space-time [3]), and it gives the R axion massive [4]. However, once such a R symmetry breaking term is introduced, we have no reason to keep R symmetry only in SUSY breaking sector.

One solution is to break R symmetry explicitly, although the SUSY breaking vacua become metastable [3,5–11]. In Ref. [5], we examined a simple SUSY breaking model without R symmetry, which has following features:

- (i) Since all interactions which are allowed by the anomalous $U(1)_A$ gauge symmetry, which has Fayet-Iliopoulos (FI) D -term [12], are introduced with $O(1)$ coefficients, R symmetry is maximally broken.
- (ii) SUSY is spontaneously broken in metastable vacua, at which approximate R symmetry appears.
- (iii) Massless R axion does not appear because of the explicit R symmetry breaking terms.

Unfortunately, it seems to be difficult to apply this SUSY breaking model to realistic scenario since the gaugino masses become much smaller than the sfermion masses because of the approximate R symmetry. The gaugino masses explicitly break the R symmetry. Even if the R symmetry is not the one of the fundamental symmetry as in our model, the gaugino mass is vanishing because of the accidental R symmetry at the SUSY breaking vacuum.

In this paper, we will point out that the sizable gaugino masses can be produced if the SUGRA effects are taken into account. The essential point is that the constant term in the superpotential, which is irrelevant to the global SUSY analysis, contributes to the SUSY breaking dynamics. This term breaks not only R symmetry explicitly, but also contributes to the SUSY breaking terms. As the result, generically, the gaugino masses become of order the gravitino mass, mediated by the extra fields, such as

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moduli fields. That is nothing but the usual results of the gravity mediation. However, since the vacuum structure is modified by including the SUGRA effects and the application of the obtained model is important, we will stress in this paper that even with approximate R symmetry in global SUSY calculation, the gaugino masses can be around the gravitino mass when the SUGRA effects are included.

In Sec. II, we review the simple SUSY breaking model without R symmetry. In Sec. III, we discuss the SUSY breaking with the SUGRA effects in the model. In Sec. IV, we have summary and discussion.

II. REVIEW OF SUSY BREAKING MODEL WITHOUT R SYMMETRY

The SUSY breaking with the FI term has been studied in the global SUSY [13–20]. In this section, we give a brief review of a SUSY breaking model without R symmetry in which the FI term of anomalous $U(1)_A$ gauge symmetry plays an important role in breaking SUSY spontaneously, following Ref. [5].

First of all, we remind you of a simple FI model [13] which has R symmetry. In this model, there are two fields, S and Θ whose $U(1)_A$ charges are $s \gg 1$ and $\theta = -1$, respectively. Note that the large charge, s , realizes the hierarchy between the SUSY breaking scale and the cutoff scale (Planck scale). If the R charges of S and Θ are 2 and 0, respectively, the generic superpotential is given as

$$W = y\Lambda^3 \frac{S}{\Lambda} \left(\frac{\Theta}{\Lambda} \right)^s \quad (1)$$

where y and Λ are the coefficient and the cutoff of the model. The potential is given as

$$V = |F_S|^2 + |F_\Theta|^2 + \frac{1}{2} D_A^2, \quad (2)$$

where the F -terms and the D -term are

$$F_S^* = -\frac{\partial W}{\partial S} = -y\Theta^s \quad (3)$$

$$F_\Theta^* = -\frac{\partial W}{\partial \Theta} = -ys\Theta^{s-1}, \quad (4)$$

$$D_A = -g(\xi^2 - |\Theta|^2 + s|S|^2), \quad (5)$$

when the Kähler potential K is taken to be canonical. Here, g and ξ^2 are the gauge coupling constant of the $U(1)_A$ gauge symmetry and the FI parameter of the FI term, respectively. Note that we usually take $\Lambda = 1$ for simplicity in this paper. The vacuum expectation values (VEVs) of these fields are determined by the minimization of the potential as

$$\langle S \rangle = 0, \quad \langle \Theta \rangle \equiv \lambda \sim \frac{\xi}{\Lambda}. \quad (6)$$

Note that R symmetry is not broken at all although SUSY is broken spontaneously. At this vacuum, the VEVs of the F -terms and the D -term are given as

$$\langle F_S \rangle \sim \lambda^s, \quad \langle F_\Theta \rangle = 0, \quad \langle D_A \rangle \sim \frac{s}{g} \lambda^{2s-2}. \quad (7)$$

λ is expected to be $\mathcal{O}(0.1)$ [5], so that the SUSY breaking scale, that is given by the F -term and the D -term, becomes much smaller than the cutoff scale.

Second, we consider another model where the R symmetry is not imposed to the above setup [5]. Then the generic superpotential is given by

$$W = W(S\Theta^s), \quad (8)$$

where $W(x)$ is a function of x and expected to be a polynomial function as $W(x) = \sum_{n=0} a_n x^n$. The coefficients a_n are expected to be of order one generically. Then SUSY vacua appear because all of the F -terms and the D -term,

$$F_S^* = -W'(S\Theta^s)\Theta^s \quad (9)$$

$$F_\Theta^* = -W'(S\Theta^s)s\Theta^{s-1}, \quad (10)$$

$$D_A = -g(\xi^2 - |\Theta|^2 + s|S|^2), \quad (11)$$

can be vanishing at the same time. Here, $W'(x) \equiv \frac{dW}{dx}$ is defined. Indeed, $W'(S\Theta^s) = 0$ and $D_A = 0$ can be satisfied by fixing two variables, $\langle S \rangle$ and $\langle \Theta \rangle$, which become of order one generically. On the other hand, as pointed out in Ref. [5], this model has metastable vacua where SUSY is spontaneously broken. The metastable vacua are near the vacua with the R symmetry in Eq. (6) as

$$\langle S \rangle \sim \frac{\Theta^{s+2}|a_2|}{s^2|a_1|} \sim \frac{1}{s^2} \lambda^{s+2}, \quad \langle \Theta \rangle \equiv \lambda \sim \frac{\xi}{\Lambda}. \quad (12)$$

The VEVs of F and D_A are given as

$$\langle F_S \rangle \sim \lambda^s, \quad \langle F_\Theta \rangle \sim \frac{1}{s} \lambda^{2s+1}, \quad \langle gD_A \rangle \sim s\lambda^{2s-2}. \quad (13)$$

It is obvious that the vacua have an approximate R symmetry when $\lambda \ll 1$ and $s \gg 1$. Note that the VEV of S is roughly proportional to the R symmetry breaking parameter a_2 . Thus, the soft SUSY breaking terms that break the R symmetry, e.g., the gaugino masses, become quite small if this model is applied to the realistic models.

This is the conclusion, based on the global SUSY analysis, where the constant term in the superpotential, a_0 , is ignored. When the SUGRA effect is taken into

consideration, we can expect that a_0 largely contributes to the gaugino masses. In the next section, we discuss this model where the R symmetry is not imposed in the SUGRA.

III. SUGRA EFFECTS

In this section, we will show that SUGRA effects are not negligible especially in the models with approximate R symmetry as in the previous section. The essential point is that the constant term a_0 of the superpotential, which breaks R symmetry, contributes to the vacua in SUGRA calculation, but not in global SUSY calculation. The VEV $\langle S \rangle$ in SUGRA calculation is proportional to a_0 , which is much larger than $\langle S \rangle$ in global SUSY calculation in Eq. (12). Therefore, the breaking effect of $U(1)_R$ is larger at vacua in SUGRA calculation than in global SUSY calculation. Moreover, if there is at least one $U(1)_A$ singlet field, then the F component of the singlet field can become sizable because the constant superpotential contributes to the F component of the singlet field. Since the singlet field can couple to superfield strength, the nonvanishing F of the singlet can contribute to gaugino masses. In addition, the F component of the moduli field, that is required to cancel the gauge anomaly in anomalous $U(1)_A$ gauge theory, can have nonvanishing VEV because it includes the term proportional to the VEV of superpotential. These contributions can give the gaugino masses around the gravitino mass.

The superpotential is the same as in Eq. (8), although a_0 is determined by $\langle V \rangle = 0$ and therefore the gravitino mass $m_{3/2}$ is fixed by $\langle V \rangle = 0$ because $a_0 = m_{3/2} M_{\text{Pl}}^2$. Here, M_{Pl} is the reduced Planck scale. We treat the cutoff scale Λ and the Planck scale differently, as in Horava-Witten theory [21] or in natural grand unified theory (GUT) [22]. Then the scalar potential is written as

$$V = V_F + V_D,$$

$$V_F = e^{K/M_{\text{Pl}}^2} \left(|D_S W|^2 + |D_\Theta W|^2 - 3 \frac{|W|^2}{M_{\text{Pl}}^2} \right), \quad (14)$$

$$V_D = \frac{g^2}{2} (\xi^2 + s|S|^2 - |\Theta|^2)^2, \quad (15)$$

where the following functions are defined:

$$K = |S|^2 + |\Theta|^2, \quad (16)$$

$$D_S W = W'(S\Theta^s)\Theta^s + \frac{S^*}{M_{\text{Pl}}^2} W, \quad (17)$$

$$D_\Theta W = W'(S\Theta^s)sS\Theta^{s-1} + \frac{\Theta^*}{M_{\text{Pl}}^2} W. \quad (18)$$

The stationary conditions for the potential give the VEVs of S and Θ as

$$\langle \Theta \rangle \sim \lambda, \quad \langle S \rangle \sim \frac{\lambda^2}{s} \frac{\Lambda}{M_{\text{Pl}}}, \quad (19)$$

and the vanishing cosmological constant $\langle V \rangle = 0$ fixes $\langle W \rangle \sim \lambda^s M_{\text{Pl}}$, which determines the gravitino mass $m_{3/2}$ as $m_{3/2} = \langle W \rangle / M_{\text{Pl}}^2 \sim \lambda^s \frac{\Lambda}{M_{\text{Pl}}}$. The VEVs of $D_S W$, $D_\Theta W$, and D_A are given as

$$\langle D_S W \rangle \sim \lambda^s, \quad \langle D_\Theta W \rangle \sim \lambda^{s+1} \frac{\Lambda}{M_{\text{Pl}}}, \quad \langle g D_A \rangle \sim s \lambda^{2s-2}. \quad (20)$$

Note that $\langle S \rangle$ is not so small at all especially when $\Lambda \sim M_{\text{Pl}}$, and therefore the VEVs break R symmetry completely. This may induce gaugino masses if this mechanism is embedded in a realistic model. When $S = S_r e^{i\frac{\phi_S}{\sqrt{2}\langle S \rangle}}$, the masses of S_r , ϕ_S and Θ are given as $\frac{sM_{\text{Pl}}}{\lambda\Lambda} m_{3/2}$, $\frac{sM_{\text{Pl}}}{\lambda\Lambda} m_{3/2}$, and Λ , respectively.

Moreover, the F component of a field Z which is neutral under $U(1)_A$ can have nonvanishing VEV because of the contribution from the constant superpotential. This also gives sizable gaugino masses which can be around the gravitino mass. We show this in an explicit model in which the neutral field Z is added to the above model. Then the scalar potential is written as

$$V = V_F + V_D,$$

$$V_F = e^{K/M_{\text{Pl}}^2} \left(|D_S W|^2 + |D_\Theta W|^2 + |D_Z W|^2 - 3 \frac{|W|^2}{M_{\text{Pl}}^2} \right), \quad (21)$$

$$V_D = \frac{g^2}{2} (\xi^2 + s|S|^2 - |\Theta|^2)^2, \quad (22)$$

where the following functions are defined:

$$K = |S|^2 + |\Theta|^2 + |Z|^2, \quad (23)$$

$$D_S W = W'(S\Theta^s)\Theta^s + \frac{S^*}{M_{\text{Pl}}^2} W, \quad (24)$$

$$D_\Theta W = W'(S\Theta^s)sS\Theta^{s-1} + \frac{\Theta^*}{M_{\text{Pl}}^2} W, \quad (25)$$

$$D_Z W = \dot{W} + \frac{Z^*}{M_{\text{Pl}}^2} W, \quad (26)$$

where the superpotential is given as

$$W = \sum_{n=0} a_n(Z) (S\Theta^s)^n = W(S\Theta^s). \quad (27)$$

Here, $W' = \frac{dW(x)}{dx}$ and $\dot{W} = \frac{\partial W}{\partial Z} = \sum_{n=0} \frac{da_n}{dZ} (S\Theta^s)^n$. The VEVs are essentially the same as the previous results

except the VEV of Z . The stationary condition $\partial V/\partial Z = 0$ determines the VEV of $D_Z W$ as

$$\langle D_Z W \rangle \sim \frac{\langle \dot{W}' \rangle}{\langle W' \rangle \langle \ddot{W} \rangle} m_{3/2}^2 M_{\text{Pl}}^2 \sim \frac{\dot{a}_1 \langle Z \rangle}{a_1 \langle Z \rangle \langle \ddot{a}_0 \rangle} m_{3/2}^2 M_{\text{Pl}}^2. \quad (28)$$

Since $\dot{a}_1 \sim a_1$ is expected, $\langle D_Z W \rangle$ becomes

$$\langle D_Z W \rangle \sim \begin{cases} m_{3/2} & (\text{when } \langle \ddot{a}_0 \rangle \sim m_{3/2} M_{\text{Pl}}^2) \\ m_{3/2}^2 M_{\text{Pl}}^2 & (\text{when } \langle \ddot{a}_0 \rangle \sim 1) \end{cases} \quad (29)$$

Note that the VEV \ddot{W} is dependent on the mechanism to realize $\langle V \rangle = 0$. For example, if $a_0(Z) = m_{3/2} M_{\text{Pl}}^2 \hat{a}_0(Z)$, where $\hat{a}_0(Z)$ is a polynomial function with $O(1)$ coefficients, then the upper result in Eq. (29) is realized, and the mass of Z becomes $\frac{M_{\text{Pl}}^2}{\Lambda} m_{3/2}$. If $a_0(Z)$ is a polynomial function with $O(1)$ coefficient whose VEV is $\langle a_0 \rangle = m_{3/2} M_{\text{Pl}}^2$, the lower result is realized and the mass of Z becomes Λ . It is important that the VEV of $D_Z W$ can be $O(m_{3/2})$ and therefore, gaugino masses can be $O(m_{3/2})$ because the neutral field Z can couple with the kinetic functions of vector multiplets.

There is another contribution to gaugino masses from the F -term of the moduli fields T , that can be $O(m_{3/2})$. Since $U(1)_A$ gauge symmetry is given by

$$V_A \rightarrow V_A + \frac{i}{2} (\tilde{\Lambda} - \tilde{\Lambda}^\dagger), \quad (30)$$

$$T \rightarrow T + \frac{i}{2} \delta_{\text{GS}} \tilde{\Lambda}, \quad (31)$$

where V_A , $\tilde{\Lambda}$, and δ_{GS} are vector multiplet of $U(1)_A$, a gauge parameter chiral superfield, and dimensionless parameter which has relations¹

$$2\pi^2 \delta_{\text{GS}} = \frac{1}{3k_A} \text{tr} Q_A^3 = \frac{1}{24} \text{tr} Q_A > 0. \quad (32)$$

The anomaly of $U(1)_A$ can be cancelled [23] via

$$\mathcal{L}_{\text{gauge}} = \frac{1}{4} \int d\theta^2 k_A T W_A^\alpha W_{A\alpha} + \text{H.c.}, \quad (33)$$

where W_A^α and k_A are the super field strength of V_A and Kac-Moody level of $U(1)_A$, respectively. The $U(1)_A$ invariant Kähler potential is given as

$$\mathcal{K} = S^\dagger e^{-2g_S V_A} S + \Theta^\dagger e^{2g V_A} \Theta + f(T + T^\dagger - \delta_{\text{GS}} V_A). \quad (34)$$

The FI term can be given as

¹The relations can be satisfied by choosing normalization factor of $U(1)_A$ gauge symmetry and/or k_A , although we do not fix these explicitly in this paper.

$$\begin{aligned} \int d^4\theta f(T + T^\dagger - \delta_{\text{GS}} V_A) &= \left(-\frac{\delta_{\text{GS}} f'}{2} \right) D_A \cdots \\ &\equiv \xi^2 D_A + \cdots. \end{aligned} \quad (35)$$

Note that $\langle f' \rangle$ must be negative to obtain positive ξ^2 .

The scalar potential is given as

$$\begin{aligned} V &= V_F + V_D, \\ V_F &= e^{K/M_{\text{Pl}}^2} \left(|D_S W|^2 + |D_\Theta W|^2 + f'^{-1} |D_T W|^2 - 3 \frac{|W|^2}{M_{\text{Pl}}^2} \right), \end{aligned} \quad (36)$$

$$V_D = \frac{g^2}{2} \left(-\frac{\delta_{\text{GS}} f'}{2} + s|S|^2 - |\Theta|^2 \right)^2, \quad (37)$$

where the following functions are defined:

$$D_S W = W'(S\Theta^s)\Theta^s + \frac{S^*}{M_{\text{Pl}}^2} W, \quad (38)$$

$$D_\Theta W = W'(S\Theta^s)sS\Theta^{s-1} + \frac{\Theta^*}{M_{\text{Pl}}^2} W, \quad (39)$$

$$D_T W = \frac{f'}{M_{\text{Pl}}^2} W, \quad (40)$$

where the superpotential is given as in Eq. (8). We expand the function f around $2T_0$ as $f(T+T^\dagger) = \sum_n \frac{b_n}{n!} (T+T^\dagger - 2T_0)^n$. The stationary condition $\partial V/\partial T = 0$ gives

$$m_{3/2}^2 \left(2f' - \frac{(f')^2 f'''}{(f'')^2} \right) + gD_A \left(\frac{\delta_{\text{GS}}}{2} f'' \right) = 0. \quad (41)$$

The second derivative of the scalar potential becomes

$$\frac{\partial^2 V}{\partial T^2} \sim \frac{\delta_{\text{GS}}^2}{4} (f'')^2 > 0. \quad (42)$$

Therefore, if

$$m_{3/2}^2 \left(2b_1 - \frac{b_1^2 b_3}{b_2^2} \right) + gD_A \left(\frac{\delta_{\text{GS}}}{2} b_2 \right) = 0 \quad (43)$$

is satisfied, the VEV $\text{Re}\langle T \rangle = T_0$ is (meta)stable. Note that the moduli can easily be stabilized because of the D -term. The scalar masses can be calculated as

$$\begin{aligned} m_T &\sim \frac{\delta_{\text{GS}}}{2} \Lambda, & m_\Theta &\sim \Lambda, \\ m_{S_r} &\sim \frac{sM_{\text{Pl}}}{\lambda\Lambda} m_{3/2}, & m_{\phi_s} &\sim \frac{sM_{\text{Pl}}}{\lambda\Lambda} m_{3/2}, \end{aligned} \quad (44)$$

except massless axion. Actually, this scalar potential has an global $U(1)$ symmetry in addition to $U(1)_A$ gauge symmetry, which transforms only S and Θ as $U(1)_A$ but not T . Because of this additional $U(1)$ symmetry, a Nambu-Goldstone boson appears. If nonperturbative interactions are allowed in the superpotential or in the Kähler potential like $\Theta e^{2T/\delta_{GS}}$ or $S e^{-2sT/\delta_{GS}}$ which break the additional global $U(1)$ symmetry, the axion becomes massive. Otherwise, this axion works as QCD axion, which may solve the strong CP problem [24]. The effective Peccei-Quinn scale becomes $F_{PQ} \sim \frac{\Lambda}{8\pi^2}$, which is around 10^{14} GeV if $\Lambda \sim \Lambda_G \sim 2 \times 10^{16}$ GeV. The F -term of T becomes

$$F_T = (f'')^{-1} D_T W = \frac{f'}{f''} \frac{W}{M_{Pl}^2}, \quad (45)$$

which gives gaugino masses as $k_A \frac{b_1}{b_2} m_{3/2}$.

In conclusion, even if the vacua determined by global SUSY calculation have approximate R symmetry, SUGRA effects can change them to the vacua without R symmetry. As the result, gaugino masses can be around gravitino mass.

IV. SUMMARY AND DISCUSSION

It is one of the important issues how a realistic SUSY breaking vacuum can be realized, in supersymmetric models. The R symmetry seems to play an important role in the SUSY breaking, but it causes the massless Goldstone boson and prevents generating the nonvanishing gaugino masses. As pointed out in Ref. [8], one realistic SUSY breaking vacuum could be realized if explicit R symmetry breaking terms are enough small for the life time of the vacuum to be longer than the age of our universe. This scenario, however, requires the explanation of the origin of the tiny R symmetry breaking terms. In addition, the gaugino actually needs large R symmetry breaking effects to gain large mass.² Even taking the SUGRA effect into consideration, this situation does not change [3]. In the SUGRA, the large constant term in the superpotential, that breaks the R symmetry, is necessary for the vanishing cosmological constant in many cases [3], so that the situation may become worse compared to the global SUSY case. Thus, we need to find the symmetry or the dynamics that can replace the role of the R symmetry, in order to lead a realistic SUSY breaking vacuum. We have discussed spontaneous SUSY breaking via the FI term in a

²Even if the R symmetry is spontaneously broken, the gaugino masses are often vanishing in the gauge mediation scenario [25].

model which has anomalous $U(1)_A$ symmetry. The R symmetry is not imposed, but an approximate R symmetry appears at the metastable SUSY breaking vacua in the global SUSY analysis. Then, the gaugino masses become much smaller than the sfermion masses in the global SUSY as shown in Ref. [5]. In this paper, we have pointed out that if the SUGRA effects are taken into account, the R symmetry is largely broken by the constant term in the superpotential at the metastable SUSY breaking vacua, and as the result, the gaugino masses can be of order the gravitino mass.

In our calculation, we have adopted the cutoff scale which can be different from the Planck scale as in natural GUT or in Horava-Witten theory. The application of this mechanism to the natural GUT is interesting. In the natural GUT, the doublet-triplet splitting problem can be solved under a reasonable assumption in which all interactions including higher dimensional interactions are introduced with $O(1)$ coefficients. An important point is that the natural GUT has the cutoff scale which is the usual GUT scale smaller than the Planck scale. As the result, sfermion masses become around 100 times larger than the gaugino masses. Namely high scale SUSY (or split SUSY) is realized in the model. The details will be discussed in a separate paper.

In the explicit model we discussed, we adopted an anomalous $U(1)_A$ gauge symmetry with FI term. The anomaly can be cancelled by Green-Schwarz mechanism [23] in which moduli plays an important role. We have shown explicitly that all scalar fields become massive except a Nambu-Goldstone boson which can solve the strong CP problem. Especially when the cutoff scale is lower than the Planck scale, these massive modes become much heavier than the gravitino mass.

To obtain the gaugino masses around 1 TeV which are the same order of the gravitino mass, the gravitino problem [26–33] becomes serious. One possible way to avoid the problem is to adopt low reheating temperature of inflation. We will not discuss this problem further in this paper.

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