Gravitational Gauge Mediation

Ryuichiro Kitano

Stanford Linear Accelerator Center, Stanford University, Stanford, CA 94309

Abstract

It is often the case that naive introduction of the messenger sector to supersymmetry breaking models causes the supersymmetry restoration. We discuss a possibility of stabilizing the supersymmetry broken vacuum by the gravitational interaction.

1 Introduction

In scenarios where the electroweak scale is stabilized by supersymmetry, it is expected that future collider experiments give us hints of how the standard model sector feels the supersymmetry breaking. It is now important to discuss possible mechanisms for the supersymmetry breaking and its mediation.

In field theory, spontaneous breaking of supersymmetry is quite easily realized. An obvious example is the Polonyi model in which a singlet chiral superfield S has a linear superpotential $W = \mu^2 S$. This model also serves as the low energy effective theory of the O'Raifeartaigh model by introducing a non-minimal Kähler potential term $K \ni -(S^\dagger S)^2 / \widetilde{\Lambda}^2$ with $\widetilde{\Lambda}$ being the mass scale of the particles which are integrated out. The same effective theory can also be realized in models with strongly coupled gauge theories such as in the IYIT model in Ref. [1, 2]. Recently it has been noticed that a wide class of supersymmetric QCD leads to the above effective theory around the meta-stable vacuum [3].

Even though the spontaneous supersymmetry breaking may be a common feature of the quantum field theory, the mediation of the supersymmetry breaking to the standard model sector is not so simple. Of course, the gravity mediation is the simplest way to communicate with the hidden sector. However, we study the possibility of the gauge mediation [4] (See [5] for earlier works.) partly because it provides us with a solution to the flavor problem.

A naive way of realizing the gauge mediation is to introduce vector-like messenger particles q and $\bar{q}$ which carry the standard model quantum numbers and assume a coupling $W \ni -\lambda S q \bar{q}$ in the superpotential. Assuming the non-vanishing vacuum expectation values in the lowest and the F-component of S, the gaugino masses are obtained by the formula $m_{1/2} = (\alpha/4\pi)(F_S/S)$ at one-loop level. However, this model has a supersymmetric and hence global minimum where $S = 0$ and $q = \bar{q} = \sqrt{\mu^2/\lambda}$. Therefore the question is whether it is possible to have a meta-stable vacuum with non-vanishing value of S [6]. For example, in the O'Raifeartaigh model, the sign of the coefficient of the Kähler term $(S^\dagger S)^2 / \widetilde{\Lambda}^2$ is negative, which stabilize the field S at the origin. Therefore, there is no meta-stable vacuum away from the origin.

Mechanisms of realizing a (local) minimum away from the origin have been discussed in the literature. In Ref. [7], it is shown that the inverted hierarchy mechanism can make a local minimum at a very large value of S if S is non-singlet under some gauge interaction. The possibility of having a (local) minimum via non-perturbative effect has also been discussed in Ref. [8] (See also [2, 9, 10, 11] for the discussion on the vacuum structure of the IYIT model.). The introduction of the bare mass term for the messenger field also makes the vacuum meta-stable

as discussed in Ref. [8, 12]. In the context of supersymmetry breaking in chiral gauge theories, it has been discussed that a runaway direction which is lifted only by non-renormalizable operators has a minimum at a very large field value [13].

In this note, we argue that the situation of supersymmetry restoration by the messenger particles can be cured once we include the supergravity effect even if the origin is the unique minimum in the Planck scale $M_{\rm Pl}$ to infinity limit. Although gravity effect is always suppressed by $1/M_{\rm Pl}$, it can be large enough to stabilize S away from the origin ($S \sim \widetilde{\Lambda}^2/M_{\rm Pl}$) when $\widetilde{\Lambda} \gtrsim 10^{13}$ GeV.

2 Meta-stable vacua in supergravity

The model we will analyze is the following:

$$K = S^\dagger S - \frac{(S^\dagger S)^2}{\widetilde{\Lambda}^2} + q^\dagger q + \bar{q}^\dagger \bar{q} \ , \tag{1}$$

$$W = \mu^2 S - \lambda S q \bar{q} + c \ . \tag{2}$$

The chiral superfield S is a singlet field and q and $\bar{q}$ are the messenger fields which carry the standard model quantum numbers, and λ is a coupling constant. The constant term c does not have any effect if we neglect the gravity interactions, but it is important for the cancellation of the cosmological constant. With neglecting the constant term c, the Lagrangian possesses R-symmetry with the charge assignment $R(S) = 2$, $R(q) = R(\bar{q}) = 0$. We first discuss the model without gravity effect ($M_{\rm Pl} \to \infty$ limit).

For $\lambda = 0$, this model breaks supersymmetry and S is stabilized at the origin, $S = 0$, by the supersymmetry breaking effect. However, by turning on the λ coupling, q and $\bar{q}$ acquire the tachyonic mass terms (B-term) which makes the vacuum unstable. The true minimum is at $S = 0$ and $q = \bar{q} = \sqrt{\mu^2/\lambda}$ where supersymmetry is unbroken. Therefore, there is no supersymmetry breaking vacuum in this model.

We assumed here that the sign of the second term in the Kähler potential is negative as it is the case in the O'Raifeartaigh model. In general, the sign can be positive and sometimes it is uncalculable when the origin of the term is a strong dynamics. If it is positive, the S field may be stabilized away from the origin and the q and $\bar{q}$ directions are also stabilized at $q = \bar{q} = 0$ by the supersymmetric mass terms. However, we consider the case with the negative sign where the origin is a stable minimum.* As we see later, the supersymmetry breaking minimum reappears when we turn on the gravity effect even in that case.

---

*In Ref. [11], the presence of a minimum at the origin is shown in the IYIT model.

Before the analysis in supergravity, we need to estimate the perturbative quantum correction to the Kähler potential for S coming from the interaction term $\lambda S q \bar{q}$, which may be more important than the gravity effect. We can explicitly calculate the term at one-loop level:

$$K_{\text{1-loop}} = -\frac{\lambda^2 N_q}{(4\pi)^2} S^\dagger S \log \frac{S^\dagger S}{\Lambda^2} , \tag{3}$$

where N_q is the number of components in q and $\bar{q}$. For example, $N_q = 5$ if q and $\bar{q}$ transforms as $\mathbf{5}$ and $\bar{\mathbf{5}}$ under $\text{SU}(5)_{\text{GUT}}$. The higher order contributions in the perturbative corrections, including the dependence on the artificial scale Λ, will be minimized by taking the running coupling constant λ to be the value near the scale S. This term also tends to make $S = 0$ stable.

Now we include the gravity effect in this model. The scalar potential of the supergravity Lagrangian is given by

$$V = e^G (G_S G_{S^\dagger} G^{SS^\dagger} + G_q G_{q^\dagger} + G_{\bar{q}} G_{\bar{q}^\dagger} - 3) + \frac{1}{2} D^2 , \tag{4}$$

where $G \equiv K + \log |W|^2$ and we took the unit of $M_{\text{Pl}} = 1$. G_X is the derivative of G with respect to the field X, and $G^{SS^\dagger}$ is the inverse of the Kähler metric. $D^2/2$ represents the D-term contributions. We can easily find the supersymmetric minimum, that is a solution of the equations:

$$G_S = G_q = G_{\bar{q}} = 0, \quad q = \bar{q} . \tag{5}$$

The solution is

$$q = \bar{q} = \sqrt{\frac{\mu^2}{\lambda}} + O\left(\frac{c}{\lambda M_{\text{Pl}}^2}\right) , \quad S = O\left(\frac{c}{\lambda M_{\text{Pl}}^2}\right) . \tag{6}$$

The gravity effect is a slight shift of the values of order $c/(\lambda M_{\text{Pl}}^2)$ which is $O(\mu^2/(\lambda M_{\text{Pl}}))$ if we assume the cancellation of the cosmological constant at the meta-stable supersymmetry breaking vacuum below.

Another minimum can be found with the assumption of $q = \bar{q} = 0$. The potential is simplified to

$$V = e^G (G_S G_{S^\dagger} G^{SS^\dagger} - 3) . \tag{7}$$

The equation $V_S = 0$ with the phenomenological requirement $V = 0$, cancellation of the cosmological constant, leads to

$$(G^{SS^\dagger})_{,S} = -(G^{SS^\dagger})^2 \left[-\frac{\kappa}{S} - \frac{4S^\dagger}{\widetilde{\Lambda}^2} \right] \simeq \frac{2\sqrt{3}}{3 M_{\text{Pl}}} , \tag{8}$$
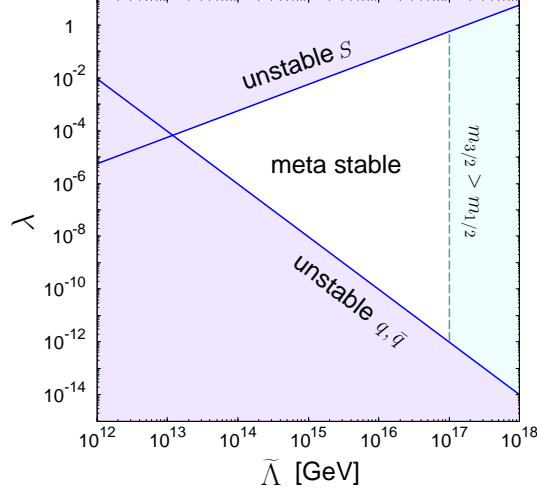
Figure 1: The region of the parameter space where the supersymmetry broken minimum is meta-stable.

where $\kappa = \lambda^2 N_q/(4\pi)^2$. For $\kappa \lesssim (\widetilde{\Lambda}/M_{\mathrm{Pl}})^2$, the minimum is at

$$S \simeq \frac{\sqrt{3}\widetilde{\Lambda}^2}{6M_{\mathrm{Pl}}} \ . \tag{9}$$

Supersymmetry is broken there with $F_S \simeq \mu^2$. By taking the limit $M_{\mathrm{Pl}} \to \infty$, this minimum is moving to $S \to 0$ and the meta-stable vacuum disappears. However, with the finite value of M_{Pl}, the supersymmetry broken and unbroken vacua are at the different places.

In the $\lambda \to 0$ limit, there is no supersymmetric vacuum as in the case without gravity, but the minimum is not at the origin ($S \sim \widetilde{\Lambda}^2/M_{\mathrm{Pl}}$) even though the sign of the $(S^\dagger S)^2$ term in Käher potential is negative. In the supergravity Lagrangian, the origin is no longer a symmetry enhanced point because the R-symmetry is explicitly broken by the c-term. By turning on the λ coupling, the supersymmetric vacuum appears near the origin, but it is separated from the supersymmetry broken minimum. The disappearance of the supersymmetry broken minimum by the finite λ coupling seen before was an artifact of the approximation $M_{\mathrm{Pl}} \to \infty$.

The stability of the vacuum can be checked by looking at the mass matrices of the S, q, and $\bar{q}$ fields. The matrices are given by

$$m_S^2 \simeq \frac{\mu^4}{\widetilde{\Lambda}^2} \cdot \begin{pmatrix} 4 & -6\kappa(M_{\mathrm{Pl}}/\widetilde{\Lambda})^2 \\ -6\kappa(M_{\mathrm{Pl}}/\widetilde{\Lambda})^2 & 4 \end{pmatrix} \ , \quad m_q^2 \simeq \begin{pmatrix} \lambda^2\widetilde{\Lambda}^4/(12M_{\mathrm{Pl}}^2) & -\lambda\mu^2 \\ -\lambda\mu^2 & \lambda^2\widetilde{\Lambda}^4/(12M_{\mathrm{Pl}}^2) \end{pmatrix} \ . \tag{10}$$

Therefore, there is a stable minimum when $\kappa \lesssim (\widetilde{\Lambda}/M_{\mathrm{Pl}})^2$ and $\widetilde{\Lambda}^4 \gtrsim \mu^2 M_{\mathrm{Pl}}^2/\lambda$.

There is another phenomenological requirement that the gauge mediation effect, the gaugino

5

masses, is of order 100 GeV. This fixes the relation between the parameters μ^2 and $\widetilde{\Lambda}$ as follows:

$$\mu^2 \simeq \left(\frac{\alpha}{4\pi}\right)^{-1} \frac{M_W \widetilde{\Lambda}^2}{M_{\text{Pl}}} \; , \tag{11}$$

where M_W is the electroweak scale. With this relation, we show in Fig. 1 the parameter region where the supersymmetry breaking vacuum is meta-stable. The gravitational stabilization of the vacuum is possible for $\widetilde{\Lambda} \gtrsim 10^{13}$ GeV. For large λ, the one-loop correction to the S potential destabilize the vacuum, and too small λ leads to the unstable q and $\bar{q}$ direction because of the small supersymmetric mass terms. The tunneling rate into the supersymmetric vacuum is small enough for $\widetilde{\Lambda} \gtrsim 10^{11}$ GeV which is satisfied in all the viable region [14]. For $\widetilde{\Lambda} \gtrsim 10^{17}$ GeV, the possible gravity mediation effect on the gaugino masses from the operator $SW^\alpha W_\alpha / M_{\text{Pl}}$ is larger than the gauge mediation. Note that this term is forbidden if we impose the approximate R-symmetry discussed before.

The mass of the S field depends on $\widetilde{\Lambda}$. With the relation in Eq. (11), we obtain

$$m_S \simeq 100 \text{ GeV} \left(\frac{\widetilde{\Lambda}}{10^{16} \text{ GeV}}\right) \; . \tag{12}$$

The gravitino mass $m_{3/2}$ is

$$m_{3/2} \simeq 1 \text{ GeV} \left(\frac{\widetilde{\Lambda}}{10^{16} \text{ GeV}}\right)^2 \; . \tag{13}$$

3   Ultraviolet completion

There are several possibilities for the underlying microscopic models which give the effective theory defined in Eqs. (1) and (2). An obvious example is the O'Raifeartaigh model as discussed before. The mass scale of the fields which are integrated out is identified with $\widetilde{\Lambda}$.

There is another interesting possibility that the scale $\widetilde{\Lambda}$ is identified with the dynamical scale of the strongly coupled gauge theory and the linear term $\mu^2 S$ originates from the mass term of the quarks in that theory. This possibility is realized quite simply in the models of Ref. [3]. In the $SO(N_c)$ gauge theory with $N_c - 4$ flavor quarks, there is a branch where the quarks confine and there is no non-perturbatively generated superpotential. If there is a mass term for the quarks T in the N_c dimensional representation, $W = mT^2$, the low energy effective theory is

$$W_{\text{eff}} = mM \; , \tag{14}$$

where M is the meson field $M \sim T^2$. Therefore there is no supersymmetric vacuum in this branch. By assuming the presence of the coupling of the messenger fields q and $\bar{q}$ to the operator

T^2, the effective superpotential is identical to that in Eq. (2) with $\mu^2 \sim m\widetilde{\Lambda}$ and $S \sim M/\widetilde{\Lambda}$. If the coupling is suppressed by the Planck scale, i.e., $W \ni T^2 q\bar{q}/M_{\mathrm{Pl}}$, the λ parameter is of order $\widetilde{\Lambda}/M_{\mathrm{Pl}}$ with which the upper limit on the λ coupling, $\kappa \lesssim (\widetilde{\Lambda}/M_{\mathrm{Pl}})^2$, is always satisfied. The stability of the vacuum is ensured by the fact that the potential grows for large S at the classical level by the mass term of T and the classical analysis is reliable for $S \gtrsim \widetilde{\Lambda}$. If this stabilization is due to the $(S^\dagger S)^2$ term in the Kähler potential, the Kähler potential in Eq. (1) is obtained. For other gauge groups such as $\mathrm{SU}(N_c)$ and $\mathrm{Sp}(N_c)$, it is suggested that there are similar vacua in models with N_c and $N_c + 1$ flavors, respectively.

For other numbers of flavors, it is shown that there are meta-stable supersymmetry broken vacua, for example, in $\mathrm{SU}(N_c)$ with $N_c + 1 \leq N_f < 3N_c/2$ [3]. However, in those cases, the relation between the dynamical scale Λ and the parameters in Eqs. (1) and (2), ($\widetilde{\Lambda}$, μ^2), is $\widetilde{\Lambda} \sim \sqrt{m\Lambda}$ and $\mu^2 \sim m\Lambda$, respectively. With this relation, $\mu^2 \sim \widetilde{\Lambda}^2$, we cannot satisfy the relation in Eq. (11).[†] Exceptions are $\mathrm{SO}(N_c)$ with $N_c - 3$ and $N_c - 2$ flavor theories where there is no non-perturbatively generated superpotential. In the $N_c - 3$ flavor model, there is a branch where low energy effective theory has superpotential [15]:

$$W = f(t)Sa^2 + \mu^2 S . \tag{15}$$

The chiral superfield a is N_f gauge singlet fields and $f(t)$ is an unknown function of $t = (\det S)(Sa^2)$ with $f(0) \neq 0$. We expect that the coupling constant between S and a^2 is $O(1)$ in low energy. In that case, the vacuum with $a = 0$ may be meta-stable for $\widetilde{\Lambda} \gtrsim 10^{16}$ GeV according to Fig. 1. In the $N_c - 2$ flavor model, low energy effective theory is $\mathrm{U}(1)$ gauge theory with a similar superpotential:

$$W = f(t)Sa^+ a^- + \mu^2 S , \tag{16}$$

where a^+ and a^- are the monopoles and $t = \det S$. The same conclusion applies in this case.

For these $N_f = N_c - 3$ and $N_f = N_c - 2$ models, it may be possible that the fields a or $a^\pm$ are actually the messenger fields q and $\bar{q}$ by gauging the subgroup of the flavor symmetry $\mathrm{SU}(N_f)$ and identify it with the standard model gauge group. However, since the meson field S is a symmetric $N_f \times N_f$ matrix which is stabilized only by the supersymmetry breaking effect, i.e., $m_S \simeq \mu^2/\widetilde{\Lambda} \lesssim 10$ TeV, it gives too large contributions to the beta function of the standard model gauge couplings. A larger structure, e.g., introduction of the partner of the unwanted light fields, is necessary for such a scenario to be viable.

[†]Strictly speaking, two scales μ^2 and $\widetilde{\Lambda}^2$ can be separated by assuming a mass hierarchy among the quarks. In that case, $\mu^2 \sim m_L \Lambda$ and $\widetilde{\Lambda}^2 \sim m_H \Lambda$ where m_L and m_H are masses of the light and heavy quarks, respectively. For $m_H \gtrsim \Lambda$, the discussion reduces to the case with fewer flavors.

Particularly interesting scale for $\widetilde{\Lambda}$ is the GUT scale $M_{\mathrm{GUT}} \sim 10^{16}$ GeV. From Eq. (11) and $\mu^2 \sim m\widetilde{\Lambda}$, the mass parameter m is $O(M_W)$ for $\widetilde{\Lambda} \sim M_{\mathrm{GUT}}$. In this case, the parameter m can be related to the μ-parameter in the minimal supersymmetric standard model, which is only the explicit mass parameter in the model.[‡] Indeed, the parameter m can really be the μ-parameter in the scenario where the Higgs fields are composite particles of the strong dynamics which we are discussing [16]. Moreover, the same dynamics can be responsible for the dynamical breaking of the gauge symmetry in grand unified theories as shown in Ref. [16]. In the conventional picture, the electroweak scale appears as a consequence of the supersymmetry breaking, and there was a puzzle that the supersymmetric parameter, the μ-parameter, must be the same size as the supersymmetry breaking parameters. This puzzle was particularly sharp in the scenario of the gauge mediation. However, in this scenario, the electroweak scale is the scale which drives the supersymmetry breaking and therefore there is no coincidence problem. There are many possibilities for the origin of the scale $O(100$ GeV$)$ such as the dynamical scale of another strongly coupled gauge theory. A more attractive possibility of relating it to the cosmological constant term, c-term in Eq. (2), is pointed out in Ref. [16].

The IYIT model also gives the same effective theory in Eqs. (1) and (2). The model is $\mathrm{Sp}(N_c)$ gauge theory with $N_f = N_c + 1$ flavors with superpotential:

$$W = ySQQ , \tag{17}$$

where S is a singlet field and an anti-symmetric $2N_f \times 2N_f$ matrix, Q is the quark in the $2N_c$ dimensional representation, and y is a coupling constant. By the quantum modified constraint, the effective superpotential is

$$W = y\Lambda^2 S , \tag{18}$$

with Λ being the dynamical scale. Therefore $\mu^2 = y\Lambda^2$. Near the origin of S, the correction to the effective Kähler potential is calculated in Ref. [11] to be

$$\delta K \sim -\frac{y^2}{(4\pi)^2}\frac{(S^\dagger S)^2}{\Lambda^2} , \tag{19}$$

which stabilize the origin of the potential. The relation $\widetilde{\Lambda} \simeq 4\pi\Lambda/y$ is obtained. In order to satisfy the relation in Eq. (11), the coupling constant is determined to be $y \simeq 10^{-4}$. With this small value of y, the field S is stabilized at a large value, $S \sim (4\pi)^2\Lambda^2/(y^2 M_{\mathrm{Pl}})$ which must be smaller than Λ/y so that the effective Kähler potential in Eq. (19) is reliable. This constraint gives an upper limit on Λ to be $\Lambda \lesssim 10^{12}$ GeV which translates into $\widetilde{\Lambda} \lesssim M_{\mathrm{Pl}}/10$. This is consistent with the region in Fig. 1.

[‡]The μ-parameter shouldn't be confused with the μ^2 term in Eq. (2). For $\widetilde{\Lambda} \sim M_{\mathrm{GUT}}$, μ^2 is intermediate scale such as $\mu^2 \sim (10^9$ GeV$)^2$ whereas the μ-parameter is always $O(100$ GeV$)$.

4 Summary

We considered the gravitational stabilization mechanism of the supersymmetry broken vacuum in a simple gauge mediation model. The gravitational interaction splits the supersymmetry broken and unbroken vacua for large enough values of the "cut-off" scale $\widetilde{\Lambda} \gtrsim 10^{13}$ GeV.

The low energy effective model we analyzed arises from many microscopic theories of supersymmetry breaking. Therefore the mechanism we discussed is applicable to a wide class of models. The model possesses the R-symmetry and that is unbroken at the origin of the field S. However, the explicit breaking of the R-symmetry in supergravity (by the cosmological constant) shifts the vacuum from the origin. There is no unwanted Goldstone mode associated with the R-symmetry breaking since the symmetry is broken explicitly [17].

Acknowledgments

References

[1] K. I. Izawa and T. Yanagida, "Dynamical Supersymmetry Breaking in Vector-like Gauge Theories," Prog. Theor. Phys. **95**, 829 (1996) [arXiv:hep-th/9602180].

[2] K. A. Intriligator and S. D. Thomas, "Dynamical Supersymmetry Breaking on Quantum Moduli Spaces," Nucl. Phys. B **473**, 121 (1996) [arXiv:hep-th/9603158].

[3] K. Intriligator, N. Seiberg and D. Shih, "Dynamical SUSY breaking in meta-stable vacua," JHEP **0604**, 021 (2006) [arXiv:hep-th/0602239].

[4] M. Dine and A. E. Nelson, "Dynamical supersymmetry breaking at low-energies," Phys. Rev. D **48**, 1277 (1993) [arXiv:hep-ph/9303230]; M. Dine, A. E. Nelson and Y. Shirman, "Low-energy dynamical supersymmetry breaking simplified," Phys. Rev. D **51**, 1362 (1995) [arXiv:hep-ph/9408384]; M. Dine, A. E. Nelson, Y. Nir and Y. Shirman, "New tools for low-energy dynamical supersymmetry breaking," Phys. Rev. D **53**, 2658 (1996) [arXiv:hep-ph/9507378].

[5] M. Dine, W. Fischler and M. Srednicki, Nucl. Phys. B **189**, 575 (1981); S. Dimopoulos and S. Raby, Nucl. Phys. B **192**, 353 (1981); M. Dine and W. Fischler, Phys. Lett. B **110**, 227

(1982); Nucl. Phys. B **204**, 346 (1982); C. R. Nappi and B. A. Ovrut, Phys. Lett. B **113**, 175 (1982); L. Alvarez-Gaume, M. Claudson and M. B. Wise, Nucl. Phys. B **207**, 96 (1982); S. Dimopoulos and S. Raby, Nucl. Phys. B **219**, 479 (1983).

[6] N. Arkani-Hamed, C. D. Carone, L. J. Hall and H. Murayama, Phys. Rev. D **54**, 7032 (1996) [arXiv:hep-ph/9607298]; I. Dasgupta, B. A. Dobrescu and L. Randall, Nucl. Phys. B **483**, 95 (1997) [arXiv:hep-ph/9607487].

[7] H. Murayama, "A model of direct gauge mediation," Phys. Rev. Lett. **79**, 18 (1997) [arXiv:hep-ph/9705271].

[8] K. I. Izawa, Y. Nomura, K. Tobe and T. Yanagida, "Direct-transmission models of dynamical supersymmetry breaking," Phys. Rev. D **56**, 2886 (1997) [arXiv:hep-ph/9705228].

[9] Y. Shirman, "Dynamical supersymmetry breaking versus run-away behavior in supersymmetric gauge theories," Phys. Lett. B **389**, 287 (1996) [arXiv:hep-th/9608147].

[10] N. Arkani-Hamed and H. Murayama, "Renormalization group invariance of exact results in supersymmetric gauge theories," Phys. Rev. D **57**, 6638 (1998) [arXiv:hep-th/9705189].

[11] Z. Chacko, M. A. Luty and E. Ponton, "Calculable dynamical supersymmetry breaking on deformed moduli spaces," JHEP **9812**, 016 (1998) [arXiv:hep-th/9810253].

[12] K. I. Izawa and T. Yanagida, "Strongly coupled gauge mediation," Prog. Theor. Phys. **114**, 433 (2005) [arXiv:hep-ph/0501254].

[13] E. Poppitz and S. P. Trivedi, "New models of gauge and gravity mediated supersymmetry breaking," Phys. Rev. D **55**, 5508 (1997) [arXiv:hep-ph/9609529]; N. Arkani-Hamed, J. March-Russell and H. Murayama, "Building models of gauge-mediated supersymmetry breaking without a messenger sector," Nucl. Phys. B **509**, 3 (1998) [arXiv:hep-ph/9701286].

[14] S. R. Coleman, "The Fate Of The False Vacuum. 1. Semiclassical Theory," Phys. Rev. D **15**, 2929 (1977) [Erratum-ibid. D **16**, 1248 (1977)]; C. G. . Callan and S. R. Coleman, "The Fate Of The False Vacuum. 2. First Quantum Corrections," Phys. Rev. D **16**, 1762 (1977); M. J. Duncan and L. G. Jensen, "Exact tunneling solutions in scalar field theory," Phys. Lett. B **291**, 109 (1992).

[15] K. A. Intriligator and N. Seiberg, "Duality, monopoles, dyons, confinement and oblique confinement in supersymmetric SO(N(c)) gauge theories," Nucl. Phys. B **444**, 125 (1995) [arXiv:hep-th/9503179].

[16] R. Kitano, "Dynamical GUT breaking and μ-term driven supersymmetry breaking," arXiv:hep-ph/0606129.

[17] T. Banks, D. B. Kaplan and A. E. Nelson, "Cosmological implications of dynamical supersymmetry breaking," Phys. Rev. D **49**, 779 (1994) [arXiv:hep-ph/9308292].