

Higgs boson mass of 125 GeV and $g - 2$ of the muon in a gaugino mediation model

Keisuke Harigaya,¹ Tsutomu T. Yanagida,¹ and Norimi Yokozaki²¹*Kavli IPMU (WPI), TODIAS, the University of Tokyo, Kashiwa, Chiba 277-8583, Japan*²*INFN, Sezione di Roma, Piazzale A. Moro 2, I-00185 Roma, Italy*

(Received 26 February 2015; published 14 April 2015)

Gaugino mediation is very attractive since it is free from the serious flavor problem in the supersymmetric standard model. We show that the observed Higgs boson mass at around 125 GeV and the anomaly of the muon $g - 2$ can be easily explained in gaugino mediation models. It should be noted that no dangerous CP violating phases are generated in our framework. Furthermore, there are large parameter regions which can be tested not only at the planned International Linear Collider but also at the coming 13–14 TeV Large Hadron Collider.

DOI: 10.1103/PhysRevD.91.075010

PACS numbers: 12.60.Jv

I. INTRODUCTION

The minimal supersymmetric standard model (MSSM) is an attractive candidate for the physics beyond the standard model (SM). The gauge coupling unification is beautifully realized with new particles predicted by supersymmetry (SUSY). The longstanding anomaly of the muon $g - 2$ [1] (see Refs. [2,3] for SM predictions) is resolved if sleptons, binos, and winos exist at around the weak scale [4,5]. Thus, low-energy SUSY, containing light SUSY particles, is especially attractive. However, there is a serious obstacle in low-energy SUSY models: the SUSY flavor problem. Without the suppression of flavor-violating soft masses, sleptons need to be much heavier than the weak scale, otherwise nonobserved flavor-violating decays, e.g., $\mu \rightarrow e\gamma$ are generated with large branching fractions [6]. Obviously, these heavy sleptons are inconsistent with the explanation of the muon $g - 2$ anomaly. Therefore, we need a SUSY breaking model which largely suppresses flavor-violating soft masses.

Gaugino mediation [7,8] provides a convincing solution to the SUSY flavor problem, where soft masses of squarks and sleptons vanish at a high-energy scale. At low-energy scales, these soft masses are generated radiatively by gaugino loops and, hence, they are flavor blind. As a result, gaugino mediation is free from the serious SUSY flavor problem.¹

In this paper, we show that the observed Higgs boson mass $m_h \approx 125$ GeV [10,11] as well as the anomaly of the muon $g - 2$ can be easily explained in our framework of gaugino mediation, although there is a tension between the muon $g - 2$ and the observed Higgs boson mass in general.

¹Also, gaugino mediation provides an attractive solution to the fine-tuning problem of the electroweak symmetry breaking scale: focus point gaugino mediation [9]. In focus point gaugino mediation, the electroweak symmetry scale is naturally explained by the SUSY particle mass scale even for few TeV gaugino masses.

Here, in our framework, we assume that two Higgs doublets couple to a SUSY breaking field, allowing enhancements of a SUSY contribution to the muon $g - 2$ with a light Higgsino and radiative corrections to the Higgs boson mass through a sizable left-right mixing of the stops. We also discuss the possibility of a large gravitino mass in comparison with MSSM soft masses, which relaxes cosmological problems.

II. GAUGINO MEDIATION MODEL

We consider the Kahler potential such that squark and slepton masses vanish at the tree level. One example is the so-called sequestering [7], which may be realized in the brane world [12]. Another example is the SUSY nonlinear sigma model [13–15] where squarks and sleptons are assumed to be pseudo-Nambu-Goldstone bosons. Due to the Nambu-Goldstone nature of squarks and sleptons, soft masses of them vanish at the tree level. As long as low-energy phenomenology is concerned, the following discussion based on the sequestered Kahler potential is essentially the same as on the Kahler potential in SUSY nonlinear sigma models.² Thus, we discuss the sequestered Kahler potential as an example. We comment on SUSY nonlinear sigma models if necessary.

The Kahler potential and the superpotential are given by

$$K = -3M_P^2 \log \left[1 - \frac{f(Z + Z^\dagger) + Q_i^\dagger Q_i + \Delta K}{3M_P^2} \right],$$

$$W = C + \tilde{\mu} H_u H_d + W_{\text{Yukawas}}, \quad (1)$$

where M_P is the reduced Planck scale, Z is a chiral superfield, f is a real function, Q_i are MSSM chiral

²There would be differences when one discusses quantum corrections to soft masses at a high-energy scale. (See the discussion in Appendix A.)

superfields, \mathcal{C} is a constant, and $\tilde{\mu}$ is the Dirac mass term of Higgsinos. ΔK is explained later. For an appropriate choice of the function f , SUSY is broken by the F term of Z , and the cosmological constant vanishes at the vacuum without introducing a linear term of Z in the superpotential [16]. (See Ref. [17] for the case of $f(Z + Z^\dagger)$ linear in $Z + Z^\dagger$.) A concrete example of f is given in Ref. [18]. In the following, we assume that SUSY is dominantly broken by the F term of Z and call Z the SUSY breaking field. We refer to this type of SUSY breaking as the gravitational SUSY breaking, following Ref. [16].

In addition to the sequestering, we have assumed the shift symmetry, $Z \rightarrow Z + ir$ with a real constant r , so that no CP violating phases arise [18].³ Actually, by phase rotations of MSSM fields and an $U(1)_R$ rotation, we can take $\tilde{\mu}$ and \mathcal{C} as real. Due to the shift symmetry and the reality of the Kahler potential, all couplings in the Kahler potential (and gauge kinetic functions, as we will see,) are real. Thus, physical CP phases vanish in our setup.

We assume that the up- and down-type Higgs couple to the SUSY breaking field through the Kahler potential,⁴

$$\Delta K = \left[c_u \frac{(Z + Z^\dagger)}{M_P} + d_u \frac{(Z + Z^\dagger)^2}{M_P^2} + \dots \right] H_u^\dagger H_u + \left[c_d \frac{(Z + Z^\dagger)}{M_P} + d_d \frac{(Z + Z^\dagger)^2}{M_P^2} + \dots \right] H_d^\dagger H_d, \quad (2)$$

where c_u , c_d , d_u , and d_d are real constants and ellipses denote terms of higher order in $Z + Z^\dagger$. In the sequestering scenario based on the brane world, these couplings can be understood by assuming that Higgs doublets live in bulk.

Hereafter, we shift Z by a constant so that the vacuum expectation value (VEV) of $Z + Z^\dagger$ vanishes and regard Z in ΔK as the shifted field. Then chiral fields with canonical kinetic terms are given by

$$Q_i^c = (1 - \langle f \rangle / 3M_P^2)^{-1/2} Q_i, \quad (3)$$

where $\langle \dots \rangle$ denotes the VEV of \dots . The Higgsino mass is then given by

$$\mu = e^{\langle K \rangle / 2M_P^2} (1 - \langle f \rangle / 3M_P^2) \tilde{\mu}. \quad (4)$$

For simplicity, we omit $\langle \rangle$ in the following.

The scalar soft mass squared for H_u and H_d are given by (see Appendix A)

³This shift symmetry automatically arises in SUSY nonlinear sigma models [19,20].

⁴If the Kahler potential contains a term $K \supset c H_u H_d$, the constant c is, in general, complex after taking $\tilde{\mu}$ real. If $c = O(1)$, this term leads to large CP violations. However, c is suppressed unless the combination $H_u H_d$ has vanishing charges under any symmetries.

$$m_{H_u}^2 = 9k^2(1 - f/3M_P^2)^2[-2d_u + c_u^2] \times m_{3/2}^2, \\ m_{H_d}^2 = 9k^2(1 - f/3M_P^2)^2[-2d_d + c_d^2] \times m_{3/2}^2, \quad (5)$$

where $m_{3/2}$ is the gravitino mass. Higgs trilinear couplings and the B_μ term are

$$A_u = -3k(1 - f/3M_P^2)c_u \times m_{3/2}, \\ A_d = -3k(1 - f/3M_P^2)c_d \times m_{3/2}, \\ B_\mu = (A_u + A_d)\mu, \quad (6)$$

where the constant k is given by

$$k = \left(\frac{\partial f}{\partial x} \right)^{-1}. \quad (7)$$

In SUSY nonlinear sigma models, those soft masses, in general, exist if Higgs doublets are not Nambu-Goldstone bosons.

Note that the B_μ term is proportional to the sum of trilinear couplings and vanishes when MSSM Higgses are also sequestered from the SUSY breaking field (i.e., $c_u = d_u = c_d = d_d = 0$). This is not the case for generic sequestering scenarios because of the VEV of the scalar auxiliary component of the supergravity multiplet. (In the conformal formulation, it is the VEV of the F term of the compensator.) In the gravitational SUSY breaking, the VEV vanishes at the vacuum [16] and, hence, the B_μ term vanishes in the sequestered limit.

Next, we consider gaugino masses. As a model of the grand unified theory (GUT), we consider the $SU(5) \times SU(3)_H \times U(1)_H$ product group unification (PGU) model [21,22]. In the PGU model, the doublet-triplet splitting problem, which is a serious problem in the minimal $SU(5)$ GUT, is solved. The gauge coupling unification is approximately maintained if the gauge coupling of $SU(3)_H \times U(1)_H$ is sufficiently stronger than that of $SU(5)$.

The relevant part of the Lagrangian is given by

$$\mathcal{L} \supset \int d^2\theta \left[\left(\frac{1}{4g_5^2} - \frac{k_5 Z}{M_P} \right) W_5 W_5 + \left(\frac{1}{4g_{3H}^2} - \frac{k_{3H} Z}{M_P} \right) W_{3H} W_{3H} + \left(\frac{1}{4g_{1H}^2} - \frac{k_{1H} Z}{M_P} \right) W_{1H} W_{1H} \right] + \text{H.c.} \quad (8)$$

Here, g_5 , g_{3H} , and g_{1H} are the gauge coupling constants of $SU(5)$, $SU(3)_H$ and $U(1)_H$ gauge interactions, respectively. W_5 , W_{3H} , and W_{1H} are the superfield field strength of each gauge multiplet. k_5 , k_{3H} , and k_{1H} are constants, which are real in order to preserve the shift symmetry of Z .⁵

⁵The shift symmetry, $Z \rightarrow Z + ir$, has a quantum gauge anomaly and, hence, the shift symmetry is maintained in a perturbative limit.

After $SU(5) \times SU(3)_H \times U(1)_H$ is broken to $SU(3)_C \times SU(2)_L \times U(1)_Y$, nonuniversal gaugino masses are generated at the GUT scale [23],

$$\begin{aligned} M_1/M_2 &= \frac{k_5 \mathcal{N} + k_{1H}}{k_5} \frac{g_{1H}^2}{g_5^2 + \mathcal{N} g_{1H}^2}, \\ M_3/M_2 &= \frac{k_5 + k_{3H}}{k_5} \frac{g_{3H}^2}{g_5^2 + g_{3H}^2}, \end{aligned} \quad (9)$$

where the real constant \mathcal{N} depends on the $U(1)_H$ charge of the GUT breaking field. In the strong coupling limit, $g_{1H}^2, g_{3H}^2 \gg 1$, the ratios of the gaugino masses are written in a simple form: $M_1/M_2 \simeq (k_5 \mathcal{N} + k_{1H})/(k_5 \mathcal{N})$ and $M_3/M_2 \simeq (k_5 + k_{3H})/k_5$.

Together with Eqs. (5) and (6), we have obtained a gaugino mediation model with nonzero soft masses of the Higgs sector and nonuniversal gaugino masses.

Before closing this section, we discuss how large gravitino mass can be naturally taken in comparison with MSSM soft masses. It is known that the gravitino as well as the SUSY breaking field cause various cosmological problems because they are easily produced in the early Universe while they are long-lived [24–27]. If the gravitino mass is large, these problems are relaxed because the gravitino and the SUSY breaking field are shorter-lived for larger masses. It would be worthwhile to consider the possibility of a large gravitino mass in comparison with MSSM soft masses.

In order to evaluate the naturalness of a large gravitino mass, let us assume that couplings of the SUSY breaking field Z with Higgs doublets (ΔK), as well as that with gauginos (k_5, k_{3H} and k_{1H}), are also absent and, hence, all soft masses vanish at the tree level.⁶ We estimate possible quantum corrections to soft masses under this assumption and require that the corrections do not upset our setup.

When MSSM soft masses vanish at the tree level, the MSSM is a supersymmetric theory, where no soft masses are generated.⁷ Thus, quantum corrections to MSSM soft masses arise only from the so-called anomaly mediation [12,28–31] or from gravitational interactions [32,33].⁸

The possible largest quantum correction is the one-loop correction to scalar soft mass squared from gravitational interactions. One-loop quantum corrections by

⁶Note that the B_μ term also vanishes in this limit in the gravitational SUSY breaking, as we have discussed. Thus, there is no so-called B_μ problem even in the sequestering limit.

⁷This argument can be invalidated by couplings between regulator fields and the SUSY breaking field Z . See the comment in Appendix A.

⁸The anomaly mediation proportional to the VEV of the scalar auxiliary component of the supergravity multiplet is determined by the superdiffeomorphism [31] and, hence, cannot be eliminated. In our setup, the VEV vanishes at the tree level and, hence, that anomaly mediation is suppressed [16].

gravitational interactions around the cutoff scale are expected to generate scalar soft mass squared as large as

$$\Delta m_{\text{scalar}}^2 \sim \frac{1}{16\pi^2} m_{3/2}^2 \left(\frac{\Lambda}{M_P} \right)^n, \quad (10)$$

where Λ and n are the cutoff of the theory and an integer, respectively. If $\Lambda \sim M_P$, the gravitino mass is at most $O(1)$ TeV for soft masses of $O(100)$ GeV. If $\Lambda \ll M_P$, however, as is assumed in the sequestering based on the brane world [12], the gravitino mass can be larger.

For some cases, one-loop corrections by gravitational interactions to scalar soft mass squared vanish [34], as is shown in Appendix A. (There, we also discuss SUSY nonlinear sigma models.) In that case, quantum corrections to scalar soft mass squared start from the two-loop level. Even if $\Lambda \sim M_P$, the two-loop corrections are, at most, comparable to possible one-loop corrections to gaugino masses, trilinear couplings, and the B_μ term. Thus, the gravitino mass of $O(10)$ TeV is possible in this case, even if $\Lambda \sim M_P$.

III. HIGGS BOSON MASS AND $g-2$ OF THE MUON

In the MSSM, the observed Higgs boson mass at around 125 GeV is explained by a large stop mass and/or a large trilinear coupling of the stops [35]. These soft masses generate large radiative corrections to the Higgs potential: $\Delta m_{H_u}^2 \sim (m_t^2 \text{ or } A_t^2)$, where m_t and A_t are the stop mass and a trilinear coupling, respectively. For $m_t, A_t \sim 1\text{--}4$ TeV, the fine-tuning of parameters in the Higgs potential is required to explain the electroweak symmetry breaking (EWSB) scale. From EWSB conditions, Z boson mass m_Z is written as

$$\begin{aligned} \frac{m_Z^2}{2} &\simeq -(m_{H_u}^2(M_{\text{GUT}}) + \Delta m_{H_u}^2 + \mu^2) \\ &\quad + (m_{H_d}^2(M_{\text{GUT}}) + \Delta m_{H_d}^2 - m_{H_u}^2(M_{\text{GUT}}) \\ &\quad - \Delta m_{H_u}^2)/\tan^2\beta + \dots, \end{aligned} \quad (11)$$

where ... indicates terms suppressed by $1/\tan^n\beta$ ($n \geq 4$). Here, $m_{H_u}(M_{\text{GUT}})$ and $m_{H_d}(M_{\text{GUT}})$ denote the soft masses of the up- and down-type Higgs at the GUT scale, respectively, and $\Delta m_{H_u}^2$ and $\Delta m_{H_d}^2$ are the radiative corrections to $m_{H_u}^2$ and $m_{H_d}^2$. To explain $m_Z \simeq 91.2$ GeV, $\Delta m_{H_u}^2 \sim (m_t^2 \text{ or } A_t^2)$ needs to be canceled by either $m_{H_u}^2(M_{\text{GUT}})$ or μ^2 .⁹ Consequently, there arise two distinct regions: a small μ region and a large μ region. If the Higgs

⁹There is an exception where no large cancellation between $m_{H_u}^2(M_{\text{GUT}})$, $\Delta m_{H_u}^2$ and μ^2 is required. For instance, in focus point gaugino mediation models [9], $\Delta m_{H_u}^2$ is small and the fine-tuning is significantly relaxed.

potential is tuned by the SUSY breaking parameter, $m_{H_u}^2(M_{\text{GUT}})$, it is likely that μ is of the order of m_Z . In contrast, if the EWSB scale is explained by the tuning of the SUSY invariant mass μ , μ needs to be as large as $\mu \sim m_{\tilde{t}}$.

A. Small μ case

When the Higgsino is light (μ is small), the SUSY contributions to the muon $g - 2$ are dominated by the wino-Higgsino-(muon-sneutrino) loop. In this case, the dominant SUSY contribution is given by [5]

$$(a_\mu)_{\tilde{W}-\tilde{H}-\tilde{\nu}_L} \simeq \frac{\alpha_2 m_\mu^2 (M_2 \mu)}{4\pi m_{\tilde{\nu}}^4} \tan \beta \cdot F_C \left(\frac{M_2^2}{m_{\tilde{\nu}}^2}, \frac{\mu^2}{m_{\tilde{\nu}}^2} \right), \quad (12)$$

where F_C is a loop function (e.g., $F_C(1, 1) = 1/2$). It is larger than the contribution from loops involving the bino and smuons [see Eq. (14)]. Therefore, the small μ allows relatively heavier electroweakinos and sleptons to explain the muon $g - 2$ than the large μ .

In Fig. 1, contours of the Higgs boson mass and the region consistent with the muon $g - 2$ are shown. The mass of the CP-even Higgs boson is calculated by using

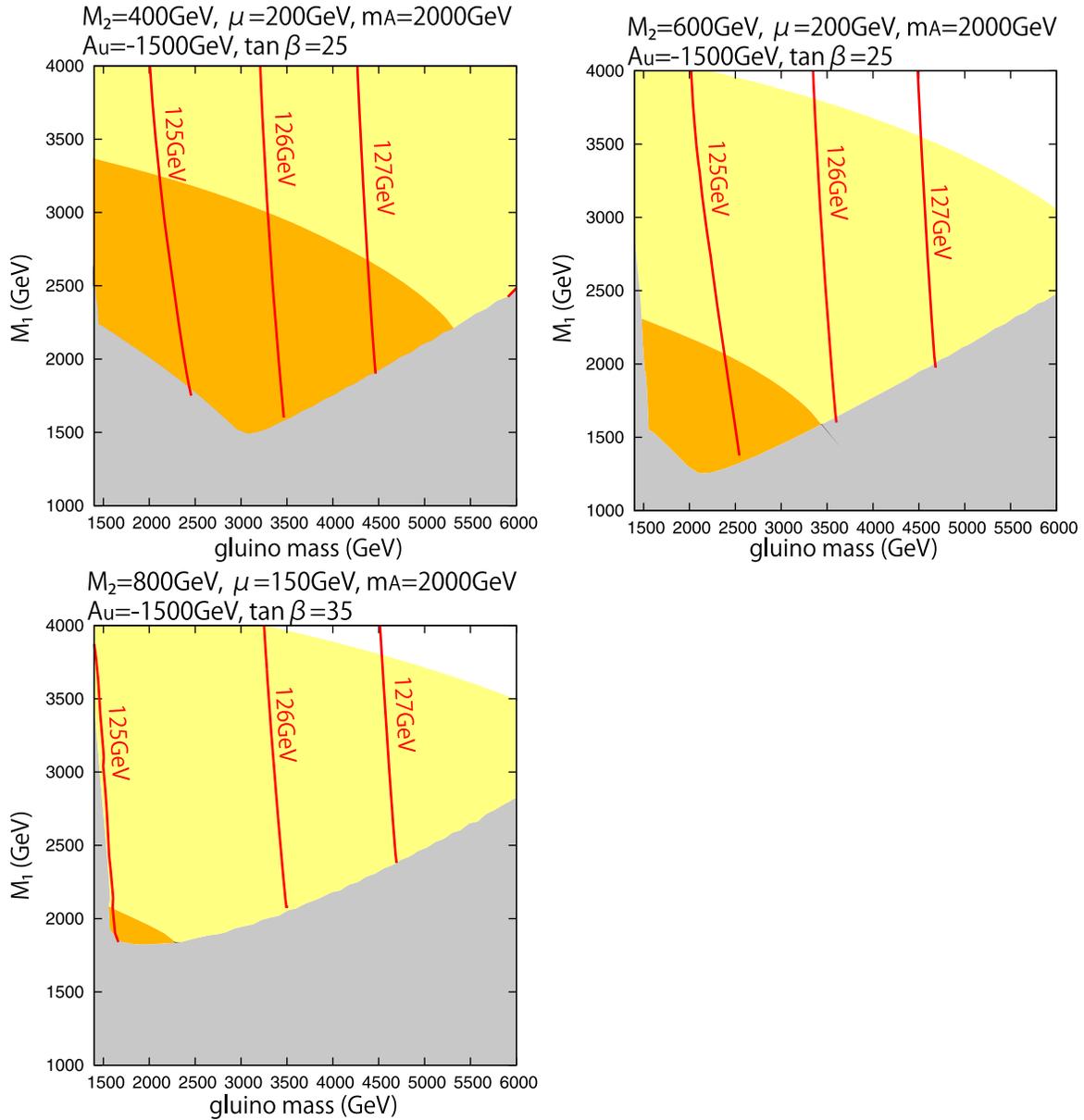


FIG. 1 (color online). The SUSY contribution to the muon $g - 2$ and the Higgs boson mass (red solid line) in the small μ case. In the orange (yellow) region, the muon $g - 2$ is explained at $1(2) \sigma$ level. In the top panels (the bottom panel), $\mu = 200$ GeV (150 GeV) and $\tan \beta = 20$ (35). Here, $\alpha_S(m_Z) = 0.1184$ and $m_t = 173.3$ GeV.

FEYNHIGGS2.10.2 [36] with the option resumming large logs. The SUSY contributions to the muon $g - 2$, Δa_μ , are also evaluated by FEYNHIGGS. The SUSY mass spectrum is evaluated by using the SUSPECT package [37] with modification suitable for our purpose. In the top panels (the bottom panel), B , $m_{H_u}^2$, and $m_{H_d}^2$ at the GUT scale are taken such that $\tan \beta = 25$ (35), $\mu = 200$ GeV (150 GeV), and the physical mass of the CP-odd Higgs boson, $m_A = 2000$ GeV, are reproduced. The trilinear couplings are taken as $A_u = -1500$ GeV, $A_d = A_e = (B_\mu/\mu) - A_u$ [see Eq. (6)]. The negative sign of A_u at the GUT scale is taken so that contributions from gaugino loops are added to A_u constructively. We discard the gray region, since the stau becomes the lightest SUSY particle (LSP) or the stop is lighter than 600 GeV [38]¹⁰; in the region of small M_1 , the stau becomes light because of a small positive radiative correction from M_1 . The region consistent with the muon $g - 2$ at 1σ (2σ) level is shown in orange (yellow). Here, the deviation of the muon $g - 2$ is evaluated as

$$(a_\mu)_{\text{EXP}} - (a_\mu)_{\text{SM}} = (26.1 \pm 8.1) \cdot 10^{-10}, \quad (13)$$

using the SM prediction in Ref. [2].

One can see that the observed Higgs boson mass and muon $g - 2$ are explained simultaneously for the physical gluino mass at around 2.5–3.0 TeV and $M_2(M_{\text{GUT}}) = 400\text{--}800$ GeV (corresponding to $M_2(m_{\text{SUSY}}) \simeq 290\text{--}620$ GeV) (see also Table I). Since squark masses are smaller than the physical gluino mass, the region consistent with the observed Higgs boson mass and the muon $g - 2$ can be tested at the 14 TeV Large Hadron Collider (LHC), through productions of the gluino and squarks [39,40]. Note that the abundance of the lightest neutralino, which is mainly composed of the Higgsino with a mixture of the wino, is much smaller than the observed value, $\Omega_{\text{DM}} h^2 \simeq 0.12$ [41,42], and hence, another candidate for dark matter is required.¹¹

In the region consistent with the muon $g - 2$, the wino has a small mass and hence the constraint from charino/neutralino searches at the LHC should be considered. The charged and neutral winos produced by the electroweak interactions decay into Higgsinos + (W , Z , h), since left-handed sleptons are heavier than the wino. This process can be examined by the chargino/neutralino searches in the final state with two or three leptons and missing transverse momentum [44,45]. So far, the constraint is not very severe, and it is difficult to give a bound on the wino mass for

¹⁰In this region with the light stop, the constraint from the inclusive $b \rightarrow s\gamma$ decay is more stringent than the LHC stop searches.

¹¹Since the spin-independent scattering cross section of this Higgsino-wino neutralino is large as a few $\times 10^{-9}$ pb, the present model is excluded [43], if the neutralino is a dominant component of dark matter.

TABLE I. The mass spectra with the small μ . Here, $\Delta\text{Br}(b \rightarrow s\gamma) = \text{Br}(b \rightarrow s\gamma)_{\text{MSSM}} - \text{Br}(b \rightarrow s\gamma)_{\text{SM}}$.

P1	
$M_1(M_{\text{GUT}})$	2200 GeV
$M_2(M_{\text{GUT}})$	400 GeV
$M_3(M_{\text{GUT}})$	1100 GeV
$A_u(M_{\text{GUT}})$	-1300 GeV
$\tan \beta$	20
μ	200 GeV
m_A	2 TeV
m_{gluino}	2.4 TeV
$m_{\tilde{q}}$	2.1 TeV
$m_{\tilde{t}_{1,2}}$	1.4, 1.8 TeV
$m_{\tilde{e}_L}(m_{\tilde{\mu}_L})$	450 GeV
$m_{\tilde{e}_R}(m_{\tilde{\mu}_R})$	836 GeV
$m_{\tilde{\tau}_1}$	361 GeV
$m_{\chi_1^0}, m_{\chi_2^0}$	179, 210 GeV
$m_{\chi_3^0}, m_{\chi_4^0}$	342, 935 GeV
$m_{\chi_1^\pm}, m_{\chi_2^\pm}$	184, 343 GeV
m_h	124.5 GeV
Δa_μ	2.20×10^{-9}
$\Delta\text{Br}(b \rightarrow s\gamma)$	-2.9×10^{-5}
P2	
$M_1(M_{\text{GUT}})$	1900 GeV
$M_2(M_{\text{GUT}})$	740 GeV
$M_3(M_{\text{GUT}})$	1400 GeV
$A_u(M_{\text{GUT}})$	-1300 GeV
$\tan \beta$	35
μ	150 GeV
m_A	2 TeV
m_{gluino}	3.0 TeV
$m_{\tilde{q}}$	2.6 TeV
$m_{\tilde{t}_{1,2}}$	1.8, 2.2 TeV
$m_{\tilde{e}_L}(m_{\tilde{\mu}_L})$	573 GeV
$m_{\tilde{e}_R}(m_{\tilde{\mu}_R})$	721 GeV
$m_{\tilde{\tau}_1}$	174 GeV
$m_{\chi_1^0}, m_{\chi_2^0}$	145, 159 GeV
$m_{\chi_3^0}, m_{\chi_4^0}$	602, 806 GeV
$m_{\chi_1^\pm}, m_{\chi_2^\pm}$	151, 602 GeV
m_h	125.3 GeV
Δa_μ	1.97×10^{-9}
$\Delta\text{Br}(b \rightarrow s\gamma)$	-2.5×10^{-5}

$\mu > 150$ GeV; the region consistent with the muon $g - 2$ is safe. At the 14 TeV LHC, the wino mass up to around 800 GeV can be excluded (discovered) for an integrated luminosity of 300(3000) fb^{-1} [40]; therefore, it is expected

that the region consistent with the muon $g - 2$ at 1σ level is tested at the 14 TeV LHC.

B. Large μ case

If the Higgsino is heavy, the bino-(L-smuon)-(R-smuon) loop dominates the SUSY contributions to the muon $g - 2$. The contribution from this loop is given by [5]

$$(a_\mu)_{\tilde{B}-\tilde{\mu}_L-\tilde{\mu}_R} \simeq \frac{3}{5} \frac{\alpha_1}{4\pi} \frac{m_\mu^2}{M_1^3} \tan\beta \cdot F_N\left(\frac{m_{\tilde{\mu}_R}^2}{M_1^2}, \frac{m_{\tilde{\mu}_L}^2}{M_1^2}\right), \quad (14)$$

where F_N is a loop function (e.g., $F_N(1, 1) = 1/6$). In this region, a small bino mass is required to enhance $(a_\mu)_{\tilde{B}-\tilde{\mu}_L-\tilde{\mu}_R}$. As a result, the stau tends to be tachyonic in a gaugino mediation model. This tachyonic stau is avoided by the positive contribution from $(m_{H_d}^2 - m_{H_u}^2)$ through the renormalization group evolution. Indeed, the one-loop renormalization group equation for the right handed slepton is given by

$$\frac{dm_{E^c}^2}{dt} \ni \frac{1}{16\pi^2} \left[-\frac{6}{5} g_1^2 (m_{H_d}^2 - m_{H_u}^2) \right]. \quad (15)$$

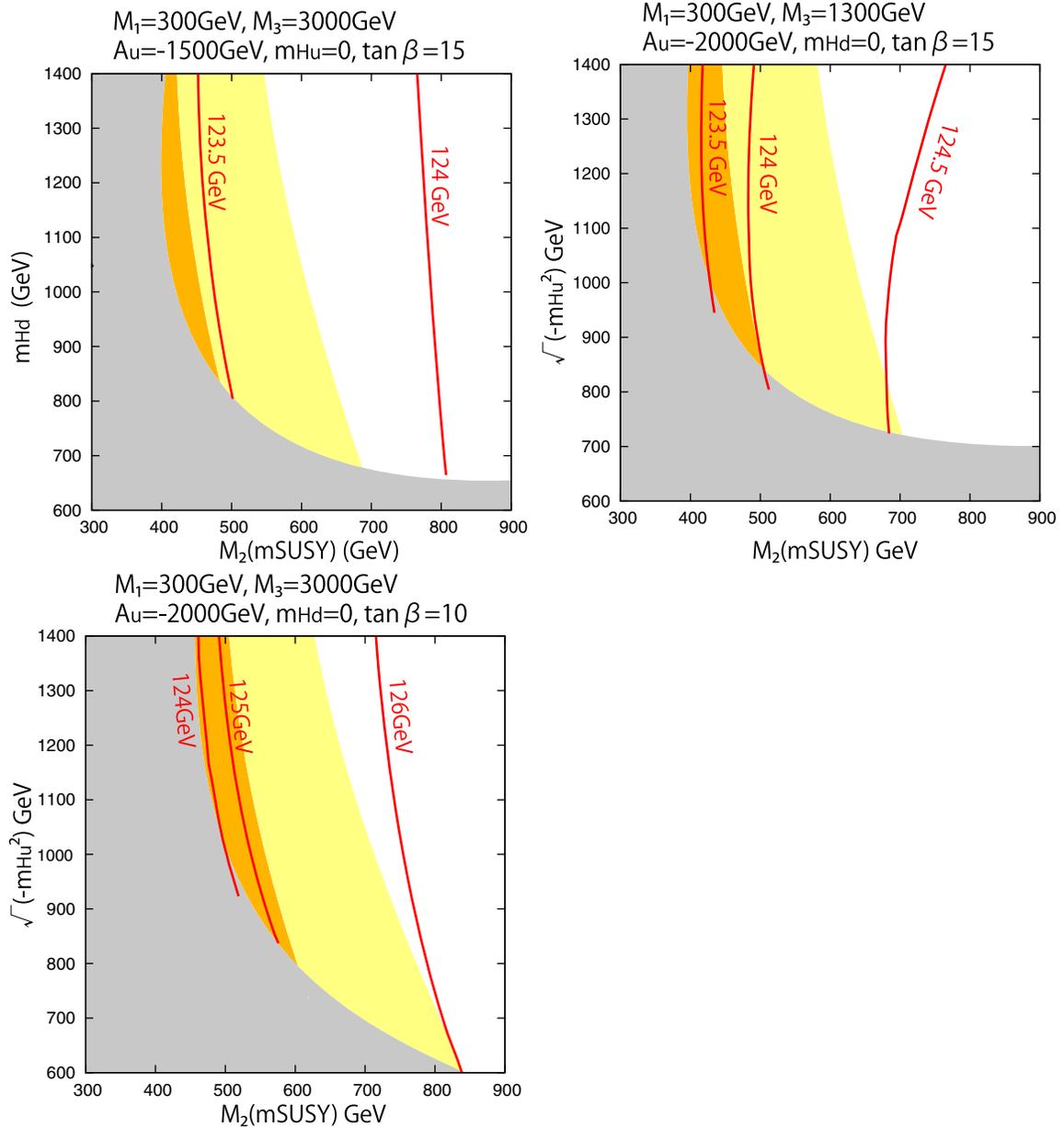


FIG. 2 (color online). The SUSY contribution to the muon $g - 2$ and the Higgs boson mass in the large μ case. The different values of M_3 , $M_3 = 1300$ GeV (top panels) and 3000 GeV (bottom panel) are taken. In the orange (yellow) region, the muon $g - 2$ is explained at $1(2)\sigma$ level. In the top-left panel, $m_{H_u} = 0$, while in the other panels, $m_{H_d} = 0$.

This contribution vanishes if $m_{H_u}^2 = m_{H_d}^2$. However, it is sizable when $m_{H_d}^2 \gg m_{H_u}^2$. In our gaugino mediation model, H_u and H_d couple to the SUSY breaking field Z ; therefore, $(m_{H_d}^2 - m_{H_u}^2)$ can be positive at the high-energy scale resulting in the positive mass squared of the stau.

In Fig. 2, we show Δa_μ and the contours of m_h in the large μ cases. In the top-left panel, $m_{H_d}^2(M_{\text{GUT}}) > 0$ and $m_{H_u}^2 = 0$ are taken, and in other two panels, $m_{H_d}^2(M_{\text{GUT}}) = 0$ and $m_{H_u}^2(M_{\text{GUT}}) < 0$ are taken. The trilinear coupling A_u is $A_u = -1500$ GeV in the top-left panel, and $A_u = -2000$ GeV in the other two panels. Other trilinear couplings satisfy the condition in Eq. (6), $A_d = A_e = (B_\mu/\mu) - A_u$. The gray region is excluded since the stau becomes the lightest SUSY particle. On the edge of the stau LSP region, i.e., for $m_{\chi_1^0} \sim m_{\tilde{\tau}_1}$, the relic abundance of the neutralino is consistent with the observed value, $\Omega_{\text{DM}} h^2 \simeq 0.12$ [41,42]. This is because the coannihilation [46] with the lighter stau reduces the relic abundance of the lightest neutralino efficiently.

It is shown that the muon $g-2$ is explained at 1σ level for $M_2(m_{\text{SUSY}}) < 500(600)$ GeV and the gluino mass of 2.8 (6.1) TeV. The calculated Higgs boson mass can be consistent with the observed mass at around 125 GeV by taking into account the uncertainties from the top pole mass and theoretical calculation of the Higgs boson mass. However, as shown below, this region is rather severely constrained from the chargino/neutralino searches at the LHC. Here, $M_3 = (1300, 3000)$ GeV corresponds to $m_{\text{gluino}} = (2.8, 6.1)$ TeV and $m_{\text{squark}} = (2.4, 5.2)$ TeV, where m_{gluino} and m_{squark} are the physical gluino mass and squark mass, respectively.

In contrast to the small μ case, the region consistent with the muon $g-2$ in the large μ case is rather severely constrained by chargino/neutralino searches at the LHC. This is because the left-handed sleptons are lighter than the wino in this region, and hence the wino can decay into an on-shell slepton and a lepton, with the slepton decaying into a neutralino and a lepton. As a result, this region is severely constrained by searches for the electroweak production of the chargino and neutralino in a final state with three leptons and missing transverse momentum [45,47]. From these searches, the wino mass is constrained to be larger than 600–700 GeV, depending on the slepton mass. Considering this constraint, the muon $g-2$ is explained at $1.5-2\sigma$ ($1-1.5\sigma$) level for $M_3 = 1300(3000)$ GeV (see also Table II). For larger M_3 , e.g., $M_3 = 5200$ GeV, the muon $g-2$ is explained at 1σ level with a chargino mass larger than 700 GeV (see P6 in Table II). Note that a large gluino mass indirectly enhances $(a_\mu)_{\tilde{B}-\tilde{\mu}_L-\tilde{\mu}_R}$ for fixed $\tan\beta$, bino, and smuon masses. Large M_3 generates large $|m_{H_u}^2|$ through radiative corrections at the two-loop level. With this large $|m_{H_u}^2|$, μ is determined to

 TABLE II. The mass spectra with the large μ . Here, $\mu \simeq m_{\chi_2^\pm}$.

P3	
$M_1(M_{\text{GUT}})$	300 GeV
$M_2(M_{\text{GUT}})$	760 GeV
$M_3(M_{\text{GUT}})$	1300 GeV
$A_u(M_{\text{GUT}})$	-2000 GeV
$\tan\beta$	15
$m_{H_u}^2(M_{\text{GUT}})$	$-6 \cdot 10^5$ GeV ²
$m_{H_d}^2(M_{\text{GUT}})$	0
m_{gluino}	2.8 TeV
$m_{\tilde{q}}$	2.5 TeV
$m_{\tilde{\tau}_{1,2}}$	1.9, 2.2 TeV
$m_{\tilde{e}_L}(m_{\tilde{\mu}_L})$	471 GeV
$m_{\tilde{e}_R}(m_{\tilde{\mu}_R})$	212 GeV
$m_{\tilde{\tau}_1}$	120 GeV
$m_{\chi_1^0}, m_{\chi_2^0}$	118, 609 GeV
$m_{\chi_1^\pm}, m_{\chi_2^\pm}$	609, 2006 GeV
m_h	124.3 GeV
Δa_μ	1.40×10^{-9}
P4	
$M_1(M_{\text{GUT}})$	300 GeV
$M_2(M_{\text{GUT}})$	780 GeV
$M_3(M_{\text{GUT}})$	3000 GeV
$A_u(M_{\text{GUT}})$	-2000 GeV
$\tan\beta$	10
$m_{H_u}^2(M_{\text{GUT}})$	$-7.6 \cdot 10^5$ GeV ²
$m_{H_d}^2(M_{\text{GUT}})$	0
m_{gluino}	6.1 TeV
$m_{\tilde{q}}$	5.2 TeV
$m_{\tilde{\tau}_{1,2}}$	4.4, 4.9 TeV
$m_{\tilde{e}_L}(m_{\tilde{\mu}_L})$	423 GeV
$m_{\tilde{e}_R}(m_{\tilde{\mu}_R})$	218 GeV
$m_{\tilde{\tau}_1}$	118 GeV
$m_{\chi_1^0}, m_{\chi_2^0}$	107, 606 GeV
$m_{\chi_1^\pm}, m_{\chi_2^\pm}$	606, 3671 GeV
m_h	125.2 GeV
Δa_μ	1.88×10^{-9}
P5	
$M_1(M_{\text{GUT}})$	300 GeV
$M_2(M_{\text{GUT}})$	900 GeV
$M_3(M_{\text{GUT}})$	3000 GeV
$A_u(M_{\text{GUT}})$	-2000 GeV
$\tan\beta$	10
$m_{H_u}^2(M_{\text{GUT}})$	0

(Table continued)

TABLE II. (Continued)

P5	
$m_{H_d}^2(M_{\text{GUT}})$	$6 \cdot 10^5 \text{ GeV}^2$
m_{gluino}	6.1 TeV
$m_{\tilde{q}}$	5.2 TeV
$m_{\tilde{t}_{1,2}}$	4.4, 4.9 TeV
$m_{\tilde{e}_L}(m_{\tilde{\mu}_L})$	515 GeV
$m_{\tilde{e}_R}(m_{\tilde{\mu}_R})$	203 GeV
$m_{\tilde{\tau}_1}$	113 GeV
$m_{\chi_1^0}, m_{\chi_2^0}$	107, 707 GeV
$m_{\chi_1^\pm}, m_{\chi_2^\pm}$	707, 3602 GeV
m_h	125.7 GeV
Δa_μ	1.41×10^{-9}
P6	
$M_1(M_{\text{GUT}})$	300 GeV
$M_2(M_{\text{GUT}})$	940 GeV
$M_3(M_{\text{GUT}})$	5200 GeV
$A_u(M_{\text{GUT}})$	-2000 GeV
$\tan \beta$	7
$m_{H_u}^2(M_{\text{GUT}})$	$-1.2 \cdot 10^6 \text{ GeV}^2$
$m_{H_d}^2(M_{\text{GUT}})$	0
m_{gluino}	10.3 TeV
$m_{\tilde{q}}$	8.7 TeV
$m_{\tilde{t}_{1,2}}$	7.4, 8.1 TeV
$m_{\tilde{e}_L}(m_{\tilde{\mu}_L})$	419 GeV
$m_{\tilde{e}_R}(m_{\tilde{\mu}_R})$	231 GeV
$m_{\tilde{\tau}_1}$	104 GeV
$m_{\chi_1^0}, m_{\chi_2^0}$	92, 708 GeV
$m_{\chi_1^\pm}, m_{\chi_2^\pm}$	708, 5828 GeV
m_h	124.8 GeV
Δa_μ	1.86×10^{-9}

be large from the EWSB condition in Eq. (11). Consequently, $(a_\mu)_{\tilde{B}-\tilde{\mu}_L-\tilde{\mu}_R}$ proportional to μ is enhanced.¹² This is the reason why favored gluino masses are relatively larger than those in the small μ case.

Finally, some mass spectra for small μ and large μ are shown in Table I and II. In small μ cases, the branching ratio of the inclusive $b \rightarrow s\gamma$ decay is enhanced for large A_t ; we calculate $\Delta\text{Br}(b \rightarrow s\gamma) = \text{Br}(b \rightarrow s\gamma)_{\text{MSSM}} - \text{Br}(b \rightarrow s\gamma)_{\text{SM}}$ in P1–P2 by using SUPERISO package [49], and demand that $-3.6 \cdot 10^{-5} <$

¹²Large $\mu \tan \beta$ generates to a charge breaking minimum in the Higgs-stau potential, which can be deeper than the EWSB minimum. Therefore, the size of $\mu \tan \beta$ is constrained from above by the stability of the EWSB minimum [48].

$\Delta\text{Br}(b \rightarrow s\gamma) < 9.2 \cdot 10^{-5}$. Here, the required range of $\Delta\text{Br}(b \rightarrow s\gamma)$ is given by the difference between the experimental value of $\text{Br}(b \rightarrow s\gamma)$ [50] and SM prediction [51] with an inclusion of 2σ error.

IV. CONCLUSION AND DISCUSSION

We have shown that the observed Higgs boson mass at around 125 GeV and the anomaly of the muon $g-2$ are explained simultaneously in our gaugino mediation models. There is no SUSY CP problem thanks to the shift symmetry of the SUSY breaking field Z , and the gravitational SUSY breaking mechanism. The Higgs doublets are assumed to couple to Z , giving nonzero soft masses of the Higgs doublets and the trilinear coupling of the stops. With this trilinear coupling, the Higgs boson mass of 125 GeV is explained relatively easily: colored SUSY particles can be lighter than 3 TeV, and they are expected to be produced at the 14 TeV LHC. Thanks to the nonzero soft masses of the Higgs doublets, the μ parameter can be small if the Higgs potential is tuned by these soft masses. In this case, the muon $g-2$ is explained at 1σ level avoiding the constraints from the chargino/neutralino searches. With the small μ , light Higgsinos as well as sleptons are targets of searches at lepton colliders such as the International Linear Collider experiments.

The muon $g-2$ is consistent with the large μ case as well. In the large μ case, the Higgs potential is tuned by this μ parameter rather than the soft masses of the Higgs doublets. Compared to the small μ case, a relatively heavy gluino is favored for the muon $g-2$, and the gluino and squarks are too heavy to be tested even at the 14 TeV LHC. However, this region can be easily covered through the chargino/slepton searches at the LHC [18].

Let us comment on possible problems concerning cosmological aspects: the gravitino problem and the Polonyi problem. By allowing small couplings between Z and Higgs doublets, the gravitino mass can be larger than $O(10)$ TeV. In this case, the cosmological gravitino problem is significantly relaxed, compared to the case with the gravitino mass of $O(1)$ TeV. The Polonyi problem caused by Z can be solved by the adiabatic solution provided that Z strongly couples to an inflaton [52,53]. Or maybe the Polonyi problem is simply absent due to an anthropic reason; note that the constraint from the big bang nucleosynthesis, which would not be avoided by an anthropic reason, is avoided for a sufficiently large mass of Z , $m_Z > O(10)$ TeV because Z decays before the big bang nucleosynthesis starts.

ACKNOWLEDGMENTS

We thank Sho Iwamoto for useful discussion. This work is supported by Grants-in-Aid for Scientific Research from the Ministry of Education, Science, Sports, and Culture (MEXT), Japan, No. 26104009 and

No. 26287039 (T. T. Y.) and also by the World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan (K. H. and T. T. Y.). The research leading to these results has received funding from the European Research Council under the European Union's Seventh Framework Programme (FP/2007-2013)/ERC Grant Agreement No. 279972 "NPFlavour" (N. Y.). The work of K. H. is supported in part by a JSPS Research Fellowship for Young Scientists.

APPENDIX: SOFT SCALAR MASSES

In this appendix, we calculate the soft scalar masses of the MSSM particles shown in Sec. II. For simplicity, we put $M_P = 1$ in this appendix.

1. Tree-level soft scalar masses

In general, the scalar potential is given by the Kahler potential $K(X^I, X^{\dagger J})$ and the superpotential $W(X^I)$ as¹³

$$V = e^K [K^{I\bar{J}} (W_I + K_I W) (W_{\bar{J}} + K_{\bar{J}} W)^\dagger - 3|W|^2]. \quad (\text{A1})$$

Here, I, J, \dots and \bar{I}, \bar{J}, \dots indicate chiral and antichiral fields, respectively. Lower indices denote derivatives with respect to the corresponding field. $K^{I\bar{J}}$ is the inverse of the matrix $K_{I\bar{J}}$.

Let us first consider the following Kahler potential and the superpotential,

$$K = g(Z + Z^\dagger) + h_i(Z + Z^\dagger) Q_i^\dagger Q_i, \quad (\text{A2})$$

$$W = \mathcal{C} + \tilde{y} Q_1 Q_2 Q_3 + \tilde{\mu} Q_4 Q_5, \quad (\text{A3})$$

where $Q_i (i = 1-5)$, \tilde{y} and $\tilde{\mu}$ are matter chiral fields, the Yukawa coupling, and the mass term, respectively. g and h_i

are real functions. Note that fields Q_i are not canonically normalized. Canonically normalized fields are given by

$$Q_i^c = h_i^{1/2} Q_i. \quad (\text{A4})$$

Then the Yukawa coupling and the mass term for the canonically normalized fields are

$$y = e^{g/2} (h_1 h_2 h_3)^{-1/2} \tilde{y}, \quad \mu = e^{g/2} (h_4 h_5)^{-1/2} \tilde{\mu}. \quad (\text{A5})$$

The potential of Z is given by

$$V(Z, Z^\dagger) = e^g |\mathcal{C}|^2 [g'^{-1} g^2 - 3]. \quad (\text{A6})$$

Vanishing of the cosmological constant requires that

$$g'^2 = 3g'' \quad (\text{A7})$$

at the vacuum.

Scalar soft masses are given by

$$\begin{aligned} V_{\text{soft}} &= m_i^2 |Q_i^c|^2 + [y A Q_1^c Q_2^c Q_3^c + B_\mu Q_4^c Q_5^c + \text{H.c.}], \\ m_i^2 &= [1 - 9(g')^{-2} (\ln h_i)'''] \times m_{3/2}^2, \\ A &= \sum_{i=1,2,3} \left(1 - \frac{3}{g'} (\ln h_i)' \right) \times m_{3/2}, \\ B_\mu/\mu &= \sum_{i=4,5} \left(1 - \frac{3}{g'} (\ln h_i)' \right) \times m_{3/2}, \end{aligned} \quad (\text{A8})$$

where primes denote the derivative with respect to $Z + Z^\dagger$. Here, we have used $e^{K/2} \mathcal{C}^* = m_{3/2}$ and Eq. (A7).

The Kahler potential discussed in Sec. II corresponds to the case with

$$\begin{aligned} g &= -3 \ln(1 - f/3), \\ h_i &= \begin{cases} (1 - f/3)^{-1} (1 + c_u(Z + Z^\dagger) + d_u(Z + Z^\dagger)^2) & \text{for } Q_i = H_u, \\ (1 - f/3)^{-1} (1 + c_d(Z + Z^\dagger) + d_d(Z + Z^\dagger)^2) & \text{for } Q_i = H_d, \\ (1 - f/3)^{-1} & \text{for others.} \end{cases} \end{aligned} \quad (\text{A9})$$

From Eqs. (A8) and (A9), we obtain the soft masses shown in Eqs. (5) and (6). Scalar soft mass squared of squarks and sleptons vanish.

2. Quantum corrections to soft scalar mass squared

We are interested in the case where the gravitino mass is far larger than the soft masses of MSSM because the

¹³We omit the D -potential, which is irrelevant for our discussion.

gravitino problem and the Polonyi problem are relaxed in this case. Let us assume vanishing soft masses at the tree level [in the sequestering discussed in the main text, $\Delta K = 0$ in Eq. (2) and $k_5 = k_{3H} = k_{1H} = 0$ in Eq. (8)] and discuss how large quantum corrections to soft masses are expected.

As we have discussed in the main text, when MSSM soft masses vanish at the tree level, renormalizable interactions in MSSM do not generate soft masses. Then quantum corrections to MSSM soft masses originate only from the

anomaly mediation or from higher-dimensional interactions such as gravitational interactions. Possibly the largest correction is the one-loop correction to scalar soft mass squared from higher-dimensional interactions, which we investigate here.

The MSSM field couples to the SUSY breaking field and the gravitino through higher-dimensional interactions. These interactions are expected to generate MSSM soft masses through quantum corrections. We treat the quantum corrections by introducing a cutoff Λ to the theory and regard the action with vanishing MSSM soft masses as a Wilsonian action at the cutoff scale. Then quantum corrections to the scalar soft mass squared are given by the following form,

$$\Delta m_{\text{scalar}}^2 \sim \frac{1}{(16\pi^2)^m} m_{3/2}^2 \left(\frac{\Lambda}{M_*}\right)^n, \quad (\text{A10})$$

where M_* is the suppression scale of higher-dimensional interactions such as the Planck scale, m is the number of loops, and n is an integer.

In the following, we show that one-loop corrections to the soft mass squared are absent when the tree-level Kahler potential is either of the following forms:

$$K = K(Z + Z^\dagger + Q_i^\dagger Q_i), \quad (\text{A11})$$

$$K = K(\ln(1 + Q_i^\dagger Q_i) + Z + Z^\dagger). \quad (\text{A12})$$

The former is the case with the sequestered form in Eq. (1) with $f(Z + Z^\dagger)$ linear in $Z + Z^\dagger$, the so-called no-scale structure.¹⁴ The latter is the case with the SUSY $CP^N \simeq SU(N+1)/SU(N) \times U(1)$ nonlinear sigma model, where N is the number of chiral superfields Q_i . The proof is parallel to the one given in Ref. [34].

In the limit of vanishing Yukawa couplings and gauge couplings, the action is invariant under the following transformation of superfields,

$$\delta Q_i = \epsilon_i, \quad \delta Z = \epsilon_i^* Q_i: \quad \text{for Eq. (A11)}, \quad (\text{A13})$$

$$\begin{aligned} \delta Q_i &= -\frac{i}{2}\theta_{i1}(1 - Q_i^2) + \frac{1}{2}\theta_{i2}(1 + Q_i^2) - i\theta_{i3}Q_i, \\ \delta Z &= -\frac{i}{2}\theta_{i1}Q_i - \frac{1}{2}\theta_{i2}Q_i + \frac{i}{2}\theta_{i3}: \quad \text{for Eq. (A12)}, \end{aligned} \quad (\text{A14})$$

where ϵ_i and θ_i are complex and real infinitesimal parameters of transformations. The latter transformation is nothing but $SU(2)$ subgroups of $SU(N+1)$.

¹⁴For the no-scale case, the tree-level potential of Z also vanishes. The potential is given by one-loop corrections and hence the scalar component of Z is lighter than the gravitino.

In both cases, quantum corrections without Yukawa or gauge coupling constants can modify the Kahler potential only to the following form consistent with the symmetries,

$$K = F(Z + Z^\dagger + Q_i^\dagger Q_i) + \dots, \quad (\text{A15})$$

where F is some real function and \dots denote terms higher order in Q_i .

Then, scalar soft mass squared appear in the scalar potential only through the following combination consistent with the symmetries,

$$V(Z + Z^\dagger + Q_i^\dagger Q_i), \quad (\text{A16})$$

which give the potential of Z for $Q_i = 0$. The scalar soft mass squared of Q_i is given by

$$\begin{aligned} \frac{\partial}{\partial Q_i^\dagger} \frac{\partial}{\partial Q_i} V(Z + Z^\dagger + Q_i^\dagger Q_i)|_{Q_i=0} \\ = V'(Z + Z^\dagger), \end{aligned} \quad (\text{A17})$$

which vanishes at the vacuum.

We have shown that the scalar soft mass squared is not generated from quantum corrections without Yukawa or gauge coupling constants. Also, as we have discussed, neither the renormalizable Yukawa or gauge interactions alone generate soft masses. Thus, the scalar soft mass squared is generated only by corrections involving both the Yukawa/gauge coupling constants and higher-dimensional interactions. Such corrections are absent at the one-loop level and are possible only from the two-loop level. Thus, quantum corrections to scalar soft mass squared are at most as large as

$$\Delta m_{\text{scalar}}^2 \sim \frac{\lambda^2}{(16\pi^2)^2} m_{3/2}^2 \left(\frac{\Lambda}{M_*}\right)^n, \quad (\text{A18})$$

where λ is the Yukawa or gauge coupling constants.

Finally, we comment on how the above discussion may be invalidated by the regularization of quantum corrections. We have argued that when MSSM soft masses vanish at the tree level, renormalizable interactions in MSSM do not generate soft masses. This argument may be invalidated by the regularization. Consider, for example, the Pauli-Villars regularization. If Pauli-Villars fields have tree-level soft masses, loop corrections from Pauli-Villars fields generate MSSM soft masses at the one-loop level [54,55]. For the above discussion to be valid, tree-level soft masses of the Pauli-Villars fields must also be absent. In the sequestering based on the brane world, this assumption holds if the Pauli-Villars fields also live on the brane of the visible sector.

- [1] G. W. Bennett *et al.* (Muon G-2 Collaboration), *Phys. Rev. D* **73**, 072003 (2006).
- [2] K. Hagiwara, R. Liao, A. D. Martin, D. Nomura, and T. Teubner, *J. Phys. G* **38**, 085003 (2011).
- [3] M. Davier, A. Hoecker, B. Malaescu, and Z. Zhang, *Eur. Phys. J. C* **71**, 1515 (2011); **72**, 1874(E) (2012).
- [4] J. L. Lopez, D. V. Nanopoulos, and X. Wang, *Phys. Rev. D* **49**, 366 (1994); U. Chattopadhyay and P. Nath, *Phys. Rev. D* **53**, 1648 (1996).
- [5] T. Moroi, *Phys. Rev. D* **53**, 6565 (1996); **56**, 4424(E) (1997).
- [6] L. J. Hall, V. A. Kostelecky, and S. Raby, *Nucl. Phys.* **B267**, 415 (1986); F. Gabbiani and A. Masiero, *Phys. Lett. B* **209**, 289 (1988); R. Barbieri and L. J. Hall, *Phys. Lett. B* **338**, 212 (1994); R. Barbieri, L. J. Hall, and A. Strumia, *Nucl. Phys.* **B445**, 219 (1995); J. Hisano, T. Moroi, K. Tobe, M. Yamaguchi, and T. Yanagida, *Phys. Lett. B* **357**, 579 (1995).
- [7] K. Inoue, M. Kawasaki, M. Yamaguchi, and T. Yanagida, *Phys. Rev. D* **45**, 328 (1992).
- [8] D. E. Kaplan, G. D. Kribs, and M. Schmaltz, *Phys. Rev. D* **62**, 035010 (2000); Z. Chacko, M. A. Luty, A. E. Nelson, and E. Ponton, *J. High Energy Phys.* 01 (2000) 003.
- [9] T. T. Yanagida and N. Yokozaki, *Phys. Lett. B* **722**, 355 (2013); *J. High Energy Phys.* 11 (2013) 020; 10 (2014) 133.
- [10] G. Aad *et al.* (ATLAS Collaboration), *Phys. Lett. B* **716**, 1 (2012); *Phys. Rev. D* **90**, 052004 (2014).
- [11] S. Chatrchyan *et al.* (CMS Collaboration), *Phys. Lett. B* **716**, 30 (2012); *Phys. Rev. D* **89**, 092007 (2014); V. Khachatryan *et al.* (CMS Collaboration), *Eur. Phys. J. C* **74**, 3076 (2014).
- [12] L. Randall and R. Sundrum, *Nucl. Phys.* **B557**, 79 (1999).
- [13] W. A. Bardeen and V. Visnjic, *Nucl. Phys.* **B194**, 422 (1982).
- [14] W. Buchmuller, S. T. Love, R. D. Peccei, and T. Yanagida, *Phys. Lett.* **115B**, 233 (1982).
- [15] W. Buchmuller, R. D. Peccei, and T. Yanagida, *Phys. Lett.* **124B**, 67 (1983).
- [16] K.-I. Izawa, T. Kugo, and T. T. Yanagida, *Prog. Theor. Phys.* **125**, 261 (2011).
- [17] E. Cremmer, S. Ferrara, C. Kounnas, and D. V. Nanopoulos, *Phys. Lett.* **133B**, 61 (1983).
- [18] S. Iwamoto, T. T. Yanagida, and N. Yokozaki, *arXiv:1407.4226*.
- [19] Z. Komargodski and N. Seiberg, *J. High Energy Phys.* 07 (2010) 017.
- [20] T. Kugo and T. T. Yanagida, *Prog. Theor. Phys.* **124**, 555 (2010).
- [21] T. Yanagida, *Phys. Lett. B* **344**, 211 (1995).
- [22] T. Hotta, K. I. Izawa, and T. Yanagida, *Phys. Rev. D* **53**, 3913 (1996).
- [23] N. Arkani-Hamed, H. C. Cheng, and T. Moroi, *Phys. Lett. B* **387**, 529 (1996).
- [24] H. Pagels and J. R. Primack, *Phys. Rev. Lett.* **48**, 223 (1982).
- [25] S. Weinberg, *Phys. Rev. Lett.* **48**, 1303 (1982).
- [26] M. Y. Khlopov and A. D. Linde, *Phys. Lett.* **138B**, 265 (1984).
- [27] G. D. Coughlan, W. Fischler, E. W. Kolb, S. Raby, and G. G. Ross, *Phys. Lett.* **131B**, 59 (1983).
- [28] G. F. Giudice, M. A. Luty, H. Murayama, and R. Rattazzi, *J. High Energy Phys.* 12 (1998) 027.
- [29] J. A. Bagger, T. Moroi, and E. Poppitz, *J. High Energy Phys.* 04 (2000) 009.
- [30] F. D'Eramo, J. Thaler, and Z. Thomas, *J. High Energy Phys.* 06 (2012) 151.
- [31] K. Harigaya and M. Ibe, *Phys. Rev. D* **90**, 085028 (2014).
- [32] I. Antoniadis and M. Quiros, *Nucl. Phys.* **B505**, 109 (1997).
- [33] M. A. Luty, *Phys. Rev. Lett.* **89**, 141801 (2002).
- [34] P. Binetruy and M. K. Gaillard, *Phys. Lett. B* **195**, 382 (1987).
- [35] Y. Okada, M. Yamaguchi, and T. Yanagida, *Prog. Theor. Phys.* **85**, 1 (1991); *Phys. Lett. B* **262**, 54 (1991); J. Ellis, G. Ridolfi and F. Zwirner, *Phys. Lett. B* **257**, 83 (1991); H. E. Haber and R. Hempfling, *Phys. Rev. Lett.* **66**, 1815 (1991); J. Ellis, G. Ridolfi and F. Zwirner, *Phys. Lett. B* **262**, 477 (1991).
- [36] S. Heinemeyer, W. Hollik, and G. Weiglein, *Comput. Phys. Commun.* **124**, 76 (2000); *Eur. Phys. J. C* **9**, 343 (1999); G. Degrassi, S. Heinemeyer, W. Hollik, P. Slavich, and G. Weiglein, *Eur. Phys. J. C* **28**, 133 (2003); M. Frank, T. Hahn, S. Heinemeyer, W. Hollik, H. Rzehak, and G. Weiglein, *J. High Energy Phys.* 02 (2007) 047.
- [37] A. Djouadi, J.-L. Kneur and G. Moultaka, *Comput. Phys. Commun.* **176**, 426 (2007).
- [38] G. Aad *et al.* (ATLAS Collaboration), *J. High Energy Phys.* 09 (2014) 015; *J. High Energy Phys.* 11 (2014) 118.
- [39] ATLAS Collaboration, Report No. ATL-PHYS-PUB-2013-002.
- [40] ATLAS Collaboration, Report No. ATL-PHYS-PUB-2014-010.
- [41] G. Hinshaw *et al.* (WMAP Collaboration), *Astrophys. J. Suppl. Ser.* **208**, 19 (2013).
- [42] P. A. R. Ade *et al.* (Planck Collaboration), *Astron. Astrophys.* **571**, A16 (2014).
- [43] D. S. Akerib *et al.* (LUX Collaboration), *Phys. Rev. Lett.* **112**, 091303 (2014).
- [44] G. Aad *et al.* (ATLAS Collaboration), *J. High Energy Phys.* 05 (2014) 071.
- [45] V. Khachatryan *et al.* (CMS Collaboration), *Eur. Phys. J. C* **74**, 3036 (2014).
- [46] K. Griest and D. Seckel, *Phys. Rev. D* **43**, 3191 (1991).
- [47] G. Aad *et al.* (ATLAS Collaboration), *J. High Energy Phys.* 04 (2014) 169.
- [48] R. Rattazzi and U. Sarid, *Nucl. Phys.* **B501**, 297 (1997); J. Hisano and S. Sugiyama, *Phys. Lett. B* **696**, 92 (2011); R. Sato, K. Tobioka, and N. Yokozaki, *Phys. Lett. B* **716**, 441 (2012).
- [49] F. Mahmoudi, *Comput. Phys. Commun.* **178**, 745 (2008); *Comput. Phys. Commun.* **180**, 1579 (2009).
- [50] Heavy Flavor Averaging Group, <http://www.slac.stanford.edu/xorg/hfag>.
- [51] M. Misiak, H. M. Asatrian, K. Bieri, M. Czakon, A. Czarnecki, T. Ewerth, A. Ferroglia, and P. Gambino *et al.*, *Phys. Rev. Lett.* **98**, 022002 (2007).
- [52] A. D. Linde, *Phys. Rev. D* **53**, R4129(R) (1996).
- [53] K. Nakayama, F. Takahashi, and T. T. Yanagida, *Phys. Rev. D* **84**, 123523 (2011).
- [54] M. K. Gaillard and B. D. Nelson, *Nucl. Phys.* **B588**, 197 (2000); P. Binetruy, M. K. Gaillard, and B. D. Nelson, *Nucl. Phys.* **B604**, 32 (2001).
- [55] J. L. Evans, M. Ibe, K. A. Olive, and T. T. Yanagida, *Eur. Phys. J. C* **74**, 2775 (2014).