Supersymmetric standard models with a pseudo-Dirac gluino from hybrid F- and D-term supersymmetry breaking

Ran Ding,[1] Tianjun Li,[2,3] Florian Staub,[4] Chi Tian,[3] and Bin Zhu[2,5]

[1]*Center for High-Energy Physics, Peking University, Beijing 100871, People's Republic of China*
[2]*State Key Laboratory of Theoretical Physics and Kavli Institute for Theoretical Physics China (KITPC), Institute of Theoretical Physics, Chinese Academy of Sciences, Beijing 100190, People's Republic of China*
[3]*School of Physical Electronics, University of Electronic Science and Technology of China, Chengdu 610054, People's Republic of China*
[4]*Theory Division, CERN, 1211 Geneva 23, Switzerland*
[5]*Institute of Physics, Chinese Academy of Sciences, Beijing 100190, People's Republic of China*

(Received 11 March 2015; published 13 July 2015)

We propose the supersymmetric Standard Models (SSMs) with a pseudo-Dirac gluino from hybrid F- and D-term supersymmetry (SUSY) breaking. Similar to the SSMs before the LHC, all the supersymmetric particles in the minimal SSM obtain the SUSY breaking soft terms from the traditional gravity mediation and have masses within about 1 TeV except gluino. To evade the LHC SUSY search constraints, the gluino also has a heavy Dirac mass above 3 TeV from D-term SUSY breaking. Interestingly, such a heavy Dirac gluino mass will not induce the electroweak fine-tuning problem. We realize such SUSY breaking via an anomalous $U(1)_X$ gauge symmetry inspired from string models. To maintain the gauge coupling unification and increase the Higgs boson mass, we introduce extra vectorlike particles. We study the viable parameter space which satisfies all the current experimental constraints and present a concrete benchmark point. This kind of model not only preserves the merits of pre-LHC SSMs such as naturalness, dark matter, etc., but also solves the possible problems in the SSMs with Dirac gauginos due to the F-term gravity mediation.

DOI: 10.1103/PhysRevD.92.015008 PACS numbers: 12.60.Jv, 12.60.Fr, 14.80.Ly

I. INTRODUCTION

It is well known that the weak scale supersymmetry (SUSY) is the most promising extension for physics beyond the Standard Model (SM) [1]. It provides a well-motivated and complete framework to understand the basic questions of TeV-scale physics: the gauge hierarchy problem is solved naturally, the lightest supersymmetric particle (LSP) such as neutralino can be a dark matter candidate, and gauge coupling unification can be realized, etc. The gauge coupling unification strongly suggests the grand unified theories (GUTs), and only the superstring theory may describe the real world. Thus, the supersymmetric SM (SSM) is also a bridge between the low-energy phenomenology and high-energy fundamental physics.

However, the discovered SM-like Higgs boson with a mass around 125 GeV [2,3] is a little bit too heavy in the minimal SSM (MSSM) since it requires the multi-TeV top squarks with small mixing or TeV-scale top squarks with large mixing [4]. Also, there exist strong constraints on the SSMs from the LHC SUSY searches. For example, the gluino mass $m_{\tilde{g}}$ and first two-generation squark mass $m_{\tilde{q}}$ should be heavier than about 1.7 TeV if they are roughly degenerate $m_{\tilde{q}} \sim m_{\tilde{g}}$, and the squark mass $m_{\tilde{q}}$ is heavier than about 850 GeV for $m_{\tilde{g}} \gg m_{\tilde{q}}$ [5]. Therefore, the naturalness of the SSMs is challenged.

The basic idea to lift Higgs mass without threatening the hierarchy problem is the introduction of additional tree-level contributions [6–12]. To escape the LHC SUSY search constraints, there are quite a few proposals: natural SUSY [13,14], compressed SUSY [15–17], stealth SUSY [18], heavy LSP SUSY [19], R-parity violation [20,21], supersoft SUSY [22–31], etc. Here, we would like to point out that all the sparticles in the SSMs can be within about 1 TeV as long as the gluino is heavier than 3 TeV, which is obviously a simple modification to the SSMs before the LHC. Also, such a heavy gluino will not induce the electroweak fine-tuning problem if it is (pseudo-)Dirac like the supersoft SUSY. However, there exist some problems for supersoft SUSY with Dirac gauginos: μ problem cannot be solved via the Giudice-Masiero mechanism [32], the D-term contribution to the Higgs quartic coupling vanishes, the right-handed slepton may be the LSP, and the scalar components of the adjoint chiral superfields might be tachyonic and then break the SM gauge symmetry, etc. [22]. The first three problems can be solved in the F-term gravity mediation, while the last problem was solved recently [31]. Therefore, we will propose the SSMs with a pseudo-Dirac gluino from hybrid

F- and D-term SUSY breaking. To be concrete, all the sparticles in the MSSM obtain SUSY breaking soft terms from the traditional gravity mediation, and only gluino receives extra Dirac mass from the D-term SUSY breaking. Especially, all the MSSM sparticles except gluino can be within about 1 TeV as the pre-LHC SSMs. The merits of this proposal are keeping the good properties of pre-LHC SSMs (naturalness, as well as explanations for the dark matter and muon anomalous magnetic moment, etc.), evading the LHC SUSY search constraints, and solving the problems in supersoft SUSY via F-term gravity mediation. We show that such SUSY breaking can be realized by an anomalous $U(1)_X$ gauge symmetry inspired from string models. To achieve the gauge coupling unification and increase the Higgs boson mass, we will introduce vectorlike particles. We shall discuss the low-energy phenomenology, and the detailed studies will be given elsewhere [33].

II. MODEL BUILDING

In order to generate the Dirac gluino mass, a chiral superfield Φ in the adjoint representation of $SU(3)_C$ is needed. To maintain the gauge coupling unification and lift the Higgs boson mass, we need to introduce some extra vectorlike particles. To avoid the Landau pole for the SM gauge couplings below the GUT scale, we only have two kinds of models: $\Delta b = 3$ and $\Delta b = 4$, where Δb is the uniform contribution to the one-loop beta functions of the SM gauge couplings from all the new particles. The additional vectorlike particles and their quantum numbers in the supersymmetric SMs with $\Delta b = 3$ and $\Delta b = 4$ are given in Tables I and II, respectively. We will study the model with $\Delta b = 3$ elsewhere (for Dirac gaugino case, see Ref. [29]). Here, we shall consider the model with $\Delta b = 4$. In this model, the $SU(2)_L \times U(1)_Y$ Dirac gaugino masses are forbidden, and the neutrino masses and mixings can be generated via type II seesaw mechanism [34]. From the string model building point of view, we usually do not have the vectorlike particles T_+ and T_- since they arise from a symmetric **15** representation of $SU(5)$. Interestingly, the symmetric **15** representation of $SU(5)$ or flipped $SU(5)$ can indeed be obtained in the type IIA orientifold on $\mathbf{T^6}/(\mathbf{Z_2} \times \mathbf{Z_2})$ with intersecting D6-branes [35,36]. The alternative way to get T_+ and T_- is to embed $SU(2)_L$ into a diagonal gauge group of $SU(2)_A \times SU(2)_B$, which was done in a particular $Z_3 \times Z_3$ orbifold of the heterotic string [37]. In this case, the type II seesaw mechanism can be realized as well [37].

TABLE I. The extra vectorlike particles and their quantum numbers in the supersymmetric SM with $\Delta b = 3$. Here, $i = 1, 2$, and we do not have to introduce S except for Dirac gaugino case since it is a SM singlet.

Particles	Quantum numbers	Particles	Quantum numbers
Φ	$(\mathbf{8}, \mathbf{1}, \mathbf{0})$	T	$(\mathbf{1}, \mathbf{3}, \mathbf{0})$
XL	$(\mathbf{1}, \mathbf{2}, \mathbf{-1/2})$	XL^c	$(\mathbf{1}, \mathbf{2}, \mathbf{1/2})$
XE_i	$(\mathbf{1}, \mathbf{1}, \mathbf{-1})$	XE_i^c	$(\mathbf{1}, \mathbf{1}, \mathbf{1})$
S	$(\mathbf{1}, \mathbf{1}, \mathbf{0})$		

TABLE II. The extra vectorlike particles and their quantum numbers in the supersymmetric SM with $\Delta b = 4$.

Particles	Quantum numbers	Particles	Quantum numbers
Φ	$(\mathbf{8}, \mathbf{1}, \mathbf{0})$		
XD	$(\mathbf{3}, \mathbf{1}, \mathbf{-1/3})$	XD^c	$(\bar{\mathbf{3}}, \mathbf{1}, \mathbf{1/3})$
T_+	$(\mathbf{1}, \mathbf{3}, \mathbf{1})$	T_-	$(\mathbf{1}, \mathbf{3}, \mathbf{-1})$

Comparing to the MSSM, the new superpotential terms with universal vectorlike particle mass are

$$W = M_V(T_+T_- + XD^cXD) + \lambda H_u T_- H_u + \lambda' H_d T_+ H_d,$$

where H_d and H_u are the MSSM Higgs fields. The λ and λ' terms will, respectively, give positive and negative contributions to the Higgs boson mass via the nondecoupling effects. Although with both terms we can still get the Higgs boson with mass around 125 GeV easily, to simplify the discussions we shall neglect the λ' term in the following. The corresponding SUSY breaking soft terms are

$$\begin{aligned} -\mathcal{L}_{\rm soft} = {} & B_T T_- T_+ + B_D XD^c XD + T_\lambda H_u T_- H_u \\ & + M_D G\Phi + \text{H.c.} + \tilde{\phi}^\dagger m_{\tilde{\phi}}^2 \tilde{\phi}, \end{aligned} \tag{1}$$

where $B_{\mu,T,D}$ are bilinear soft terms, $m_{\tilde{\phi}}^2$ are soft scalar masses, and M_D is the Dirac gluino mass.

III. SUSY BREAKING

To realize the hybrid F- and D-term SUSY breaking, we shall consider the anomalous $U(1)_X$ gauge symmetry inspired from string models [38]. Unlike Ref. [38], we introduce two SM singlet fields S and S' with $U(1)_X$ charges 0 and -1 and assume that all the SM particles and vectorlike particles are neutral under $U(1)_X$. In general, there could exist other exotic particles Q_i^X with $U(1)_X$ charges q_i^X from any real string compactification. Thus, the $U(1)_X$ D-term potential is

$$V_D = \frac{g_X^2}{2} D^2 = \frac{g_X^2}{2}\left(\sum\nolimits_i q_i^X |Q_i^X|^2 - |S'|^2 + \xi\right)^2, \tag{2}$$

where for example, in the heterotic string compactification [39], the Fayet-Iliopoulos term is given by

$$\xi = \frac{g_X^2 \mathrm{Tr} q^X}{384\pi^2} M_{\rm Pl}^2, \tag{3}$$

where $M_{\rm Pl}$ is the reduced Planck scale.

To achieve gravity mediation, we consider the following superpotential from the instanton effect which breaks $U(1)_X$:

$$W_{\text{instanton}} = M_I S S'. \tag{4}$$

This is the key difference between our scenario and that in Ref. [38] where the superpotential is $U(1)_X$ invariant and then one cannot realize the traditional gravity mediation. Minimizing the potential, we obtain

$$\langle S \rangle = 0, \qquad \langle S' \rangle^2 = \xi - M_I^2/g_X^2, \qquad \langle F_{S'} \rangle = 0, \tag{5}$$

$$\langle F_S \rangle = M_I \sqrt{\xi - M_I^2/g_X^2}, \qquad \langle D \rangle = M_I^2/g_X^2. \tag{6}$$

Because S is neutral under $U(1)_X$, the traditional gravity mediation can be realized via the nonzero F_S. The Dirac mass for gluino/Φ and soft scalar masses for Φ and $T_{+/-}$ can be generated, respectively, via the following operators [31]:

$$\int d^2\theta \left(\frac{\bar{D}^2 D^\alpha V'}{M_*} W_{3,\alpha} \Phi + \frac{\bar{D}^2 (D^\alpha V' D_\alpha X')}{M_*} X'' \right), \tag{7}$$

where we neglect the coefficients for simplicity, X' and X'' can both be Φ as well as, respectively, be $T_{+/-}$ and $T_{-/+}$, and M_* can be the reduced Planck scale for gravity mediation or the effective messenger scale. In addition, like the above second kind of operators, we can have the similar operators for H_d/H_u and XD^c/XD. Although such operators for XD^c/XD are fine, the operators for H_d/H_u must be suppressed. Otherwise, we will not have electroweak symmetry breaking due to large soft masses for H_d and H_u. For simplicity, we shall assume that these operators are suppressed due to the localizations of the particles in the extra space dimensions in type IIA/B string constructions, or the suppressed couplings with messengers. To be concrete, in the type IIA orientifold on $\mathbf{T^6}/(\mathbf{Z_2} \times \mathbf{Z_2})$ with intersecting D6-branes, all the particles except vector multiplets arise from the intersections of the D6-branes. The Yukawa couplings in the intersecting D-brane worlds arise from open string world-sheet instantons that connect three D-brane intersections [40–42]. For a given triplet of intersections, the minimal world-sheet action, which contributes to the Yukawa coupling, is weighted by a factor $e^{-A_{abc}}$, where A_{abc} is the world-sheet area of the triangles bounded by the branes a, b, and c [40–42]. Similar results hold for the four D-brane intersections (four-point interactions) [43] as well as the E_2 instanton effects [44,45]. Therefore, such operators for H_d/H_u and XD^c/XD can be suppressed easily by adjusting the world-sheet areas due to the exponential suppressions. On the other hand, even if we do not consider the explanations from the type IIA/B string constructions, the fine-tuning measures for these coupling hierarchies are about 25–70 since the soft masses for Φ and $T_{+/-}$ can be about 3–5 TeV. Such fine-tuning measures are similar to the following SUSY electroweak fine-tuning.

Let us consider two cases for M_*: (i) We choose $M_* = M_{\text{Pl}}$, $M_I = 10^8$ GeV, $\text{Tr}q^X = 2$, and $g_X = 10^{-3}/a$ with a a real number. So we get $D = 10^{22}/a^2$ GeV^2 and $F_S = 5.5a \times 10^{21}$ GeV^2. For $a = 2^{-1/2}$, we have $D/F_S = 5.1$; i.e., the Dirac gluino mass and the soft scalar masses of $T_{+/-}$ and Φ are about 5.1 times larger than the gravity mediation via F_S. This case may be realized in type IIA/B compactifications with the D-branes wrapping the large cycles but not in the heterotic string compactifications since g_X is small. (ii) We choose $M_I = 1.25 \times 10^5$ GeV, $\text{Tr}q^X = 2$, and $g_X = 0.8$, which may be realized in heterotic string as well. So we have $D = 2.44 \times 10^{10}$ GeV^2 and $F_S = 5.5 \times 10^{21}$ GeV^2. Thus, we need the effective messenger scale M_* around 10^6 GeV to realize the relatively heavy masses for Dirac gluino and scalar components of $T_{+/-}$ and Φ. In our model, the vectorlike particles like XD^c/XD can be messengers.

IV. PHENOMENOLOGY STUDY

First, with two-loop renormalization group equations (RGEs) for gauge couplings and two-loop RGEs for Yukawa couplings [46,47], we present gauge coupling unification in Fig. 1 for $M_V = M_D = 5$ TeV, and the GUT scale is around 10^{17} GeV. To avoid the Landau pole problem for gauge couplings, we need $M_V \geq 3$ TeV and $M_D \geq 3$ TeV. Thus, the contribution to Higgs boson mass from $\lambda H_u T_- H_u$ will be suppressed. In our model, we can have $m_{T_+} \gg M_V$, and then there exists a nondecoupling effect as in the Dirac next-to-MSSM (NMSSM) [48,49]. The Higgs boson mass is increased by

$$\Delta m_h^2 = \lambda_{\text{eff}}^2 \sin^4\beta v^2, \tag{8}$$

where $\tan\beta \equiv \langle H_u \rangle / \langle H_d \rangle$ and

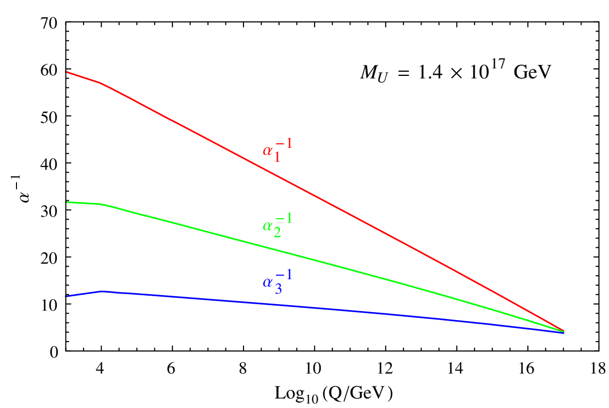

FIG. 1 (color online). Gauge coupling unification for $M_V = M_D = 5$ TeV.

TABLE III. The input parameter ranges or values used in our scans. All the mass parameters are given in appropriate power of TeV. Here, M_i are gaugino masses, μ is the bilinear Higgs mass in the superpotential and B_μ is the corresponding soft mass. We consider the universal scalar mass for the left- and right-handed squarks (sleptons) $\tilde{Q} \in \{\tilde{q}, \tilde{d}, \tilde{u}\}$ ($\tilde{L} \in \{\tilde{l}, \tilde{e}\}$) and the degenerated first and second generations. We choose $M_3 = 0.6$ TeV and the vanishing trilinear soft terms for three generations.

$\tan\beta$	$\lambda_{\rm eff}$	μ	B_μ	M_1	M_2	M_D	$m_{\tilde{Q},1\&2}$	$m_{\tilde{Q},3}$	$m_{\tilde{L},1\&2}$	$m_{\tilde{L},3}$	m_Φ
[5, 60]	[0.1, 0.7]	[0.3, 1]	[10^{-3}, 1]	[0.01, 0.1]	[0.5, 1]	[3, 5]	[0.8, 0.9]	[0.4, 0.7]	[0.1, 0.5]	[0.07, 0.16]	[$\sqrt{3}$, $\sqrt{5}$]

$$\lambda_{\rm eff}^2 \equiv \lambda^2 (m_{T_+}^2/(M_V^2 + m_{T_+}^2)). \tag{9}$$

Unlike the Dirac NMSSM, this contribution does not vanish at large $\tan\beta$ limit, which is properly accommodated with some interesting low-energy constraints such as the following Δa_μ.

V. NUMERICAL RESULTS

For the numerical studies, we are going to study the effective theory at the SUSY scale after integrating out the vectorlike particles. We implement our model in the Mathematica package SARAH [50–56]. SARAH is used in a second step to generate the various relevant outputs necessary for our analysis: we use the Fortran modules for SPHENO [57,58] to calculate the mass spectra and precision observables, and the model files for CALCHEP [59] which are used together with MICROMEGAs [60,61] to calculate the dark matter relic density and direct detection rates.

We consider all the current experimental constraints from the LEP, LHC, and B physics experiments, etc. The Higgs mass range is taken from 123 to 127 GeV. Also, the SM prediction for the anomalous magnetic moment of the muon [62,63] has a discrepancy with the experimental results [64,65] as follows:

$$\Delta a_\mu \equiv a_\mu(\text{exp}) - a_\mu(\text{SM}) = (28.6 \pm 8.0) \times 10^{-10}. \tag{10}$$

In our scans, the input parameter ranges or values are given in Table III. In Fig. 2, we present the Higgs mass versus $\lambda_{\rm eff}$ to show the large impact of the nondecoupling effect, in addition to the other constraints, especially the allowed range of Δa_μ. We see that for moderate values of $\lambda_{\rm eff}$ around 0.2–0.3, the Higgs mass falls into the desirable range, unlike the Dirac NMSSM. Another interesting property is that the electroweak symmetry breaking can be realized even in the range of small μ, which alleviates the following fine-tuning problem.

VI. FINE-TUNING

Because we discuss the simple low-energy phenomenology here, we consider the low-energy fine-tuning measure defined in Ref. [66] as follows:

$$\Delta_{\rm EW} = \frac{2}{M_Z^2} \max(C_{H_d}, C_{H_u}, C_\mu, C_{B_\mu}, C_{\delta m_{H_u}^2}), \tag{11}$$

where

$$C_{H_d} = \left|\frac{m_{H_d^2}}{\tan^2\beta - 1}\right|, \qquad C_{H_u} = \left|\frac{m_{H_u^2}\tan^2\beta}{\tan^2\beta - 1}\right|, \tag{12}$$

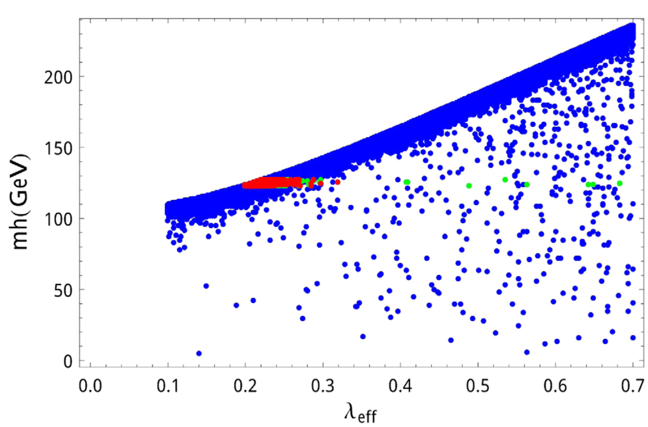

FIG. 2 (color online). The Higgs mass versus $\lambda_{\rm eff}$. The blue points provide the spectra without tachyons. In addition to satisfying the Higgs mass requirement, the green and red points have Δa_μ within 3σ and 1σ ranges, respectively.

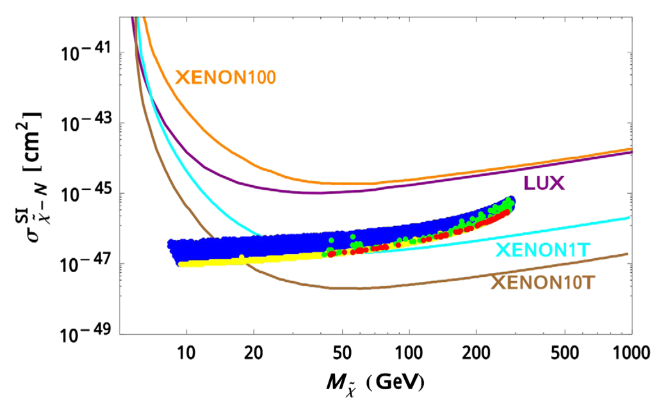

FIG. 3 (color online). The spin-independent LSP neutralino-nucleon cross section versus the LSP mass. The blue points have the particle spectra without tachyons. The yellow points satisfy the Higgs mass requirement and have Δa_μ within 3σ range. The green points have the correct relic density given in Eq. (15). And the red points satisfy all the current constraints.

TABLE IV. The particle spectrum (in GeV) for a benchmark point with pseudo-Dirac gluino masses 2927 and 3470 GeV for $\tan\beta = 29$, $M_1 = 0.21$ TeV, $\mu = 0.5$ TeV, $B_\mu = 0.02$ TeV^2, $M_2 = 0.5$ TeV, $M_3 = 0.6$ TeV, $M_D = 3$ TeV, $\lambda_{\rm eff} = 0.22$, $m_\Phi = 1.92$ TeV, $m_{\tilde{Q},1\&2} = 0.6$ TeV, $m_{\tilde{L},1\&2} = 0.8$ TeV $m_{\tilde{L},3} = 0.26$ TeV, $m_{\tilde{Q},3} = 0.55$ TeV. In this benchmark point, we have $\Delta_{\rm EW} = 60.4$, $\Omega_{\tilde{\chi}_1^0} h^2 = 0.1187$, $\Delta a_\mu = 9.96 \times 10^{-10}$, and the spin independent cross section $\sigma^{\rm SI}_{\tilde{\chi}-N} = 2.85 \times 10^{-46}$ cm^2.

$\tilde{\chi}_i^0$	$\tilde{\chi}_i^\pm$	$\tilde{\nu}_e, \tilde{\nu}_\tau$	$\tilde{e}_R, \tilde{e}_L$	$\tilde{\tau}_i$	$\tilde{u}_R, \tilde{u}_L$	$\tilde{t}_i$	$\tilde{d}_R, \tilde{d}_L$	$\tilde{b}_i$	h^0	$H^{0,\pm}/A^0$
(204,446,502,561)	(446,561)	(800,257)	(802,805)	(211,309)	(956,958)	(920,927)	(957,962)	(897,938)	124.8	$\simeq$705

$$C_\mu = |\mu^2|, \qquad C_{B_\mu} = |B_\mu|, \tag{13}$$

$$C_{\delta m^2_{H_u}} = \frac{(\lambda M_V)^2}{16\pi^2} \log\left(\frac{M_V^2 + m^2_{T_+}}{M_V^2}\right). \tag{14}$$

Compared to Ref. [66] we have additional $C_{\delta m^2_{H_u}}$ from the triplet threshold corrections to $m_{H_u^2}$. We find that the entire fine-tuning measure is given by C_μ while the other terms $C_{H_{d,u}}$, C_{B_μ} and $C_{\delta m^2_{H_u}}$ are negligible. In particular, the fine-tuning measure can be as low as 50 for the viable parameter space, even if the threshold effects at large M_V and M_D are considered. Since our MSSM sparticles except the gluino can be within about 1 TeV while gluino is Dirac, it seems that the fine-tuning measure from high-energy definition [67,68] will be small as well, which will be studied elsewhere. However, it might happen that an additional source of fine-tuning is needed to generate the correct hierarchy in the soft SUSY breaking terms to obtain sfermion masses like the ones of Table III together with a sizable $\lambda_{\rm eff}$ via RGE running when considering a constrained model.

VII. DARK MATTER

For simplicity, we concentrate on the LSP neutralino-stau coannihilation scenario here. To achieve this goal, we choose the following input parameter values or ranges: $\mu = 0.5$ TeV, $B_\mu = 0.15$ TeV^2, $M_2 = 0.5$ TeV, $M_3 = 0.6$ TeV, $M_D = 3$ TeV, $\lambda_{\rm eff} = 0.22$, $m_\Phi = 2$ TeV, $m_{\tilde{Q},1\&2} = m_{\tilde{L},1\&2} = 1$ TeV, $m_{\tilde{Q},3} = 0.404$ TeV, $5 < \tan\beta < 30$, 10 GeV $< M_1 <$ 300 GeV, 90 GeV $< m_{\tilde{L},3} <$ 300 GeV. All the other parameters are taken as in Table III. We use the relatively large values for μ and M_2 of 500 GeV to suppress the Higgsino and wino components of the LSP neutralino. This reduces the direct detection rates since the coupling to the Z boson is highly reduced. Moreover, we need a small mass splitting between the light stau and LSP neutralino to get an efficient coannihilation and to soften the LEP bounds on SUSY searches: for $m_{\tilde{\tau}_1} - m_{\tilde{\chi}_1^0} >$ 7 GeV a limit of $m_{\tilde{\tau}_1} > 87$ GeV is present while for nearly degenerated staus and neutralinos this limit becomes much weaker [69,70]. Finally, the fine-tuning measure is still small: $\Delta_{\rm EW} \simeq 60$.

As the preferred range for the LSP neutralino relic density, we consider the 2σ interval combined range from Planck + WP + highL + BAO [71]

$$0.1153 < \Omega_{\rm CDM} h^2 < 0.1221. \tag{15}$$

In Fig. 3, we show the results for spin-independent LSP neutralino-nucleon cross section. Because the masses of the first two generations of squarks have been fixed at 1 TeV and the Higgsino component in the LSP is heavily suppressed, the constraints from direct detection searches are easily evaded for all the considered points. The spin-independent cross sections are about one or 2 orders of magnitude below the current best limit provided by the LUX experiment [72]. Especially, the points with the LSP masses above 20 (15) GeV are within the reach of the projected XENON1T (XENON10T) sensitivity [73]. Also, we find that the current constraints on spin-dependent cross sections are much weaker at the moment. To be concrete, in Table IV, we present a viable benchmark point whose MSSM particles except gluino are within 1 TeV.

VIII. CONCLUSION

We have proposed the SSMs with a pseudo-Dirac gluino from hybrid F- and D-term SUSY breaking, which can be achieved via an anomalous $U(1)_X$ gauge symmetry inspired from string models. All the MSSM particles obtain the SUSY breaking soft terms from the traditional gravity mediation and can have masses within about 1 TeV except gluino. To escape the LHC SUSY search constraints and avoid the electroweak fine-tuning problem, the gluino also has a heavy Dirac mass above 3 TeV from D-term SUSY breaking. To realize the gauge coupling unification and lift the Higgs boson mass, we introduced extra vectorlike particles. We have studied the viable parameter space which satisfies all the current experimental constraints and given a concrete benchmark point. This kind of model keeps the merits of pre-LHC SSMs and solves the possible problems in the super-soft SUSY.

ACKNOWLEDGMENTS

This research was supported in part by the Natural Science Foundation of China under Grants No. 10821504, No. 11075194, No. 11135003, No. 11275246, and No. 11475238 and by the National Basic Research Program of China (973 Program) under Grant No. 2010CB833000 (T. L.).

[1] S. P. Martin, arXiv:hep-ph/9709356.
[2] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[3] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[4] M. Carena, S. Gori, N. R. Shah, and C. E. Wagner, J. High Energy Phys. 03 (2012) 014.
[5] G. Aad *et al.* (ATLAS Collaboration), J. High Energy Phys. 09 (2014) 176.
[6] J. Erler, P. Langacker, and T. Li, Phys. Rev. D **66**, 015002 (2002).
[7] U. Ellwanger, C. Hugonie, and A. M. Teixeira, Phys. Rep. **496**, 1 (2010).
[8] U. Ellwanger and C. Hugonie, Mod. Phys. Lett. A **22**, 1581 (2007).
[9] E. Ma, Phys. Lett. B **705**, 320 (2011).
[10] Y. Zhang, H. An, X.-d. Ji, and R. N. Mohapatra, Phys. Rev. D **78**, 011302 (2008).
[11] M. Hirsch, M. Malinsky, W. Porod, L. Reichert, and F. Staub, J. High Energy Phys. 02 (2012) 084.
[12] G. G. Ross and K. Schmidt-Hoberg, Nucl. Phys. **B862**, 710 (2012).
[13] S. Dimopoulos and G. Giudice, Phys. Lett. B **357**, 573 (1995).
[14] A. G. Cohen, D. Kaplan, and A. Nelson, Phys. Lett. B **388**, 588 (1996).
[15] T. J. LeCompte and S. P. Martin, Phys. Rev. D **84**, 015004 (2011).
[16] T. J. LeCompte and S. P. Martin, Phys. Rev. D **85**, 035023 (2012).
[17] H. K. Dreiner, M. Kramer, and J. Tattersall, Europhys. Lett. **99**, 61001 (2012).
[18] J. Fan, M. Reece, and J. T. Ruderman, J. High Energy Phys. 11 (2011) 012.
[19] T. Cheng, J. Li, and T. Li, J. Phys. G **42**, 065004 (2015).
[20] C. Csaki, Y. Grossman, and B. Heidenreich, Phys. Rev. D **85**, 095009 (2012).
[21] C. Csaki, E. Kuflik, and T. Volansky, Phys. Rev. Lett. **112**, 131801 (2014).
[22] P. J. Fox, A. E. Nelson, and N. Weiner, J. High Energy Phys. 08 (2002) 035.
[23] K. Benakli and M. Goodsell, Nucl. Phys. **B816**, 185 (2009).
[24] K. Benakli and M. Goodsell, Nucl. Phys. **B840**, 1 (2010).
[25] G. D. Kribs and A. Martin, Phys. Rev. D **85**, 115014 (2012).
[26] K. Benakli, M. D. Goodsell, and F. Staub, J. High Energy Phys. 06 (2013) 073.
[27] G. D. Kribs and A. Martin, arXiv:1308.3468.
[28] E. Bertuzzo, C. Frugiuele, T. Gregoire, and E. Ponton, J. High Energy Phys. 04 (2015) 089.
[29] K. Benakli, M. Goodsell, F. Staub, and W. Porod, Phys. Rev. D **90**, 045017 (2014).
[30] P. Diener, J. Kalinowski, W. Kotlarski, and D. Stckinger, J. High Energy Phys. 14 (2014) 124.
[31] A. E. Nelson and T. S. Roy, Phys. Rev. Lett. **114**, 201802 (2015).
[32] G. Giudice and A. Masiero, Phys. Lett. B **206**, 480 (1988).
[33] R. Ding, T. Li, F. Staub, C. Tian, and B. Zhu (to be published).
[34] W. Konetschny and W. Kummer, Phys. Lett. **70B**, 433 (1977).
[35] M. Cvetic, I. Papadimitriou, and G. Shiu, Nucl. Phys. **B659**, 193 (2003); **696**, 298 (2004).
[36] C. M. Chen, T. Li, and D. V. Nanopoulos, Nucl. Phys. **B751**, 260 (2006).
[37] P. Langacker and B. D. Nelson, Phys. Rev. D **72**, 053013 (2005).
[38] G. Dvali and A. Pomarol, Phys. Rev. Lett. **77**, 3728 (1996).
[39] M. Cvetic, L. L. Everett, and J. Wang, Phys. Rev. D **59**, 107901 (1999).
[40] G. Aldazabal, S. Franco, L. E. Ibanez, R. Rabadan, and A. M. Uranga, J. High Energy Phys. 02 (2001) 047.
[41] C. M. Chen, T. Li, V. E. Mayes, and D. V. Nanopoulos, Phys. Lett. B **665**, 267 (2008).
[42] C. M. Chen, T. Li, V. E. Mayes, and D. V. Nanopoulos, Phys. Rev. D **77**, 125023 (2008).
[43] C. M. Chen, T. Li, V. E. Mayes, and D. V. Nanopoulos, Phys. Rev. D **78**, 105015 (2008).
[44] R. Blumenhagen, M. Cvetic, and T. Weigand, Nucl. Phys. **B771**, 113 (2007).
[45] L. E. Ibanez and A. M. Uranga, J. High Energy Phys. 03 (2007) 052.
[46] S. P. Martin and M. T. Vaughn, Phys. Rev. D **50**, 2282 (1994).
[47] M. D. Goodsell, J. High Energy Phys. 01 (2013) 066.
[48] X. Lu, H. Murayama, J. T. Ruderman, and K. Tobioka, Phys. Rev. Lett. **112**, 191803 (2014).
[49] A. Kaminska, G. G. Ross, K. Schmidt-Hoberg, and F. Staub, J. High Energy Phys. 06 (2014) 153.
[50] F. Staub, arXiv:0806.0538.
[51] F. Staub, Comput. Phys. Commun. **182**, 808 (2011).
[52] F. Staub, T. Ohl, W. Porod, and C. Speckner, Comput. Phys. Commun. **183**, 2165 (2012).
[53] F. Staub, Comput. Phys. Commun. **184**, 1792 (2013).
[54] F. Staub, Comput. Phys. Commun. **185**, 1773 (2014).
[55] W. Porod, F. Staub, and A. Vicente, Eur. Phys. J. C **74**, 2992 (2014).
[56] M. D. Goodsell, K. Nickel, and F. Staub, Eur. Phys. J. C **75**, 32 (2015).
[57] W. Porod, Comput. Phys. Commun. **153**, 275 (2003).
[58] W. Porod and F. Staub, Comput. Phys. Commun. **183**, 2458 (2012).
[59] A. Belyaev, N. D. Christensen, and A. Pukhov, Comput. Phys. Commun. **184**, 1729 (2013).
[60] G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, Comput. Phys. Commun. **185**, 960 (2014).
[61] G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, Comput. Phys. Commun. **192**, 322 (2015).
[62] M. Davier, A. Hoecker, B. Malaescu, and Z. Zhang, Eur. Phys. J. C **71**, 1515 (2011).
[63] K. Hagiwara, R. Liao, A. D. Martin, D. Nomura, and T. Teubner, J. Phys. G **38**, 085003 (2011).
[64] G. Bennett *et al.* (Muon G-2 Collaboration), Phys. Rev. D **73**, 072003 (2006).
[65] G. Bennett *et al.* (Muon g-2 Collaboration), Phys. Rev. D **80**, 052008 (2009).
[66] H. Baer, V. Barger, P. Huang, D. Mickelson, A. Mustafayev, and X. Tata, Phys. Rev. D **87**, 035017 (2013).
[67] J. R. Ellis, K. Enqvist, D. V. Nanopoulos, and F. Zwirner, Mod. Phys. Lett. A **01**, 57 (1986).
[68] R. Barbieri and G. Giudice, Nucl. Phys. **B306**, 63 (1988).

[69] H. K. Dreiner, S. Heinemeyer, O. Kittel, U. Langenfeld, A. M. Weber, and G. Weiglein, Eur. Phys. J. C **62**, 547 (2009).
[70] J. Beringer *et al.* (Particle Data Group), Phys. Rev. D **86**, 010001 (2012).
[71] P. Ade *et al.* (Planck Collaboration), Astron. Astrophys. **571**, A16 (2014).
[72] D. Akerib *et al.* (LUX Collaboration), Phys. Rev. Lett. **112**, 091303 (2014).
[73] XENON Collaboration (unpublished).