PHYSICAL REVIEW D **86,** 115014 (2012)

Compact supersymmetry

Hitoshi Murayama,[1,2,3] Yasunori Nomura,[1,2] Satoshi Shirai,[1,2] and Kohsaku Tobioka[3,4]

[1]*Berkeley Center for Theoretical Physics, Department of Physics, University of California, Berkeley, California 94720, USA*
[2]*Theoretical Physics Group, Lawrence Berkeley National Laboratory, Berkeley, California 94720, USA*
[3]*Kavli Institute for the Physics and Mathematics of the Universe (WPI), Todai Institutes for Advanced Study, University of Tokyo, Kashiwa 277-8583, Japan*
[4]*Department of Physics, University of Tokyo, Hongo, Tokyo 113-0033, Japan*

(Received 30 June 2012; published 7 December 2012)

Supersymmetry broken geometrically in extra dimensions naturally leads to a nearly degenerate spectrum for superparticles, ameliorating the bounds from the current searches at the LHC. We present a minimal such model with a single extra dimension, and show that it leads to viable phenomenology despite the fact that it essentially has two fewer free parameters than the conventional constrained minimal supersymmetric standard model. The theory does not suffer from the supersymmetric flavor or *CP* problem because of the universality of geometric breaking, and it automatically yields near-maximal mixing in the scalar top sector with $|A_t| \approx 2m_{\tilde{t}}$ to boost the Higgs boson mass. Despite the rather constrained structure, the theory is less fine-tuned than many supersymmetric models.

DOI: 10.1103/PhysRevD.86.115014 PACS numbers: 12.60.Jv, 11.25.Wx, 14.80.Ly, 14.80.Rt

I. INTRODUCTION

Supersymmetry has widely been regarded as the prime candidate for physics beyond the standard model [1]. It can explain the dynamical origin of electroweak symmetry breaking through renormalization group effects and provide a natural candidate for the cosmological dark matter in its simplest incarnations. In particular, it stabilizes the large hierarchy between the electroweak scale $\approx$ TeV and the quantum gravity scale $\approx 10^{15}$ TeV against radiative corrections to the Higgs mass parameter. Barring a fine-tuning among parameters of the theory, this consideration strongly suggests the existence of superparticles below $\approx$ TeV. The mass spectrum of superparticles has been mostly discussed within the "minimal supergravity" or constrained minimal supersymmetric standard model (CMSSM) framework [2], which typically generates a widely spread spectrum leading to experimentally identifiable large visible and missing energies.

However, no experimental hints have been seen so far at the Large Hadron Collider (LHC), which has led to substantial anxiety in the community. Moreover, the suggested mass of 125 GeV for the Higgs boson by the LHC data [3] is not easily accommodated in the minimal supersymmetric standard model (MSSM), where one has to rely on radiative corrections to push the Higgs boson mass beyond the tree-level upper bound of $m_Z \simeq 91$ GeV. These requirements push the scalar quark masses well beyond the TeV within the CMSSM.

There are three main suggestions to allow for supersymmetry without the signal so far, within the context of *R*-parity conserving supersymmetry. One is to simply accept a fine-tuning to maintain the hierarchy against radiative corrections, at a level significantly worse than a percent. Quite often, the anthropic principle is brought in to justify this level of fine-tuning [4]. The second is to keep superparticles relevant to the Higgs mass parameter below TeV while to assume all other superparticles well beyond TeV [5]. The third is to assume that all superparticles are nearly degenerate, making them somewhat hidden from experimental searches because of low Q values in visible and missing energies. The last option, however, has been discussed only phenomenologically [6], lacking theoretical justifications based on simple and explicit models of supersymmetry breaking.

In this Letter, we point out that the third possibility of a nearly degenerate superparticle spectrum is quite automatic when supersymmetry is broken by boundary conditions in compact extra dimensions, the so-called Scherk-Schwarz mechanism [7,8]. With the simplest extra dimension—the $S^1/\mathbb{Z}_2$ orbifold—the mechanism has a rather simple structure [9]. In particular, locating matter and Higgs fields in the bulk and on a brane, respectively, and forbidding local-parity violating bulk mass parameters for the matter fields, the theory has only four parameters relevant for the spectrum of superparticles: the compactification scale $1/R$, the 5D cutoff scale Λ ($> 1/R$), the supersymmetry-breaking twist parameter α ($\in [0, \frac{1}{2}]$), and the supersymmetric Higgs mass μ.

Using the common notation in the MSSM, the spectrum of superparticles is given at the compactification scale $\approx 1/R$ as

$$M_{1/2} = \frac{\alpha}{R}, \qquad m^2_{\tilde{Q},\tilde{U},\tilde{D},\tilde{L},\tilde{E}} = \left(\frac{\alpha}{R}\right)^2, \qquad m^2_{H_u,H_d} = 0,$$
$$A_0 = -\frac{2\alpha}{R}, \qquad \mu \neq 0, \qquad B = 0, \tag{1}$$

at tree level. While these masses receive radiative corrections from physics at and above $1/R$, they are under control because of the symmetries in higher-dimensional spacetime, and thus can naturally be small. Therefore, in this limit, the theory essentially has only three free parameters:

$$\frac{1}{R}, \qquad \frac{\alpha}{R}, \qquad \mu. \tag{2}$$

This rather compact set of parameters gives all the superparticle as well as the Higgs boson masses.

Even though Eq. (2) gives two less parameters than in the traditional CMSSM framework, we show that it still leads to viable phenomenology. In addition, it solves the flavor problem that often plagues models of supersymmetry breaking, because the geometry is universal to all scalar particles and hence respects a large flavor symmetry. The problem of accommodating a large enough Higgs boson mass is ameliorated by the near degeneracy between $\tilde{t}_L$ and $\tilde{t}_R$, and $|A_t| \approx 2m_{\tilde{t}}$. And the degenerate spectrum at tree level automatically achieves a compact spectrum that allows superparticles to be hidden from the current searches even when they are below TeV.

This article is organized as follows. We first review the basics of supersymmetry breaking by boundary conditions in the $S^1/\mathbb{Z}_2$ orbifold and present the simplest model we study. We then present the low-energy spectrum of superparticles and discuss its phenomenology. We also provide benchmark points useful for further phenomenological studies of the model.

II. SUPERSYMMETRY BREAKING BY BOUNDARY CONDITIONS

We consider a single compact extra dimension with the coordinate y identified under T: $y \to y + 2\pi R$ and P: $y \to -y$. These two operations satisfy the algebra $PTP = T^{-1}$ and $P^2 = 1$, and the resulting extra dimension is an interval $y \in [0, \pi R]$: the $S^1/\mathbb{Z}_2$ orbifold.

We consider a supersymmetric $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge theory in this spacetime, with the gauge and three generations of matter supermultiplets propagating in the bulk. The boundary conditions for these fields are given such that the $SU(2)_R$ doublets in these multiplets transform as

$$P = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \qquad T = \begin{pmatrix} \cos(2\pi\alpha) & -\sin(2\pi\alpha) \\ \sin(2\pi\alpha) & \cos(2\pi\alpha) \end{pmatrix}, \tag{3}$$

under P and T (see Ref. [9] for details). In this paper we consider the case $\alpha \ll 1$.[1]

The boundary conditions of Eq. (3) leave only the MSSM gauge and matter fields below the compactification scale $1/R$. Specifically, the matter supermultiplets yield three generations of quarks and leptons as the zero modes, while their superpartners obtain the common soft mass of α/R. (Here, we have assumed that there are no 5D bulk mass terms for the matter multiplets.[2]) The gauge supermultiplets give massless standard model gauge fields and gauginos of mass α/R. We therefore obtain the first two expressions in Eq. (1). [The Kaluza-Klein excitations form $N = 2$ supermultiplets and have masses $\approx n/R$ $(n = 1, 2, \ldots)$, with supersymmetry-breaking mass splitting of order α/R.]

The Higgs chiral superfields H_u and H_d are located on one of the branes at $y = 0$. The Yukawa couplings and μ term can then be written on that brane:

$$\mathcal{L}_{\rm brane} = \delta(y) \int d^2\theta (y_U^{ij} Q_i U_j H_u + y_D^{ij} Q_i D_j H_d + y_E^{ij} L_i E_j H_d + \mu H_u H_d). \tag{4}$$

This leads to the other four expressions in Eq. (1). (Here, we have simply assumed the existence of the μ term on the brane. We leave discussions of its origin to future work.)

Note that the degeneracy among the three generations of scalars is automatic because of the geometric nature of the supersymmetry breaking. In addition, since α and μ can always be taken real by phase redefinitions of fields associated with R and Peccei-Quinn rotations, there is no physical phase in $M_{1/2}$, A_0, μ, or B. Therefore, the flavor problem as well as the CP problem are automatically solved in this model.

The expressions in Eq. (1) receive corrections from physics above and at $1/R$. In the 5D picture, corrections above $1/R$ come from brane-localized kinetic terms for the gauge and matter supermultiplets, and affect $M_{1/2}$, $m_{\tilde{f}}^2 \equiv m^2_{\tilde{Q},\tilde{U},\tilde{D},\tilde{L},\tilde{E}}$, and A_0. These terms have tree-level contributions at Λ and radiative ones between $1/R$ and Λ. From dimensional analysis, the size of the radiative contributions is

$$\frac{\delta M_{1/2}}{M_{1/2}}, \qquad \frac{\delta m_{\tilde{f}}^2}{m_{\tilde{f}}^2}, \qquad \frac{\delta A_0}{A_0} \approx O\left(\frac{g^2, y^2}{16\pi^2} \ln(\Lambda R)\right). \tag{5}$$

Moreover, it is (technically) natural to assume that the tree-level contributions do not exceed the radiative ones with $\ln(\Lambda R) \to O(1)$. Therefore, with this assumption, the corrections to Eq. (1) from physics above $1/R$ are always negligible for $\Lambda R \lesssim 16\pi^2$, i.e., when our effective higher dimensional field theory is valid. (The same can also be seen in the 4D picture. In this picture, $N = 2$ supersymmetry existing for the $n > 0$ modes leads to nontrivial cancellations of the corrections to $M_{1/2}$, $m_{\tilde{f}}^2$, and A_0 from these modes. To see the cancellations for the gauge multiplets, the effect of anomalies must be taken into account correctly. The explicit demonstration of these nontrivial cancellations will be given elsewhere.)

[1]The twist parameter α in the boundary conditions is equivalent to an F-term vacuum expectation value of the radion superfield [10], which can be generated dynamically through a radion stabilization mechanics and hence can be naturally small.

[2]This assumption can be justified by a local parity in the bulk; see, e.g., Ref. [11].

The corrections from physics at $1/R$ arise from nonlocal operators in 5D. They affect all the supersymmetry-breaking masses and are of order $1/16\pi^2$. Here we calculate only the contributions to the Higgs mass parameters, which could potentially affect the analysis of electroweak symmetry breaking. By choosing the renormalization scale to be $1/(2\pi R)$, we find these corrections are

$$\delta m_{H_u}^2 = \left(-\frac{33y_t^2}{8\pi^2} + \frac{9(g_2^2+g_1^2/5)}{16\pi^2}\right)\left(\frac{\alpha}{R}\right)^2,$$
$$\delta m_{H_d}^2 = \frac{9(g_2^2+g_1^2/5)}{16\pi^2}\left(\frac{\alpha}{R}\right)^2, \qquad (6)$$
$$\delta B = \left(\frac{9y_t^2}{8\pi^2} - \frac{3(g_2^2+g_1^2/5)}{8\pi^2}\right)\frac{\alpha}{R},$$

where we have included only the contributions from the top-Yukawa coupling, y_t, and $SU(2)_L$ and $U(1)_Y$ gauge couplings, g_2 and g_1 [in the $SU(5)$ normalization].

In summary, the low-energy superparticle masses are obtained by evolving down Eqs. (1) and (6) defined at the renormalization scale

$$\mu_{\rm RG} = \frac{1}{2\pi R}, \qquad (7)$$

using the MSSM renormalization group equations. Incidentally, the gravitino mass is $m_{3/2} = \alpha/R$, generated by the supersymmetry-breaking twist in the fifth dimension.

III. SUPERPARTICLE SPECTRUM

Following the procedure described above, we calculate the MSSM mass spectrum using SOFTSUSY 3.3.1 [12] and the lightest Higgs boson mass using FEYNHIGGS 2.8.6 [13]. In Fig. 1, we plot the contours of the mass of the lightest Higgs boson, M_H, and the fine-tuned parameter, defined by $\Delta^{-1} = \min_x |\partial \ln m_Z^2/\partial \ln x|^{-1}$ with $x = \alpha, \mu, 1/R, y_t, g_3, \ldots,$ in the $1/R$-α/R plane. (The fine-tuning parameter is determined mostly by $x = \mu$.) In the calculation, we have used the top-quark mass of $m_t = 173.2$ GeV [14]. Varying it by 1σ, $\Delta m_t = \pm 0.9$ GeV, affects the Higgs boson mass by $\Delta M_H \approx \pm 1$ GeV. Also, theoretical errors in M_H are expected to be about $|\Delta M_H| \approx 2$–3 GeV [15], so that the regions with $M_H \gtrsim 121$–123 GeV in the plot are not necessarily incompatible with the 125 GeV Higgs boson hinted at the LHC [3]. Indeed, using the recently released program H3m [16], which includes a partial three-loop effect, we find that the corrections to M_H from higher order effects are positive and of order a few GeV in most of the parameter region in the plot.

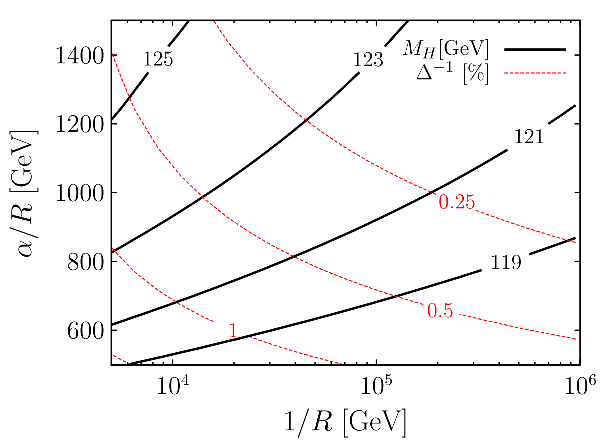

FIG. 1 (color online). The lightest Higgs boson mass M_H (in GeV) and the fine-tuned parameter Δ^{-1}. Note that there is an approximately 3 GeV systematic error in theoretical computation of M_H.

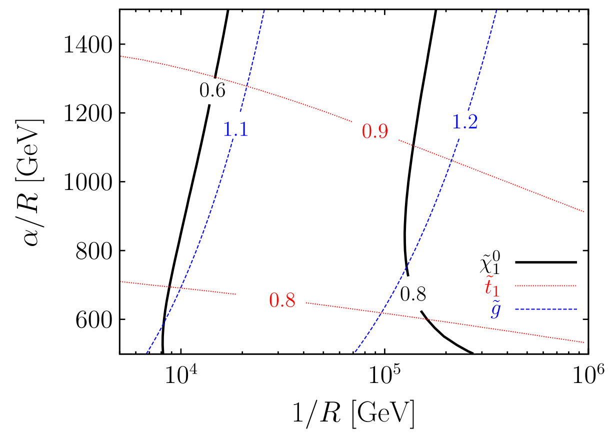

FIG. 2 (color online). Masses of the lightest neutralino $\tilde{\chi}_1^0$, the lighter top squark $\tilde{t}_1$, and the gluino $\tilde{g}$ normalized to α/R.

In Fig. 2, the masses of selected superparticles (the lightest neutralino $\tilde{\chi}_1^0$, the lighter top squark $\tilde{t}_1$, and the gluino $\tilde{g}$) are shown. The masses of the first and second generation squarks are almost the same as the gluino mass. The masses of the electroweak superparticles are close to α/R, except for the lightest two neutralinos $\tilde{\chi}_{1,2}^0$ and the lighter chargino $\tilde{\chi}_1^+$, which are Higgsino-like (and thus close in mass) in most of the parameter space. We find that the masses of the superparticles are degenerate at a 10% level, except possibly for the Higgsinos, which can be significantly lighter (up to a factor of ≈ 2).

IV. EXPERIMENTAL LIMITS

As we have seen, the model naturally predicts a degenerate mass spectrum for superparticles. This has strong implications on supersymmetry searches at the LHC. Because of the mass degeneracy, production of high p_T jets and large missing energy is suppressed. Therefore, typical searches, based on high p_T jets and large missing energy, are less effective for the present model.

To estimate the number of supersymmetric events, we have used ISAJET 7.72 [17] for the decay table of superparticles, HERWIG 6.520 [18] for the generation of supersymmetric events, ACERDET 1.0 [19] for the detector simulation, and NLL-fast [20] for the estimation of the production cross

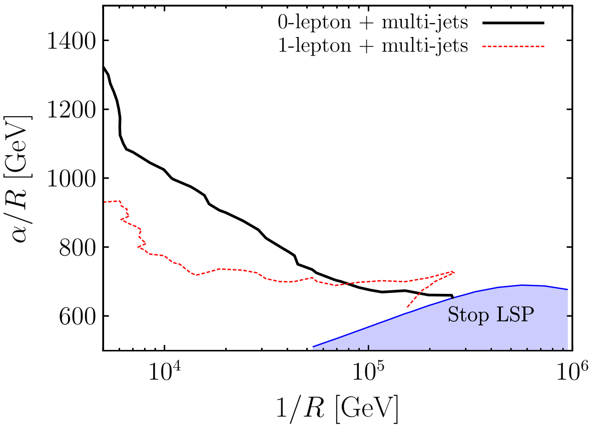

FIG. 3 (color online). The current LHC constraint on the model.

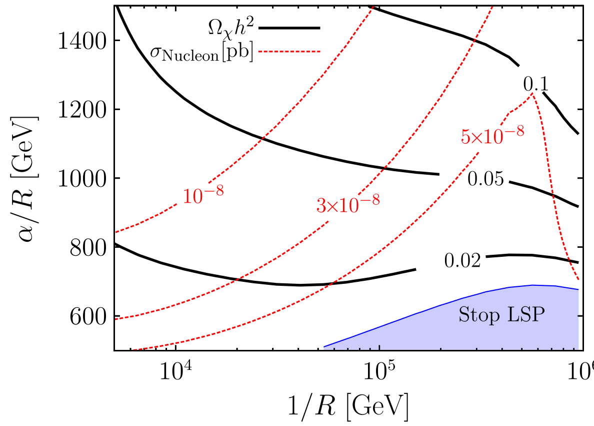

FIG. 4 (color online). The thermal relic abundance, $\Omega_\chi h^2$, and the spin-independent cross section with a nucleon, $\sigma_{\rm Nucleon}$, of $\tilde{\chi}_1^0$. The solid (black) lines are the contours of $\Omega_\chi h^2$, while the dotted (red) lines are those of $\sigma_{\rm Nucleon}$.

section including next-to-leading order QCD corrections and the resummation at next-to-leading-logarithmic accuracy. To constrain the parameter space, we compare the obtained event numbers with the results of ATLAS searches for multijets plus large missing energy with and without a lepton at $L = 4.7\ {\rm fb}^{-1}$ at $\sqrt{s} = 7$ TeV [21,22]. In Fig. 3, we show the resulting LHC constraint on the model. Other searches such as those for b-jets and/or multileptons are less effective. We find that for $1/R \gtrsim 10^5$ GeV, the case that $m_{\tilde{g}} \simeq m_{\tilde{q}} \lesssim 1$ TeV is still allowed. This constraint is significantly weaker than that on the CMSSM, which excludes $m_{\tilde{g}} \lesssim 1.4$ TeV for $m_{\tilde{g}} \simeq m_{\tilde{q}}$ [21]. (We have checked that our naive method of estimating the LHC constraints adopted here reproduces this bound for the CMSSM spectra.)

We note that since B is not a free parameter in the present model, $\tan\beta$ is determined by the electroweak symmetry breaking condition. We typically find $\tan\beta \sim 4$–10. This allows for the model to avoid the constraint from $b \to s\gamma$, despite the large A terms.

The contribution of the Kaluza-Klein states to the electroweak precision parameters bounds $1/R \gtrsim$ a few TeV [23]. Since we consider the region $1/R \gtrsim 10$ TeV in this paper, however, the model is not constrained by the electroweak precision data.

V. DARK MATTER

In the present model, the dark matter candidate is the lightest neutralino $\tilde{\chi}_1^0$, whose dominant component is the Higgsino. In Fig. 4, we show the thermal relic abundance, $\Omega_\chi h^2$, and the spin-independent cross section with a nucleon, $\sigma_{\rm Nucleon}$, of $\tilde{\chi}_1^0$, assuming R-parity conservation. To estimate these, we have used micrOMEGAs 2.4 [24]. For the strange quark form factor we have adopted $f_s = 0.02$, suggested by lattice calculations [25], instead of the default value of micrOMEGAs ($f_s = 0.26$). As seen in Fig. 4, the thermal relic abundance of $\tilde{\chi}_1^0$ is much smaller than the observed dark matter density $\Omega_{\rm DM} h^2 \simeq 0.1$, unless $\tilde{\chi}_1^0$ is rather heavy $\sim$ TeV (in the upper-right corner of the plot). Therefore, in most parameter regions, $\tilde{\chi}_1^0$ cannot be the dominant component of dark matter if only the thermal relic abundance is assumed. It must be produced nonthermally to saturate $\Omega_{\rm DM} h^2$, or some other particle(s), e.g., the axion/axino, must make up the rest.

It is natural, however, to expect that at least the thermal abundance of $\tilde{\chi}_1^0$ remains as a (sub)component of dark matter. In this case, direct and indirect signatures of the relic neutralino are expected. To discuss the direct-detection signal, it is useful to define

$$\sigma_{\rm Nucleon}^{\rm eff} \equiv \sigma_{\rm Nucleon} \frac{\min\{\Omega_\chi, \Omega_{\rm DM}\}}{\Omega_{\rm DM}}, \qquad (8)$$

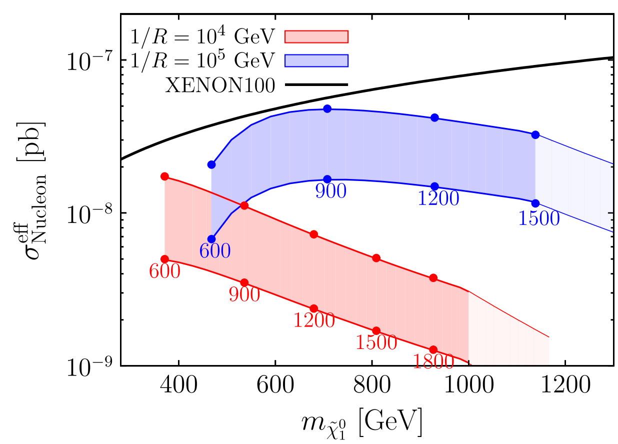

FIG. 5 (color online). Effective dark matter—nucleon cross section for $1/R = 10^4$ GeV (lower, red shaded) and 10^5 GeV (upper, blue shaded). In each region, the upper and lower borders correspond to $f_s = 0.26$ and 0.02, respectively, and the dots represent the corresponding values of α/R. (The very light shaded regions are those in which the thermal abundance exceeds $\Omega_{\rm DM}$.) The solid (black) line shows the current upper bound from XENON100.

TABLE I. Phenomenologically viable mass spectrum of the benchmark points (in GeV). Point 1: $1/R = 10^4$ GeV, $\alpha/R = 1400$ GeV, and Point 2: $1/R = 10^5$ GeV, $\alpha/R = 800$ GeV.

Particle	Point 1	Point 2	Particle	Point 1	Point 2
$\tilde{g}$	1494	949	...	...	...
$\tilde{u}_L$	1467	939	$\tilde{u}_R$	1459	925
$\tilde{d}_L$	1469	942	$\tilde{d}_R$	1458	924
$\tilde{b}_2$	1460	924	$\tilde{b}_1$	1430	875
$\tilde{t}_2$	1557	988	$\tilde{t}_1$	1267	681
$\tilde{\nu}$	1411	822	$\tilde{\nu}_\tau$	1410	822
$\tilde{e}_L$	1413	826	$\tilde{e}_R$	1406	812
$\tilde{\tau}_2$	1417	823	$\tilde{\tau}_1$	1402	809
$\tilde{\chi}_1^0$	767	630	$\tilde{\chi}_2^0$	777	671
$\tilde{\chi}_3^0$	1384	755	$\tilde{\chi}_4^0$	1410	821
$\tilde{\chi}_1^\pm$	771	642	$\tilde{\chi}_2^\pm$	1409	817
h^0	125	120	H^0	819	718
A^0	819	717	$H^\pm$	822	722

which is the quantity to be compared with the dark matter—nucleon cross section in the usual direct-detection exclusion plots (which assume $\Omega_\chi = \Omega_{\rm DM}$). In Fig. 5, we plot $\sigma^{\rm eff}_{\rm Nucleon}$ as a function of $m_{\tilde{\chi}_1^0}$ for $1/R = 10^4$ GeV and 10^5 GeV. To represent the uncertainty from the nucleon matrix element, we show both the $f_s = 0.02$ and 0.26 cases. We also present the current upper bound on $\sigma^{\rm eff}_{\rm Nucleon}$ from XENON100 [26]. We find that improving the bound by 1 or 2 orders of magnitude will cover a significant portion of the parameter space of the model.

VI. FINAL REMARKS

In this article, we pointed out that supersymmetry broken by boundary conditions in extra dimensions leads naturally to a nearly degenerate superparticle spectrum, ameliorating the limits from experimental searches. We presented the simplest such model in the $S^1/\mathbb{Z}_2$ orbifold, and showed that it leads to viable phenomenology despite the fact that it essentially has two fewer free parameters than the CMSSM: $1/R$, α/R, and μ. In Table I we give two representative points in the parameter space, which can serve benchmark points for further phenomenological studies.

The theory presented here can be extended in several different ways. An interesting one is to introduce a singlet field S together with superpotential interactions on the $y = 0$ brane: $\lambda S H_u H_d + f(S)$, where $f(S)$ is a polynomial of S with the simplest possibility being $f(S) = -\kappa S^3/3$. This allows for an extra contribution to the Higgs boson mass from λ, and it can make the lightest neutralino (which would now contain a singlino component as well) saturate the observed dark matter abundance without resorting to non-thermal production. Detailed studies of this possibility will be presented elsewhere.

ACKNOWLEDGMENTS

This work was supported in part by the DOE under Contract No. DE-AC02-05CH11231. The work of H. M. was also supported by the NSF under Grant No. PHY-1002399, the JSPS Grant (C) No. 23540289, the FIRST program "Subaru Measurements of Images and Redshifts (SuMIRe)," CSTP, and by WPI, MEXT, Japan. The work of Y. N. was supported in part by the NSF under Grant No. PHY-0855653. The work of K. T. was supported in part by the Grant-in-Aid for JSPS Fellows, and K. T. acknowledges the financial support of Friends of Todai, Inc., to conduct research in Berkeley.

[1] For a review, see, e.g., S. P. Martin, in *Perspectives on Supersymmetry II*, edited by G. L. Kane (World Scientific, Singapore, 2010), p. 1.
[2] A. H. Chamseddine, R. Arnowitt, and P. Nath, Phys. Rev. Lett. **49**, 970 (1982); R. Barbieri, S. Ferrara, and C. A. Savoy, Phys. Lett. **119B**, 343 (1982); L. Hall, J. Lykken, and S. Weinberg, Phys. Rev. D **27**, 2359 (1983).
[3] ATLAS Collaboration, Phys. Lett. B **710**, 49 (2012); S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **710**, 26 (2012).
[4] N. Arkani-Hamed and S. Dimopoulos, J. High Energy Phys. 06 (2005) 073; L. J. Hall and Y. Nomura, J. High Energy Phys. 01 (2012) 082.
[5] S. Dimopoulos and G. F. Giudice, Phys. Lett. B **357**, 573 (1995); A. Pomarol and D. Tommasini, Nucl. Phys. **B466**, 3 (1996); A. G. Cohen, D. B. Kaplan, and A. E. Nelson, Phys. Lett. B **388**, 588 (1996).
[6] See, e.g., T. J. LeCompte and S. P. Martin, Phys. Rev. D 85, 035023 (2012); Report No. ATLAS-CONF-2011-155, and references therein.
[7] J. Scherk and J. H. Schwarz, Phys. Lett. **82B**, 60 (1979); Nucl. Phys. **B153**, 61 (1979).
[8] A. Pomarol and M. Quirós, Phys. Lett. B **438**, 255 (1998); I. Antoniadis, S. Dimopoulos, A. Pomarol, and M. Quirós, Nucl. Phys. **B544**, 503 (1999); R. Barbieri, L. J. Hall, and Y. Nomura, Phys. Rev. D **63**, 105007 (2001).
[9] R. Barbieri, L. J. Hall, and Y. Nomura, Phys. Rev. D **66**, 045025 (2002); Nucl. Phys. **B624**, 63 (2002).
[10] D. Marti and A. Pomarol, Phys. Rev. D **64**, 105025 (2001); D. E. Kaplan and N. Weiner, arXiv:hep-ph/0108001.
[11] R. Barbieri, L. J. Hall, G. Marandella, Y. Nomura, T. Okui, S. J. Oliver, and M. Papucci, Nucl. Phys. **B663**, 141 (2003).
[12] B. C. Allanach, Comput. Phys. Commun. **143**, 305 (2002).

[13] S. Heinemeyer, W. Hollik, and G. Weiglein, Comput. Phys. Commun. **124**, 76 (2000).
[14] The Tevatron Electroweak Working Group, arXiv:1107.5255.
[15] B. C. Allanach, A. Djouadi, J. L. Kneur, W. Porod, and P. Slavich, J. High Energy Phys. 09 (2004) 044.
[16] P. Kant, R. V. Harlander, L. Mihaila, and M. Steinhauser, J. High Energy Phys. 08 (2010) 104.
[17] F. E. Paige, S. D. Protopopescu, H. Baer, and X. Tata, arXiv:hep-ph/0312045.
[18] G. Corcella, I. G. Knowles, G. Marchesini, S. Moretti, K. Odagiri, P. Richardson, M. H. Seymour, and B. R. Webber, J. High Energy Phys. 01 (2001) 010; arXiv:hep-ph/0210213.
[19] E. Richter-Wąs, arXiv:hep-ph/0207355.
[20] http://web.physik.rwth-aachen.de/service/wiki/bin/view/Kraemer/SquarksandGluinos.
[21] ATLAS Collaboration, Report No. ATLAS-CONF-2012-033.
[22] ATLAS Collaboration, Report No. ATLAS-CONF-2012-041.
[23] See, e.g., A. Delgado and M. Quirós, Nucl. Phys. **B607**, 99 (2001); A. Delgado, A. Pomarol, and M. Quiros, J. High Energy Phys. 01 (2000) 030.
[24] G. Bélanger, F. Boudjema, P. Brun, A. Pukhov, S. Rosier-Lees, P. Salati, and A. Semenov, Comput. Phys. Commun. **182**, 842 (2011); G. Bélanger, F. Boudjema, A. Pukhov, and A. Semenov, Comput. Phys. Commun. **149**, 103 (2002); Comput. Phys. Commun. **174**, 577 (2006).
[25] H. Ohki *et al.* (JLQCD Collaboration), Phys. Rev. D **78**, 054502 (2008); D. Toussaint and W. Freeman (MILC Collaboration), Phys. Rev. Lett. **103**, 122002 (2009).
[26] E. Aprile *et al.* (XENON100 Collaboration), Phys. Rev. Lett. **107**, 131302 (2011).