Implication of a Higgs boson at 125 GeV within the stochastic superspace framework

Manimala Chakraborti,[1,*] Utpal Chattopadhyay,[1,†] and Rohini M. Godbole[2,‡]

[1]*Department of Theoretical Physics, Indian Association for the Cultivation of Science,*
2A & B Raja S. C. Mullick Road, Jadavpur, Kolkata 700 032, India
[2]*Centre for High Energy Physics, Indian Institute of Science, Bangalore 560 012, India*

We revisit the issue of considering stochasticity of Grassmannian coordinates in $N = 1$ superspace, which was analyzed previously by Kobakhidze *et al.* In this stochastic supersymmetry (SUSY) framework, the soft SUSY breaking terms of the minimal supersymmetric Standard Model (MSSM) such as the bilinear Higgs mixing, trilinear coupling, as well as the gaugino mass parameters are all proportional to a single mass parameter ξ, a measure of supersymmetry breaking arising out of stochasticity. While a nonvanishing trilinear coupling at the high scale is a natural outcome of the framework, a favorable signature for obtaining the lighter Higgs boson mass m_h at 125 GeV, the model produces tachyonic sleptons or staus turning to be too light. The previous analyses took Λ, the scale at which input parameters are given, to be larger than the gauge coupling unification scale M_G in order to generate acceptable scalar masses radiatively at the electroweak scale. Still, this was inadequate for obtaining m_h at 125 GeV. We find that Higgs at 125 GeV is highly achievable, provided we are ready to accommodate a nonvanishing scalar mass soft SUSY breaking term similar to what is done in minimal anomaly mediated SUSY breaking (AMSB) in contrast to a pure AMSB setup. Thus, the model can easily accommodate Higgs data, LHC limits of squark masses, WMAP data for dark matter relic density, flavor physics constraints, and XENON100 data. In contrast to the previous analyses, we consider $\Lambda = M_G$, thus avoiding any ambiguities of a post-grand unified theory physics. The idea of stochastic superspace can easily be generalized to various scenarios beyond the MSSM.

I. INTRODUCTION

Low energy supersymmetry (SUSY) [1–4] has been one of the most promising candidates for a theory of fundamental particles and interactions going beyond the Standard Model (SM); the so-called BSM physics. The minimal extension of the SM including SUSY, namely, the minimal supersymmetric Standard Model (MSSM), extends the particle spectrum of the SM by one additional Higgs doublet and the supersymmetric partners of all the SM particles—the sparticles. The supersymmetric extension of the SM provides a particularly elegant solution to the problem of stabilizing the electroweak (EW) symmetry breaking scale against large radiative correction and keeps the Higgs ''naturally'' light. In fact, a very robust upper limit on the mass of the lightest Higgs boson is perhaps one of the important predictions of this theory. Further, this upper limit is linked in an essential way to the values of some of the SUSY breaking parameters in the theory. In addition, in *R*-parity conserving SUSY, the lightest supersymmetric particle (LSP) emerges as the natural candidate for the dark matter (DM), the existence of which has been proved beyond any doubt in astrophysical experiments. The search for evidence of the realization of this symmetry

in nature (in the context of high energy collider experiments, precision measurements at the high intensity *B* factories, and in the DM detection experiments) has therefore received enormous attention of particle physicists, perhaps only next to the Higgs boson. The recent observation of a boson with mass around ~125 GeV at the Large Hadron Collider (LHC) [5] and the rather strong lower limits on the masses of sparticles that possess strong interactions that the LHC searches have yielded [6] necessitates careful studies of the MSSM in the context of all the recent low energy data. In these studies, it is also very important to seek suitable guiding principles which could possibly reduce the associated large number of SUSY breaking parameters of MSSM. Thus, looking for modes of specific SUSY breaking mechanisms that involve only a few input quantities given at a relevant scale can be useful. Here, the soft SUSY breaking parameters at the electroweak scale are found via renormalization group (RG) analyses. Apart from the simplicity of having a few parameters as input, such schemes create challenging balancing acts. On the one hand, various soft breaking masses and couplings become correlated with one another in such schemes. On the other hand, overall one has to accommodate a large number of very stringent low energy constraints in a comprehensive model consisting of only a few parameters. A simple and well-motivated example of a SUSY model is the minimal supergravity (mSUGRA) [7]. Here, SUSY is broken spontaneously in a hidden sector and the breaking

*tpmc@iacs.res.in
†tpuc@iacs.res.in
‡rohini@cts.iisc.ernet.in

is communicated to the observable sector where MSSM resides via Planck mass suppressed supergravity interactions. The model involves soft SUSY breaking parameters like (i) universal gaugino mass parameter $m_{\frac{1}{2}}$, (ii) the universal scalar mass parameter m_0, (iii) the universal trilinear coupling A_0, and (iv) the universal bilinear coupling B_0, all given at the gauge coupling unification scale. In addition, one has the superpotential related Higgsino mixing parameter μ_0 with its associated sign parameter. The two radiative electroweak symmetry breaking (REWSB) conditions may then be used so as to replace B_0 and μ_0 with the Z-boson mass M_Z and $\tan\beta$, the ratio of Higgs vacuum expectation values. Similar to mSUGRA, one has other SUSY breaking scenarios like the gauge mediated SUSY breaking and models with anomaly mediated SUSY breaking (AMSB), etc. [2,3]. Apart from direct collider physics data, one has to satisfy constraints from flavor changing neutral current (FCNC), as well as flavor conserving phenomena like the anomalous magnetic moment of muon, constraints like electric dipole moments associated with CP violations, or to check whether there is a proper amount of dark matter content in R-parity conserving scenarios or proper neutrino masses in R-parity violating scenarios [1–3]. A single model is yet to be found that can adequately explain various stringent experimental results, and at the same time possesses a sufficient degree of predictiveness. However, it is always important to continue the quest of a simple model and check the degree of agreement with low energy constraints.

In this work, we pursue a predictive theory of SUSY breaking by considering a field theory on a superspace where the Grassmannian coordinates are essentially fluctuating/stochastic [8–10]. We note that in a given SUSY breaking scenario, our limitation of knowing the actual mechanism of breaking SUSY is manifested in the soft parameters. Here, in a stochastic superspace framework we assume that a manifestation of an unknown but fundamental mechanism of SUSY breaking may effectively lead to stochasticity in the Grassmannian parameters of the superspace. With a suitably chosen probability distribution, this causes a given Kähler potential and a superpotential to lead to soft breaking terms that carry signatures of the stochasticity. As we will see, the SUSY breaking is parametrized by ξ which is nothing but $1/\langle \bar{\theta}\,\bar{\theta}\rangle$, where the symbol $\langle\rangle$ refers to averaging over the Grassmannian coordinates. The other scale that is involved is Λ. Values of various soft parameters at this scale are the input parameters of the scheme. The values of the same at the electroweak scale are then obtained from these input values by using the renormalization group evolution. Considering the superpotential of MSSM, the soft terms obtained are readily recognized as the ones supplied by the externally added soft SUSY breaking terms of constrained MSSM (CMSSM) [3], except that the model is unable to produce a scalar mass soft term [8]. Reference [8] used Λ and ξ as

free parameters while analyzing the low energy signatures within MSSM. Λ was chosen between M_G to M_P, the scale of gauge coupling unification and the Planck mass scale, respectively. The model as given in Ref. [8] is called the stochastic supersymmetric model (SSM) and it is characterized by universal gaugino mass parameter $m_{\frac{1}{2}}$, universal trilinear soft SUSY breaking parameter A_0, and universal bilinear soft SUSY breaking parameter B_0, all being related to ξ, the parameter related to SUSY breaking. We note that with the bilinear soft SUSY parameter being given, $\tan\beta$, becomes a derived quantity.

However, as already mentioned, in spite of the fact that the SSM generates soft SUSY breaking terms, it produces no scalar mass soft term. Scalar masses start from zero at the scale Λ and renormalization group evolution is used to generate scalar masses at the electroweak scale M_Z. Scalar masses at M_Z severely constrain the model because typically scalars are very light. In particular, quite often sleptons turn to be the lightest supersymmetric particles or even become tachyonic over a large part of the parameter space. This is partially ameliorated when one takes the high scale Λ to be larger than the gauge coupling unification scale $M_G \sim 2 \times 10^{16}$ GeV. However, in spite of obtaining valid parameter space that would provide us with a lightest neutralino as a possible dark matter candidate in R-parity preserving framework, we must note that the low values that one obtains for the masses of the first two generation of squarks are hardly something of an advantage in view of the constraints coming from FCNC as well as those from the LHC data [6]. On the other hand, SSM has a natural advantage of being associated with a non-vanishing trilinear coupling that is favorable to produce a relatively light spectra for a given value of Higgs boson mass m_h. The recent announcement from the CMS and ATLAS Collaborations of the LHC experiment about the discovery of a Higgs-like boson at ~ 125 GeV [5] thus makes this model potentially attractive. However, as explored in Ref. [10], SSM as such is unable to accommodate such a large m_h in spite of having a built-in feature of having a nonvanishing A_0. It could at most reach 116 GeV for m_h [9,10] and the constraint due to Br($B_s \rightarrow \mu^+\mu^-$) as used in Ref. [10] was much less stringent in comparison to the same of present day [11]. Furthermore, it is also important to investigate the effect of the direct detection rate of dark matter as constrained by the recent XENON100 data [12].

In this analysis, we would like to give all the input parameters at the grand unification scale M_G, the scale at which the Standard Model gauge group, namely, $SU(3) \times SU(2) \times U(1)$, comes into existence. Any evolution above M_G would obviously demand choosing a suitable gauge group; a question that is not going to be addressed in this work. In this way we would like to avoid unknown issues arising out of a post-grand unified theory (GUT) [13] physics. However, we would rather try to meet the

phenomenological demand of confronting the issue of sleptons becoming tachyonic or avoiding scalar masses to become light in general in a *minimal modification* by considering an externally given scalar mass soft parameter m_0 as a manifestation of an additional origin of SUSY breaking. It would be useful to have the first two generations of scalar masses adequately heavy so as to overcome the FCNC related constraints and LHC data [6] on squark masses. Additionally, this will also be consistent with having the lighter Higgs boson mass (m_h) to be in the vicinity of 125 GeV. We will henceforth denote the model as Mod-SSM.

We may note that traditionally minimal versions of models of SUSY breaking have been extended for phenomenological reasons. It is also true that extending a minimal model often lowers predictiveness and may even cause partial dilution of the main motivations associated with the building of the model. For example, considering nonuniversal gaugino or scalar mass scenarios may be more suitable than CMSSM or mSUGRA so as to obtain a relatively lighter spectra in the context feasibility of exploring via LHC. Another example may be given in the context of the minimal AMSB model. As we know, a pure AMSB scenario [14] is associated with form invariance of the renormalization group equations (RGE) of scalar masses and absence of flavor violation. However, it produces tachyonic sleptons. In the minimal AMSB model [3,15], one introduces an additional common mass parameter m_0 for all the scalars of the theory. This ameliorates the tachyonic slepton problem but it is true that we sacrifice the much cherished feature of form invariance and accept some degree of flavor violations at the end. We would like to explore a nonminimal scenario of stochastic supersymmetry model, namely, Mod-SSM in this spirit. We particularly keep in mind that the stochastic supersymmetry formalism may be used not only within the MSSM framework but it may be extended to superpotentials beyond that of the MSSM [9]. The fact that the model with a minimal modification can easily accommodate the recent Higgs boson mass range by its generic feature of having a nonvanishing trilinear coupling parameter makes it further attractive. We believe that our approach of considering an additional SUSY breaking scalar mass term is justified for phenomenological reasons.

Thus, in Mod-SSM we consider the input of soft term parameters $m_{\frac{1}{2}}$, B_0, and A_0 (all are either proportional to $|\xi|$ or ξ^*) and a universal scalar mass soft parameter m_0, all being given at a suitable scale which we simply choose as the gauge coupling unification scale M_G considering the LEP data on gauge couplings could be a hint of the existence of grand unification. As in Ref. [8], we would also restrict ξ to be real, either positive or negative.

We clearly like to emphasize that the SSM framework produces no scalar mass soft term at a scale Λ. With Λ set to the gauge coupling unification scale M_G, SSM is

plagued with tachyonic sleptons. With $\Lambda > M_G$, one can avoid tachyonic scalar, but there is no scope of obtaining the currently accepted Higgs boson mass. In Mod-SSM, we consider $\Lambda = M_G$ and add the extra scalar mass term for an additional SUSY breaking effect. With the trilinear coupling parameter A_0 and the bilinear coupling parameter B_0 becoming correlated with $m_{\frac{1}{2}}$, Mod-SSM parameter space is essentially a subset of the one for the CMSSM. In Mod-SSM, the SSM inspired nonvanishing A_0 parameter is suitable for producing an appropriately large loop correction to the lighter Higgs boson mass while keeping the overall sparticle spectra relatively at a lower range. CMSSM, on the other hand, does not pinpoint with any fundamental physical distinction/motivation for such a zone of parameter space that is capable of producing the correct Higgs mass with similarly smaller sparticle mass scale. Even if we consider the stochastic superspace model in its original form as only a toy idea, we believe that it may be worthwhile to explore it with minimal modifications in regard to the current Higgs boson mass, as well as other phenomenological constraints.

II. STOCHASTIC GRASSMANNIAN COORDINATES AND SUSY BREAKING

As seen in Ref. [8], we consider an $N = 1$ superspace where the Grassmannian coordinates θ and $\bar{\theta}$ are taken to be stochastic in nature. One starts with identifying the terms involving superfields in the superpotential and the kinetic energy terms that could be used to construct the SUSY invariant Lagrangian density for a given model. Each term is then multiplied with a probability distribution function $\mathcal{P}(\theta, \bar{\theta})$ and integrated over the Grassmannian coordinates appropriately. $\mathcal{P}(\theta, \bar{\theta})$ can be expanded into terms involving θ and $\bar{\theta}$ which obviously has a finite number of terms because of the Grassmannian nature of θ and $\bar{\theta}$. One then imposes the normalization condition $\int d^2\theta d^2\bar{\theta} \mathcal{P}(\theta, \bar{\theta}) = 1$ and vanishing of Lorentz nonscalar moments like $\langle\theta\rangle$, $\langle\bar{\theta}\rangle$, $\langle\theta\bar{\theta}\rangle$, $\langle\theta^2\bar{\theta}\rangle$, and $\langle\theta\bar{\theta}^2\rangle$. The stochasticity parameter ξ is defined as $\langle\theta\theta\rangle = 1/\xi^*$. Here, ξ is a complex parameter with mass dimension unity. As computed in Ref. [8], and as worked out in this analysis explicitly in the Appendix, the above leads to the following Hermitian probability distribution:

$$\mathcal{P}(\theta, \bar{\theta})|\xi|^2 = \tilde{\mathcal{P}}(\theta, \bar{\theta})$$
$$= 1 + \xi^*(\theta\theta) + \xi(\bar{\theta}\,\bar{\theta}) + |\xi|^2(\theta\theta)(\bar{\theta}\,\bar{\theta}).$$
$$(1)$$

For the simple case of a Wess-Zumino type of scenario [3] where the kinetic term is obtained from $\Phi^\dagger\Phi$ and the superpotential is given as $W = \frac{1}{2}m\Phi^2 + \frac{1}{3}h\Phi^3$, where Φ is a chiral superfield, one finds that the effect of stochasticity as described above leads to the following SUSY breaking term:

$$- L_{\text{soft}} = \frac{1}{2} \xi^* m \phi^2 + \frac{2}{3} \xi^* h \phi^3 + \text{H.c.} \quad (2)$$

Applying the stochasticity idea to the superpotential of MSSM, along with considering the effect on the gauge kinetic energy function, the above formalism leads to the following tree level soft SUSY breaking parameters to be given at the high scale Λ:

(a) universal gaugino mass parameter $m_{\frac{1}{2}} = \frac{1}{2}|\xi|$,

(b) universal trilinear soft parameter $A_0 = 2\xi^*$,

(c) universal bilinear Higgs soft parameter $B_0 = \xi^*$.

Recall that there is no scalar mass soft SUSY breaking term in SSM. For convenience, we take ξ to be a real positive number with an additional input $\text{sign}(\xi)$. If we count the universal gaugino mass parameter $m_{\frac{1}{2}}$ as the independent parameter, we have

$$A_0 = \text{sign}(\xi).4m_{\frac{1}{2}}, \qquad B_0 = \text{sign}(\xi).2m_{\frac{1}{2}}. \quad (3)$$

As has already been discussed before, we introduce a nonvanishing scalar mass parameter m_0 and fix Λ at M_G.[1] Thus with the above extension, the input quantities for the stochastic SUSY model are

$$m_{\frac{1}{2}}, \qquad m_0, \qquad \text{sign}(\mu), \quad \text{and} \quad \text{sign}(\xi).$$

We note that the model quite naturally is associated with nonvanishing trilinear soft breaking terms. As we will see, this is quite interesting in view of the recent LHC announcement for the Higgs mass range centering around 125 GeV [5]. In this analysis, we will discuss only the case of $\xi < 0$ because the other sign of ξ does not produce a spectra compatible with the dark matter relic density constraint.

The requirement of the REWSB then results in the following relations at the electroweak scale:

$$\mu^2 = -\frac{1}{2} M_Z^2 + \frac{m_{H_D}^2 - m_{H_U}^2 \tan^2\beta}{\tan^2\beta - 1} + \frac{\Sigma_1 - \Sigma_2 \tan^2\beta}{\tan^2\beta - 1}, \quad (4)$$

and

$$\sin 2\beta = 2B\mu/(m_{H_D}^2 + m_{H_U}^2 + 2\mu^2 + \Sigma_1 + \Sigma_2), \quad (5)$$

where Σ_i denote the one-loop corrections [16,17]. Here, B refers to the value of bilinear Higgs coupling at the electroweak scale which has to be consistent with its given value B_0 at M_G. B_0 is determined via $m_{\frac{1}{2}}$ apart from a sign of the stochasticity parameter as mentioned before. Consequently, $\tan\beta$ is a derived quantity in the model. B_0 at the scale M_G and B at the electroweak scale are
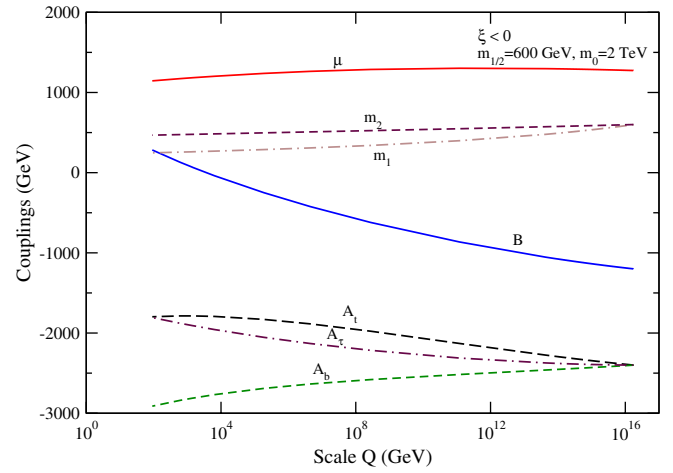

FIG. 1 (color online). Evolution of a few relevant couplings for a specimen input of $m_{1/2} = 600$ GeV, $m_0 = 2$ TeV, and $\mu > 0$. With $\xi < 0$, one has $B_0 = -2m_{1/2}$ and $A_0 = -4m_{1/2}$. For a valid parameter point within the model with $\mu > 0$, we require $B = B(M_Z) > 0$ a necessity in order to have a positive $\sin 2\beta$ from Eq. (5).

connected via the following RGE written here at the one-loop level:

$$\frac{dB}{dt} = \left(3\tilde{\alpha}_2 \tilde{m}_2 + \frac{3}{5} \tilde{\alpha}_1 \tilde{m}_1\right) + (3Y_t A_t + 3Y_b A_b + Y_\tau A_\tau), \quad (6)$$

where $t = ln(M_G^2/Q^2)$ with Q being the renormalization scale. $\tilde{\alpha}_i = \alpha_i/(4\pi)$ for $i = 1, 2, 3$ refer to scaled gauge coupling constants (with $\alpha_1 = \frac{5}{3} \alpha_Y$) and $\tilde{m}_i$ for $i = 1, 2, 3$ are the running gaugino masses. Y_i are the squared Yukawa couplings, e.g., $Y_t \equiv y_t^2/(4\pi)^2$ where y_t is the top Yukawa coupling. In this analysis, the value of $\tan\beta$ is determined via Eqs. (4)–(6) along with $B_0 = \text{sign}(\xi)2m_{\frac{1}{2}}$ at the scale M_G. We use SuSpect [18] for solving the RGEs and obtaining the spectra. The code takes $\tan\beta$ as an input quantity. Hence, we implement a self-consistent method of solution that starts from a guess value of $\tan\beta$ resulting into a $B(M_G)$ that in general would not agree with the input of B_0. Use of a Newton-Raphson root finding scheme ensures a fast convergence toward the correct value of $\tan\beta$ when $B(M_G)$ matches with the input of B_0. Here we stress that we do not encounter any parameter point with multiple values of $\tan\beta$ in our analysis.[2]

III. RESULTS

The fact that the model has $\tan\beta$ as a derived quantity necessitates studying the behavior of the evolution of the bilinear Higgs parameter B. Figure 1 shows the evolution of a few relevant couplings for a specimen input of

---

[1]We note that a vanishing scalar mass parameter at a post-GUT scale with RG evolution corresponding to an appropriate gauge group would indeed generate nonvanishing scalar mass terms at the unification scale M_G [13].

[2]See Ref. [19] for such a general possibility in REWSB where B_0 is given as an input.
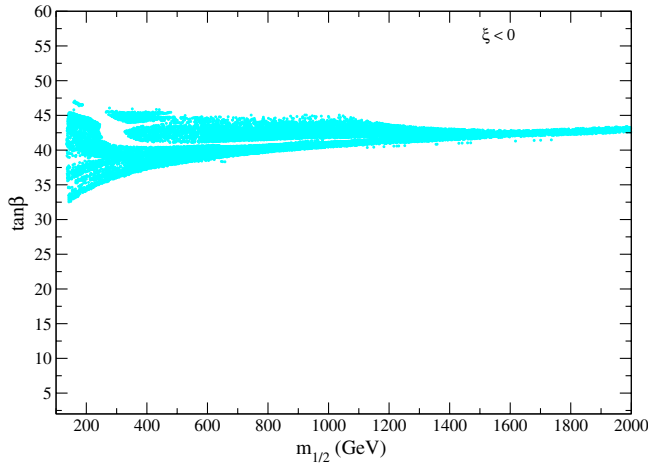
FIG. 2 (color online).   Scatter plot of parameter points in the $m_{\frac{1}{2}} - \tan\beta$ plane when m_0 and $m_{\frac{1}{2}}$ are scanned up to 7 and 2 TeV, respectively. Here we only consider the validity of REWSB constraints of Eqs. (4) and (5). The values of $\tan\beta$ that satisfy the REWSB constraint vary from 32 to 48. The white region inside the shaded (blue-green) region corresponds to invalid parameter points where no consistent solution satisfying REWSB could be found even after trying with a large number of iterations. The spread of $\tan\beta$ for a given $m_{\frac{1}{2}}$ arises from variation of m_0. This spread decreases as $m_{\frac{1}{2}}$ becomes larger because of decreased sensitivity on m_0 while satisfying REWSB.

$m_{1/2} = 600$ GeV, $m_0 = 2$ TeV, and $\mu > 0$ in Mod-SSM. For $\xi < 0$, both B_0 as well as A_0 are negative, namely, $B_0 = -2m_{1/2}$ and $A_0 = -4m_{1/2}$. For a valid parameter point within the model with $\mu > 0$, we require $B(M_Z) > 0$ a necessity in order to have a positive $\sin 2\beta$ from Eq. (5).[3] We note that the denominator in the right-hand side of Eq. (5) is the square of pseudoscalar Higgs mass which needs to be positive. The fact that B is originally negative at M_G and has to change to a positive value at M_Z puts a strong constraint on the parameter space of the model. Numerically, this results in $\tan\beta$ assuming large values. In regard to the evolution of A parameters, we defer our discussion until Fig. 3.

Figure 2 shows a scatter plot of parameter points in the $\tan\beta - m_{\frac{1}{2}}$ plane that satisfy the REWSB constraints of Eqs. (4) and (5). As m_0 is varied up to 7 TeV and $m_{\frac{1}{2}}$ up to 2 TeV, $\tan\beta$ is seen to have a range of 32 to 48. The spread of $\tan\beta$ for a given $m_{\frac{1}{2}}$ arises from variation of m_0. For smaller values of $m_{\frac{1}{2}}$, there is a larger dependence of m_0 on $\tan\beta$. Hence there is a larger spread of $\tan\beta$ when m_0 is varied. For larger $m_{\frac{1}{2}}$, the valid solution of $\tan\beta$ has a lesser dependence on m_0. Hence, for larger $m_{\frac{1}{2}}$ values the spread in values of $\tan\beta$ decreases. The white region embedded within the blue-green area corresponds to parameter points

_____

[3]$\tan\beta$ which is the ratio of two vacuum expectation values is positive. Hence, $\sin 2\beta = \frac{2\tan\beta}{(1+\tan^2\beta)}$ is also positive.
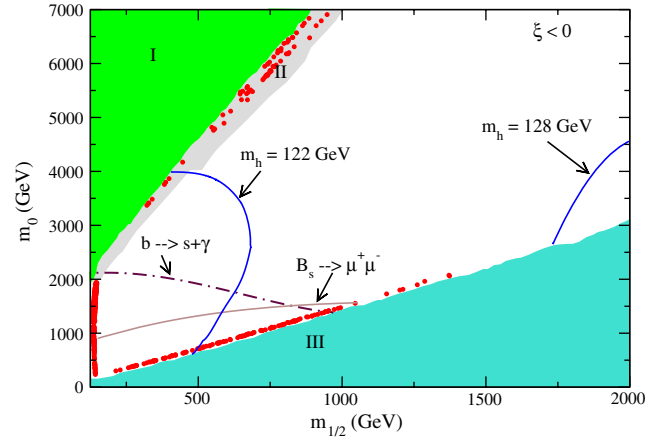

FIG. 3 (color online).   Constraints shown in the $m_{\frac{1}{2}} - m_0$ plane for $\mu > 0$ and $\xi < 0$ for Mod-SSM. A_0 and B_0 in the model satisfy $A_0 = \text{sign}(\xi).4m_{\frac{1}{2}}$ and $B_0 = \text{sign}(\xi).2m_{\frac{1}{2}}$. $\tan\beta$ becomes a derived quantity that varies between 32 and 48. The Higgs boson limits are shown as two solid blue lines. $\text{Br}(b \rightarrow s\gamma)$ limit is shown as a maroon dot-dashed line. The lower part corresponds to discarded region via Eq. (9) where the branching ratio goes below the lower limit of the constraint. $\text{Br}(B_s \rightarrow \mu^+\mu^-)$ limit is shown as a brown solid line of which the lower region exceeds the upper limit of Eq. (10). The top green region (I) corresponds to discarded zone via REWSB. The bottom blue-green region (III) refers to the zone where stau becomes LSP or tachyonic. The gray region (II) has discontinuous patches of valid parameter zones, the details of which are mentioned in the text. Red points/areas falling in region II satisfy the WMAP-7 data only for the upper limit of Eq. (12). Typically, the red points bordering region III and some part of the extreme left red points (for very small $m_{\frac{1}{2}}$) satisfy both the upper and the lower limits of Eq. (12).

with no valid solution satisfying REWSB. Typically, the two REWSB conditions are nonlinear in nature with respect to $m_{\frac{1}{2}}$ and m_0. The code unsuccessfully tries with a large number of iterations to find consistent μ^2 and m_A^2 solutions for parameter points within the white region (for all m_0). Thus, a lack of a valid $\tan\beta$ for any value of m_0 results in the above white region. It is worth mentioning that the range of valid $\tan\beta$ is much larger for the case of $\xi > 0$ where B stays positive throughout the range from M_G to M_Z. This is unlike the case of $\xi < 0$ under discussion, where B is negative at M_G and necessarily has to become positive at the electroweak scale, thus adding stringency to $\tan\beta$ in its range. However, as already mentioned, we will not discuss the case of $\xi > 0$ further because of the resulting overabundance of dark matter for the entire parameter space for this sign of ξ. We will study now the effect of low energy constraints particularly in the context of the recent discovery of the Higgs-like boson [5]. Figure 3 shows the result in the $m_{\frac{1}{2}} - m_0$ plane for $\xi < 0$. A sufficiently nonvanishing A_t (see, for example, Fig. 1) helps in producing a large loop correction to the lighter CP-even Higgs boson h. This can be understood by

looking at the expression for the dominant part of loop correction to the Higgs boson mass coming from the top-stop sector [20–22]

$$\Delta m_h^2 = \frac{3\bar{m}_t^4}{2\pi^2 v^2 \sin^2\beta} \left[\log \frac{M_S^2}{\bar{m}_t^2} + \frac{X_t^2}{2M_S^2} \left(1 - \frac{X_t^2}{6M_S^2} \right) \right]. \quad (7)$$

Here, $M_S = \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}$, $X_t = A_t - \mu \cot\beta$, $v = 246$ GeV, and $\bar{m}_t$ is the running top-quark mass that also takes into account QCD and electroweak corrections. The loop correction is maximized if $X_t = \sqrt{6}M_S$. Clearly, a nonvanishing A_0 can be useful to increase Δm_h^2 so that m_h reaches the LHC specified zone without a need to push up the average sparticle mass scale by a large amount. As particularly mentioned in Ref. [8], the SSM is additionally attractive in this context since it naturally possesses a nonvanishing and large $|A_0|$. We further note that in this model, the lighter Higgs boson h has couplings similar to those in the Standard Model because the CP-odd Higgs boson mass (m_A) is in the decoupling zone [23]. The ATLAS and CMS results for the possible Higgs boson masses are $126.0 \pm 0.4(\text{stat}) \pm 0.4(\text{syst})$ GeV and $125.3 \pm 0.4(\text{stat}) \pm 0.5(\text{syst})$ GeV, respectively [5]. In regard to MSSM light Higgs boson mass, we note that there is about a 3 GeV uncertainty arising out of uncertainties in the top-quark mass, renormalization scheme, as well as scale dependence and uncertainties in higher order loop corrections up to three loop [24–28]. Hence, in this analysis we consider the following limits for m_h:

$$122 \text{ GeV} < m_h < 128 \text{ GeV}. \quad (8)$$

We will outline the other relevant limits used in this analysis. A SUSY model parameter space finds a strong constraint from $\text{Br}(b \rightarrow s\gamma)$. In the SM, the principal contribution that almost saturates the experimental value comes from the loop comprised of top quark and W boson [29]. In the MSSM, principal contributions arise from loops containing top quark and charged Higgs bosons, and the same containing top squarks and charginos [30]. The chargino loop contributions are proportional to $A_t\mu$, and this may cause cancellations or enhancements between the principal terms of the MSSM contribution depending on the sign of $A_t\mu$. Similar to mSUGRA, both the SSM and Mod-SSM with $\xi < 0$ also typically have $A_t < 0$. With $\mu > 0$, this primarily means cancellation between the chargino and the charged Higgs contributions. This leads to a valid region for $\text{Br}(b \rightarrow s\gamma)$ for a larger sparticle mass scale compared to the case of $\mu < 0$. $\text{Br}(b \rightarrow s\gamma)$ constraint thus favors the positive sign of μ by allowing larger areas of parameter space. We consider the experimental value $\text{Br}(b \rightarrow s\gamma) = (355 \pm 24 \pm 9) \times 10^{-6}$ [31]. This results in the following 3σ level zone as used in this analysis:

$$2.78 \times 10^{-4} < \text{Br}(b \rightarrow s\gamma) < 4.32 \times 10^{-4}. \quad (9)$$

The above constraint is displayed as a maroon dot-dashed line. The left region of this line would be a discarded zone.

Next, the fact that the stochastic model with $\xi < 0$ selects appreciably large values for $\tan\beta$ necessitates checking the $B_s \rightarrow \mu^+\mu^-$ limit. This is required because $B_s \rightarrow \mu^+\mu^-$ increases with $\tan\beta$ as $\tan^6\beta$ and decreases with increase in m_A, the mass of pseudoscalar Higgs boson as m_A^{-4} [32]. We use the recent experimental limit from LHCb [11]: $\text{Br}(B_s \rightarrow \mu^+\mu^-)_{\text{exp}} = (3.2^{+1.4}_{-1.2}(\text{stat.})^{+0.5}_{-0.3}(\text{syst.})) \times 10^{-9}$. This is contrasted with the SM evaluation $\text{Br}(B_s \rightarrow \mu^+\mu^-)_{\text{SM}} = (3.23 \pm 0.27) \times 10^{-9}$ [33]. As in Ref. [34], combining the errors of the LHCb data along with that of the SM result, one finds the following:

$$0.67 \times 10^{-9} < \text{Br}(B_s \rightarrow \mu^+\mu^-) < 6.22 \times 10^{-9}. \quad (10)$$

The upper limit of $\text{Br}(B_s \rightarrow \mu^+\mu^-)$ is shown as a brown solid line going across Fig. 3. Parameter values in the region below this curve lead to values of $\text{Br}(B_s \rightarrow \mu^+\mu^-)$ higher than the above limit.

We also compute $\text{Br}(B \rightarrow \tau\nu_\tau)$ in this analysis. The SUSY contribution to $\text{Br}(B \rightarrow \tau\nu_\tau)$ is typically effective for large $\tan\beta$ and small charged Higgs boson mass scenarios [35]. The experimental data from *BABAR* [36] reads $\text{Br}(B^+ \rightarrow \tau^+\nu_\tau) = (1.83^{+0.53}_{-0.49}(\text{stat.}) \pm 0.24(\text{syst.})) \times 10^{-4}$. The recent result from Belle [37] for $B^- \rightarrow \tau^-\bar{\nu}_\tau$ that used the hadronic tagging method is given by $\text{Br}(B^- \rightarrow \tau^-\bar{\nu}_\tau) = (0.72^{+0.27}_{-0.25}(\text{stat.}) \pm 0.11(\text{syst.})) \times 10^{-4}$. The same branching ratio from Belle extracted by using a semileptonic tagging method is $(1.54^{+0.38}_{-0.37}(\text{stat.})^{+0.29}_{-0.31}(\text{syst.})) \times 10^{-4}$ [38]. We use Ref. [39] for the result of averaging all of the recent Belle and *BABAR* data which is $\text{Br}(B \rightarrow \tau\nu)_{\text{exp}} = (1.16 \pm 0.22) \times 10^{-4}$. The SM result strongly depends on the CKM element $|V_{ub}|$ and the B-meson decay constant. We use $\text{Br}(B \rightarrow \tau\nu)_{\text{SM}} = (0.97 \pm 0.22) \times 10^{-4}$ [39]. Using the above theoretical and experimental errors appropriately, we obtain the following:

$$R_{(B \rightarrow \tau\nu_\tau)} = \frac{\text{Br}(B \rightarrow \tau\nu_\tau)_{\text{SUSY}}}{\text{Br}(B \rightarrow \tau\nu_\tau)_{\text{SM}}} = 1.21 \pm 0.30. \quad (11)$$

This translates into $0.31 < R_{(B \rightarrow \tau\nu_\tau)} < 2.10$ at 3σ. Here, $\text{Br}(B \rightarrow \tau\nu_\tau)_{\text{SUSY}}$ denotes the branching ratio in a SUSY framework, of course including the SM contribution. In general, we find that the model parameter space of Mod-SSM is not constrained by $B \rightarrow \tau\nu_\tau$ since charged Higgs bosons are sufficiently heavy.

We have not, however, included the constraint from muon $g - 2$ in this analysis, considering the tension arising out of large deviation from the SM value, uncertainty in hadronic contribution evaluations, and accommodating SUSY models in view of the LHC sparticle mass lower limits [40].

We now explore the cosmological constraint for neutralino dark matter relic density [41]. At 3σ, the WMAP-7 data [42] are considered as shown below

$$0.094 < \Omega_{\tilde{\chi}_1^0} h^2 < 0.128. \quad (12)$$

The conclusions in regard to the relic density constraint is additionally found to be sensitive on the top-quark mass in this model. We divide the dark matter analysis into two parts depending on (a) the top-quark pole mass set at 173.3 GeV, and (b) using a spread of top-quark pole mass within its range $m_t = 173.3 \pm 2.8$ GeV following the result of the recent analysis performed in Ref. [43]. In this context, we note that the experimental value as measured by the CDF and D0 Collaborations of Tevatron is $m_t^{\mathrm{exp}} = 173.2 \pm 0.9$ GeV [44].[4]

A. Analysis with $m_t = 173.3$ GeV: Underabundant LSP

The lightest neutralino, the LSP of the model, is typically highly bino dominated except in a few regions where the Higgsino mixing parameter μ turns out to be small. The parameter points within all of the white region in Fig. 3 have bino-dominated LSP. At this point we note that the implementation of the REWSB conditions as manifest in Eqs. (4) and (5) has to be done by keeping in mind (i) positivity of $\sin 2\beta$, (ii) positivity of μ^2, and (iii) the requirement of B_0, A_0 being related to $m_{\frac{1}{2}}$ as given by Eq. (3), as well as requiring that the lighter chargino mass lower limit is respected. All these requirements lead the top green shaded region (labeled as I) to be a discarded zone. On the other hand, the gray shaded region (II) has discontinuous zones of valid parameter points shown in red. The red points have considerably small values of μ, thus giving the LSP a large degree of Higgsino mixing. Apart from the red points, there are no solutions in the gray areas of region II. The need to satisfy the point (i) as explained above along with the requirement to satisfy the condition (iii) (implemented via a Newton-Raphson method of finding the correct $\tan\beta$) stringently negates the existence of solution zones within region II. As a result, either there are solutions with appreciably small μ or no solution at all within this region. This on the other hand leads to a large amount of $\tilde{\chi}_1^0 - \tilde{\chi}_1^\pm$ coannihilation. Such degrees of coannihilations indeed cause the LSP to be only a subdominant component of DM. Thus, in this part of the analysis, we consider the possibility of an underabundant dark matter candidate and ignore the lower limit of the WMAP-7 data. Typically we see that the relic density falls below the lower limit of Eq. (12) in region II by an order of magnitude. Such underabundant LSP scenarios have been discussed in several works [45].

The lower shaded region III is disallowed as the mass of the stau ($\tilde{\tau}_1$) turns negative or it is the LSP. Typically the red strip near region III refers to the LSP-stau coannihilation[5] zone where the relic density can be consistent with both the upper and lower limits of Eq. (12). Quite naturally,

the coannihilation may be stronger, and this would additionally produce some underabundant DM points for this zone. Finally, the leftmost red region (with very small $m_{\frac{1}{2}}$) satisfying WMAP-7 data is discarded by all other constraints. We note that only a small region satisfying the Higgs mass bound is discarded via Br($b \to s\gamma$). On the other hand, the Higgs mass bound line of 122 GeV supersedes the constraint imposed by the recent data on inclusive search for SUSY by the ATLAS experiment [6]. The most potent constraint to eliminate a large region of parameter space with m_0 up to 1.5 TeV or so is due to Br($B_s \to \mu^+\mu^-$) data Eq. (10). The constraint is effective simply because of the large values of $\tan\beta$ involved in the model. The discarded part of parameter space via the above constraint includes a large zone that satisfies the dark matter limit via LSP-stau coannihilation.

We will now describe the spin-independent direct detection scattering cross section for scattering of LSP with proton. The scalar cross section depends on t-channel Higgs exchange diagrams and s-channel squark diagrams. Unless the squark masses are close to that of the LSP, the Higgs exchange diagrams dominate [47]. We note that for the cases of parameter points with $\Omega_{\tilde{\chi}} h^2 < (\Omega_{\mathrm{CDM}} h^2)_{\mathrm{min}}$, where $(\Omega_{\mathrm{CDM}} h^2)_{\mathrm{min}}$ refers to the lower limit of Eq. (12), one must appropriately include the fraction of local DM density contributed by the specific candidate of DM under discussion while evaluating the event rate.

This translates into multiplying $\sigma^{\mathrm{SI}}_{p\tilde{\chi}_1^0}$ for such underabundant scenarios by ρ_χ/ρ_0. Here, ρ_χ is the actual DM density contributed by the specific DM candidate contributing to ρ_0 where the latter is the local dark matter density. We thus use $\rho_\chi = \rho_0 \zeta$ where $\zeta = \Omega_{\tilde{\chi}} h^2/(\Omega_{\mathrm{CDM}} h^2)_{\mathrm{min}}$. On the other hand, ζ is simply 1 for abundant or overabundant dark matter cases. Thus, we conveniently define $\zeta = \min\{1, \Omega_{\tilde{\chi}} h^2/(\Omega_{\mathrm{CDM}} h^2)_{\mathrm{min}}\}$ [48]. Figure 4 shows the rescaled cross section as computed via micrOMEGAs version 2.4 [49]. We wish to emphasize that while some region of parameter space where the LSP typically has a large degree of Higgsino component is eliminated via XENON100 data, as announced in the summer of 2012 [12], a large section of parameter space remains to be explored via future direct detection of DM experiments. This consists of both types of coannihilation zones; namely, the chargino as well as the stau coannihilation zones. We must also keep in mind the issue of theoretical uncertainties, particularly the hadronic uncertainties in evaluating $\sigma^{\mathrm{SI}}_{p\tilde{\chi}_1^0}$. The strangeness content of nucleon finds a large reduction in the evaluation of relevant couplings via lattice calculations [50]. This is not incorporated in our computation while using micrOMEGAs to calculate the cross section. Thus, the above itself will cause a reduction of $\sigma^{\mathrm{SI}}_{p\tilde{\chi}_1^0}$ by almost an order of magnitude. There is also an appreciable amount of uncertainty of the local dark matter density [51]. All these points need

---

[4]For bottom quark mass we have used $m_b^{\overline{\mathrm{MS}}}(m_b) = 4.19$ GeV.
[5]See, for example, works in Ref. [46] for various annihilation processes in relation to SUSY parameter space in general.
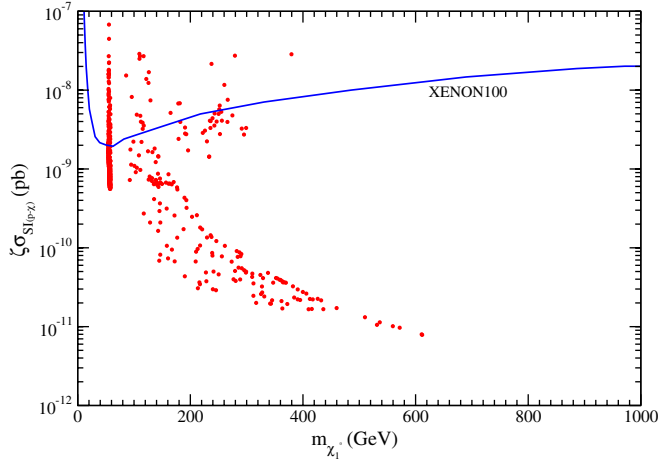
FIG. 4 (color online). Scaled spin-independent $\tilde{\chi}_1^0 - p$ scattering cross section vs LSP mass. The scaling factor is given as $\zeta = \Omega_{\tilde{\chi}} h^2/(\Omega_{\text{CDM}} h^2)_{\text{min}}$, where $(\Omega_{\text{CDM}} h^2)_{\text{min}}$ refers to the lower limit of Eq. (12).
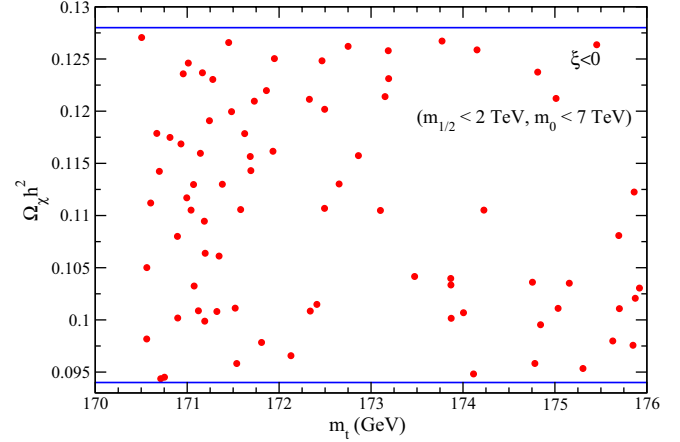

FIG. 5 (color online). Relic density satisfied points shown in red that fall within the lower and upper limits of WMAP-7 data for the neutralino relic density, when m_t is varied by 2.8 GeV on either side of 173.3 GeV. Here $m_{\frac{1}{2}}$ and m_0 are scanned up to 2 and 7 TeV, respectively, for $\mu > 0$ and $\xi < 0$. The two blue lines are the WMAP-7 limits of Eq. (12). The value of $m_t = 173.3$ GeV as used in Fig. 3 visibly falls in a disfavored zone in the context of obtaining the correct relic density, particularly toward the lower limit of Eq. (12).

to be kept in mind while evaluating the implications of Fig. 4 for SUSY models.

B. Analysis with $m_t = 173.3 \pm 2.8$ GeV: LSP of right abundance

It is to be noted that in the gray area (shown as region II) of Fig. 3, the existence of a valid solution depends very critically on the parameters of the model. It was found that either (i) we obtain a very small μ (barely satisfying the lighter chargino mass lower limit) in the gray area that would only provide us with extreme coannihilation between lighter chargino and LSP leading to underabundance of DM, or (ii) we find no valid solution at all. The sensitivity arises from the stringency of satisfying REWSB on the parameter space for this region. In other words, for a given $m_{\frac{1}{2}}$, a small change in assumed m_0 for a valid parameter point would produce a small change in $\tan\beta$ in the *white* region producing most probably another valid parameter point. However, the change may not be allowed via REWSB, particularly Eq. (5) if the parameter point is considered in the *gray* region. Equation (5) means $\sin 2\beta$ needs to be a positive quantity less than unity. This typically becomes a severe constraint even if the condition of satisfying the lighter chargino mass lower limit is met. With the top-quark mass having a strong influence on REWSB, it may be useful to investigate whether varying m_t may extract newer valid points within the gray region that would have a suitable μ so as to satisfy a well-tempered [52] LSP situation. This will then have the right abundance of dark matter satisfying both the lower and upper limits of DM in Eq. (12).

The present experimental data on top-quark mass read $m_t^{\text{exp}} = 173.2 \pm 0.9$ GeV [44]. Recently, Ref. [43] predicted the pole mass of top quark to be $m_t^{\text{pole}} = 173.3 \pm 2.8$ GeV. The analysis used the next-to-next-to-leading

order (NNLO) in the QCD prediction of the inclusive $pp \rightarrow t\bar{t} + X$ cross section and the Tevatron and LHC data of the the same cross section. The comparison between the experimental and theoretical results helped extract the top-quark mass in the modified minimal subtraction ($\overline{\text{MS}}$) scheme. This was then used to compute the pole mass m_t^{pole}. We now extend our analysis by investigating the effect of varying top-quark pole mass within the above range ($m_t^{\text{pole}} = 173.3 \pm 2.8$ GeV) on the solution space of the model. Indeed, we will see that even with a variation of 0.9 GeV, the range of experimental error would be enough to have a substantial effect on the conclusions.

Figure 5 shows the scattered points that satisfy the WMAP-7 given range of relic density for the aforesaid variation of $m_t \equiv m_t^{\text{pole}}$. Here, $m_{\frac{1}{2}}$ and m_0 are varied up to 2 and 7 TeV, respectively. We see from this figure that the central value of $m_t (= 173.3)$ GeV as considered in Fig. 3 is indeed the one which has the least amount of possibility to satisfy the WMAP-7 data. The occurrence of points which satisfies WMAP-7 data is particularly rare around this value of m_t, near the lower part of the limit of relic density. We already found that the gray region (region II) of Fig. 3 is a sensitive zone because of REWSB, where μ can be quite small. For such small values of μ, one can only expect a large degree of $\tilde{\chi}_1^0 - \tilde{\chi}_1^\pm$ coannihilation which results into very small relic density. The latter goes below the lower limit of Eq. (12). Hence no red point exists near the bottom blue line of Fig. 5 for this value of m_t. On the contrary, a value of m_t less than 1 GeV from the central
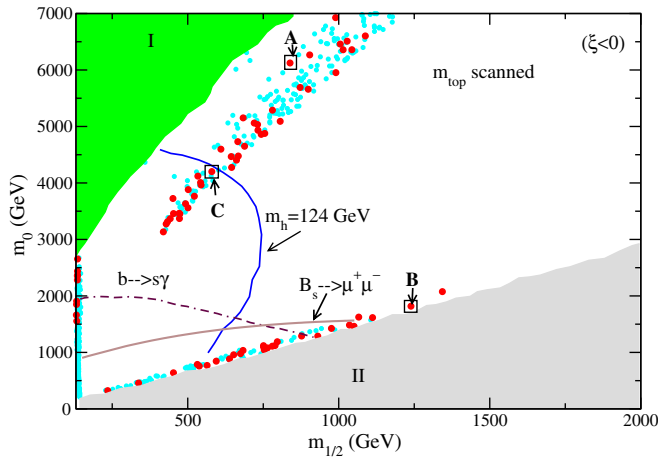
FIG. 6 (color online). Scattered points in the $m_{\frac{1}{2}} - m_0$ plane for $\mu > 0$ and $\xi < 0$ when top-quark mass is varied by 2.8 GeV on either side of 173.3 GeV. Region I is the same as that in Fig. 3 except that it is now a discarded zone for all values of m_t within its limit. The region II is discarded via stau turning the LSP or turning itself tachyonic for all m_t. The constraints from Br($b \to s\gamma$) and Br($B_s \to \mu^+\mu^-$) are as shown. The lines show the boundary of purely discarded zones irrespective of variation of m_t within its range. There is no region toward the left of the blue line labeled as $m_h = 124$ GeV for which m_h may become larger than 124 GeV irrespective of values of m_t. Three benchmark points A, B, and C are shown corresponding to Table I.

value, which is only within the experimental error of m_t^{exp}, would cause to have a LSP with correct abundance for DM. The most favored zone for m_t, however, would be from 171 to 172 GeV for satisfying the relic density limits. In fact, a reduced degree of sensitivity for satisfying REWSB is the reason for obtaining a well-tempered LSP while considering a top-quark mass little away from the central value.

Figure 6 shows the effect of scanning the top-quark mass on the $m_{\frac{1}{2}} - m_0$ plane. Region I shown in green is a disallowed area of parameter space via REWSB similar to Fig. 3, except that it is now an invalid region for all values of m_t within its limit. Thus, this region is smaller in extension than the corresponding region of Fig. 3. The region II is disallowed because of stau turning tachyonic or the least massive. The constraints from Br($b \to s\gamma$) and Br($B_s \to \mu^+\mu^-$) are as shown. The lines denote the boundary of purely discarded zones irrespective of variation of m_t within its range. The blue line for $m_h = 124$ GeV means that $m_h < 124$ GeV for all the region left of the line irrespective of the value of m_t within the range. The ATLAS specified limit [6] for squarks also falls well within this left zone. The neutralino relic density satisfied areas are below region I and above region II. The *red* points satisfy both the limits of the WMAP-7 data, and thus correspond to having the right degree of abundance of DM. On the other hand, we have also shown *blue-green* points that only satisfy the upper limit of the WMAP-7 data. In this part of the analysis, we consider the
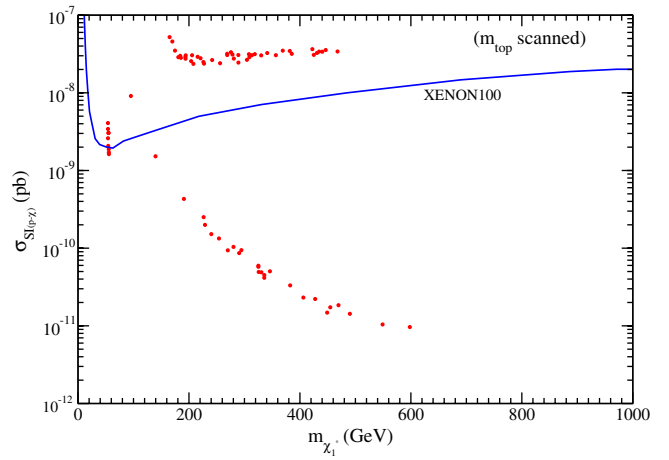

FIG. 7 (color online). Spin-independent LSP-proton scattering cross section vs LSP mass when m_t is scanned along with $m_{\frac{1}{2}}$ and m_0 for $\mu > 0$ and $\xi < 0$. Only WMAP-7 satisfied points (for both the lower and the upper limit) are shown along with the XENON100 exclusion limit.

LSP to have the correct abundance so as to be a unique candidate for DM.

Finally, we show the effect of varying m_t on the spin-independent LSP-proton scattering cross section in Fig. 7. Only the WMAP-7 satisfied points are shown (in red). Considering an order of magnitude of uncertainty (reduction), primarily because of the issue of strangeness content of nucleon as well as astrophysical uncertainties as mentioned before, we believe that the recent XENON100 data still can accommodate Mod-SSM even while considering the LSP as a unique candidate of DM.

Table I shows three benchmark points of the model. The top-quark mass m_t is as shown for each of the cases. Points A and B correspond to values of m_t which are entirely within the experimental error. The points A and C correspond to the upper region of Fig. 6, and these two points correspond to the hyperbolic branch (HB)/focus point (FP) zone [53,54]. These points are associated with a large degree of $\tilde{\chi}_1^0 - \tilde{\chi}_1^\pm$ coannihilation. The degree of agreement between the desired value of B_0 and the one obtained via the Newton-Raphson iteration may be seen from the fifth and sixth rows. We allowed a maximum deviation of 1% in the iterative procedure while scanning the parameter space. A majority of the points are consistent within 0.1% in this regard. The charginos and the neutralinos are relatively lighter for points A and C in comparison to those in point B. The scalar particles, on the other hand, are relatively lighter for point B. Undoubtedly, like many SUSY models analyzed or reanalyzed after the Higgs boson discovery, the spectra is on the heavier side. The fact that A_0 is nonvanishing and adequately large helps in reducing the average sparticle mass to a great extent compared to vanishing A_0 scenarios satisfying the current lighter Higgs boson limit. The m_h for point C, however, goes below the assumed limit of Eq. (8). However, we still believe

TABLE I. Spectra of three specimen parameter points A, B, and C as shown in Fig. 6. Results with marginal deviation from the assumed limits are shown in italics (red). Muon $g-2$ is not imposed as a constraint in this analysis. $B_0 = B(M_G)$ has two entries. The first one is the desired value while the second one is the value obtained with a suitable $\tan\beta$ as found from the Newton-Raphson root finding scheme. See text for further details.

Parameter	A	B	C
m_t	173.10	173.87	171.58
$m_{1/2}$	838.78	1239.16	579.69
m_0	6123.75	1817.69	4200.55
$(A_0 = -4m_{1/2})$	-3355.13	-4956.64	-2318.75
$(B_0 = -2m_{1/2})$	-1677.56	-2478.32	-1159.37
B_0 (*as output*)	-1683.56	-2478.32	-1160.27
$\tan\beta$ (*as output*)	45.86	40.92	45.11
$\mathrm{sgn}(\mu)$	1	1	1
μ	403.86	2508.85	310.43
$m_{\tilde{g}}$	2145.53	2727.64	1525.80
$m_{\tilde{u}_L}$	6247.84	2994.87	4292.70
$m_{\tilde{t}_1}, m_{\tilde{t}_2}$	3758.76, 4376.60	1333.10, 2078.60	2587.56, 3026.20
$m_{\tilde{b}_1}, m_{\tilde{b}_2}$	4397.10, 4886.58	2054.58, 2339.25	3037.22, 3381.67
$m_{\tilde{e}_L}, \tilde{\nu}_e$	6119.53, 6119.05	1983.72, 1982.21	4197.61, 4196.89
$m_{\tilde{\tau}_1}, m_{\tilde{\nu}_\tau}$	4750.43, 5482.91	549.62, 1536.41	3281.31, 3770.25
$m_{\tilde{\chi}_1^\pm}, m_{\tilde{\chi}_2^\pm}$	406.25, 741.03	1038.00, 2491.42	304.84, 518.70
$m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}_2^0}$	356.62, 417.02	548.95, 1038.00	241.80, 316.70
$m_{\tilde{\chi}_3^0}, m_{\tilde{\chi}_4^0}$	424.40, 741.06	2489.56, 2491.14	322.35, 518.80
$m_A, m_{H^\pm}$	2573.69, 2573.69	1967.83, 1967.50	1846.04, 1846.04
m_h	124.42	126.55	121.57
$\Omega_{\tilde{\chi}_1} h^2$	0.1105	0.1002	0.1106
$BF(b \to s\gamma)$	3.23×10^{-4}	2.96×10^{-4}	3.15×10^{-4}
$BF(B_s \to \mu^+\mu^-)$	2.98×10^{-9}	5.23×10^{-9}	2.89×10^{-9}
$R_{(B \to \tau\nu)}$	0.98	0.98	0.97
Δa_μ	5.78×10^{-11}	1.65×10^{-10}	1.22×10^{-10}
$\sigma_{p\chi}^{\mathrm{SI}}$	3.05×10^{-8}	1.04×10^{-11}	2.64×10^{-8}

that it is within an acceptable zone considering the various uncertainties to compute m_h as mentioned before. Points A and C have larger spin-independent scattering cross section $\sigma_{p\chi}^{\mathrm{SI}}$ than the XENON100 limit. But, we believe this is within an acceptable limit considering the existing uncertainties arising out of strangeness content of nucleon, as well as those from astrophysical origins particularly from local DM density. Finally, it is also possible to satisfy all the limits in full subject to a 5%–10% heavier spectra and/or considering a multicomponent DM scenario.

IV. CONCLUSION

In the SSM [8], one assumes that a manifestation of an unknown but a fundamental mechanism of SUSY breaking may effectively lead to stochasticity in the Grassmannian parameters of the superspace. With a suitable probability distribution decided out of physical requirements, stochasticity in Grassmannian coordinates for a given Kähler potential and a superpotential may lead to well-known soft breaking terms. When applied to the superpotential of the MSSM, the model leads to soft breaking terms like

the bilinear Higgs coupling term, a trilinear soft term, as well as a gaugino mass term, all related to a parameter ξ, the scale of SUSY breaking. The other scale considered in the model of Ref. [8] at which the input quantities are given is Λ, where the latter can assume a value between the gauge coupling unification scale M_G and the Planck mass scale M_P. There is an absence of a scalar mass soft term at Λ in the original model that leads to stau turning to be the LSP or even turning itself tachyonic if Λ is chosen as M_G. This is only partially ameliorated when Λ is above M_G. However, the model because of its nonvanishing trilinear parameter is potentially accommodative to have a larger lighter Higgs boson mass via large stop scalar mixing. As recently shown in Ref. [10], the model in spite of its nice feature of bringing out the desired soft breaking terms of MSSM is not able to produce m_h above 116 GeV, and it has a LSP which is only a subdominant component of DM.

In this work, a minimal modification (referred to as Mod-SSM) is made by allowing a nonvanishing single scalar mass parameter m_0 as an explicit soft breaking term. Phenomenologically, the above addition is similar to what was considered in minimal AMSB while confronting the

issue of sleptons turning tachyonic in a pure AMSB framework. The modified model successfully accommodates the lighter Higgs boson mass near 125 GeV. Additionally, it can accommodate the stringent constraints from the dark matter relic density, $Br(B_s \rightarrow \mu^+ \mu^-)$, $Br(b \rightarrow s\gamma)$, and XENON100 data on direct detection of dark matter. A variation of top-quark mass within its allowed range is included in the analysis, and this shows the LSP to be a suitable candidate for dark matter satisfying both the limits of WMAP-7 data. Finally, we remind that the idea of stochastic superspace can easily be generalized to various scenarios beyond the MSSM.

ACKNOWLEDGMENTS

APPENDIX

Using Ref. [8], we consider the following Hermitian probability distribution:

$$
\begin{aligned}
\mathcal{P}(\theta, \bar{\theta}) &= A + \theta^\alpha \Psi_\alpha + \bar{\theta}_{\dot{\alpha}} \bar{\Xi}^{\dot{\alpha}} + \theta^\alpha \theta_\alpha B + \bar{\theta}_{\dot{\alpha}} \bar{\theta}^{\dot{\alpha}} C \\
&\quad + \theta^\alpha \sigma^\mu_{\ \alpha\dot{\beta}} \bar{\theta}^{\dot{\beta}} V_\mu + \theta^\alpha \theta_\alpha \bar{\theta}_{\dot{\alpha}} \bar{\Lambda}^{\dot{\alpha}} + \bar{\theta}_{\dot{\alpha}} \bar{\theta}^{\dot{\alpha}} \theta^\alpha \Sigma_\alpha \\
&\quad + \theta^\alpha \theta_\alpha \bar{\theta}_{\dot{\alpha}} \bar{\theta}^{\dot{\alpha}} D.
\end{aligned}
\tag{A1}
$$

Here, A, B, C, D, and V_μ are complex numbers. Ψ, $\bar{\Xi}$, $\bar{\Lambda}$, and Σ are Grassmann numbers.

In order to arrive at the results of Ref. [8], we use the following [3]: $d^2\theta = -\frac{1}{4} d\theta^\alpha d\theta_\alpha$, $d^2\bar{\theta} = -\frac{1}{4} d\bar{\theta}_{\dot{\alpha}} d\bar{\theta}^{\dot{\alpha}}$, $d^4\theta = d^2\theta d^2\bar{\theta}$, $\int d^2\theta(\theta\theta) = 1$, $\int d^2\bar{\theta}(\bar{\theta}\bar{\theta}) = 1$, $\int d^2\theta = \int d^2\bar{\theta} = 0$, and $\int d^2\theta \theta^\alpha = \int d^2\bar{\theta} \bar{\theta}_{\dot{\alpha}} = 0$.

Normalization: First, $\mathcal{P}(\theta, \bar{\theta})$ should satisfy the normalization condition $\int d^2\theta d^2\bar{\theta} \mathcal{P}(\theta, \bar{\theta}) = 1$. All the terms except the one with the coefficient D vanishes in $\int d^2\theta d^2\bar{\theta} \mathcal{P}(\theta, \bar{\theta})$. Thus, $D = 1$.

Vanishing fermionic moments: Next, we require vanishing of moments of fermionic type because of the requirement of the Lorentz invariance. The fact that $\langle\theta^\beta\rangle = \int d^2\theta d^2\bar{\theta} \theta^\beta \mathcal{P}(\theta, \bar{\theta}) = 0$ means $\Sigma_\alpha = 0$. Similarly, $\langle\bar{\theta}_{\dot{\beta}}\rangle = 0$ means $\bar{\Lambda}^{\dot{\alpha}} = 0$, and $\langle\theta^\beta \bar{\theta}_{\dot{\gamma}}\rangle = 0$ leads to $V_\mu = 0$. Finally, $\langle\theta^2 \bar{\theta}_{\dot{\beta}}\rangle = 0$ gives $\bar{\Xi}^{\dot{\alpha}} = 0$, and $\langle\theta^\beta \bar{\theta}^2\rangle = 0$ gives $\Psi_\alpha = 0$.

Bosonic moments: We now compute the bosonic moments.

$$
\begin{aligned}
\langle\theta\theta\rangle &= \int d^2\theta d^2\bar{\theta} \theta^\beta \theta_\beta \mathcal{P}(\theta, \bar{\theta}) \\
&= \int d^2\theta d^2\bar{\theta}(\theta\theta)(\bar{\theta}\bar{\theta})C = C.
\end{aligned}
$$

Similarly, $\langle\bar{\theta}\bar{\theta}\rangle = B$. Calling $B = 1/\xi$, one has $B = C^* = 1/\xi$. Furthermore, $\langle\theta\theta\bar{\theta}\bar{\theta}\rangle = A$. The fact that $\langle\theta\theta\bar{\theta}\bar{\theta}\rangle = \langle\theta\theta\rangle\langle\bar{\theta}\bar{\theta}\rangle$ leads to $A = 1/|\xi|^2$.

Thus we find the following Hermitian probability measure for the stochastic Grassmann variables:

$$
\begin{aligned}
\mathcal{P}(\theta, \bar{\theta})|\xi|^2 &= \tilde{\mathcal{P}}(\theta, \bar{\theta}) \\
&= 1 + \xi^*(\theta\theta) + \xi(\bar{\theta}\bar{\theta}) + |\xi|^2(\theta\theta)(\bar{\theta}\bar{\theta}).
\end{aligned}
\tag{A2}
$$

We consider the Wess-Zumino model with a single chiral superfield Φ. Φ has the following expansion [3]:

$$
\begin{aligned}
\Phi &= \phi(x) - i\theta\sigma^\mu\bar{\theta}\partial_\mu\phi(x) - \frac{1}{4}\theta^2\bar{\theta}^2\partial_\mu\partial^\mu\phi(x) \\
&\quad + \sqrt{2}\theta\psi(x) + \frac{i}{\sqrt{2}}\theta^2\partial_\mu\psi(x)\sigma^\mu\bar{\theta} + \theta^2 F(x).
\end{aligned}
\tag{A3}
$$

Correspondingly, for $\Phi^\dagger$ we have

$$
\begin{aligned}
\Phi^\dagger &= \phi^*(x) + i\theta\sigma^\mu\bar{\theta}\partial_\mu\phi^*(x) - \frac{1}{4}\theta^2\bar{\theta}^2\partial_\mu\partial^\mu\phi^*(x) \\
&\quad + \sqrt{2}\bar{\theta}\bar{\psi}(x) - \frac{i}{\sqrt{2}}\bar{\theta}^2\theta\sigma^\mu\partial_\mu\bar{\psi}(x) + \bar{\theta}^2 F^*(x).
\end{aligned}
\tag{A4}
$$

The kinetic term of the Lagrangian $\mathcal{L}$ is $[\Phi^\dagger\Phi]_D$. One finds

$$
\begin{aligned}
\Phi^\dagger\Phi &= |\phi|^2 + \sqrt{2}\theta\psi\phi^* + \sqrt{2}\bar{\theta}\bar{\psi}\phi + \theta^2\phi^*F + \bar{\theta}^2 F^*\phi \\
&\quad + 2\bar{\theta}\bar{\psi}\theta\psi + i\sqrt{2}\theta^2\bar{\theta}\bar{\sigma}^\mu\psi[\partial_\mu]\phi^* + \sqrt{2}\theta^2\bar{\theta}\bar{\psi}F \\
&\quad - 2i\theta\sigma^\mu\bar{\theta}\phi^*[\partial_\mu]\phi + i\sqrt{2}\bar{\theta}^2\theta\sigma^\mu\bar{\psi}[\partial_\mu]\phi \\
&\quad + \sqrt{2}\bar{\theta}^2\theta\psi F^* + \theta^2\bar{\theta}^2\Big(F^*F + \frac{1}{2}\partial_\mu\phi^*[\partial^\mu]\phi \\
&\quad - \frac{1}{2}\phi^*[\partial_\mu]\partial^\mu\phi + i\psi\sigma^\mu[\partial_\mu]\bar{\psi}\Big).
\end{aligned}
\tag{A5}
$$

Here, $X[\partial^\mu]Y = \frac{1}{2}(X\partial_\mu Y - Y\partial_\mu X)$. Upon vanishing the appropriate surface terms, the D term in particular reads

$$
[\Phi^\dagger\Phi]_D = F^*F + \partial_\mu\phi^*\partial^\mu\phi + \frac{i}{2}(\psi\sigma^\mu\partial_\mu\bar{\psi} - \partial_\mu\psi\sigma^\mu\bar{\psi}).
\tag{A6}
$$

Next, we consider the superpotential W given by $W = \frac{1}{2}m\Phi^2 + \frac{1}{3}h\Phi^3$. One has

$$
\begin{aligned}
\Phi^2 &= \phi^2 + 2\sqrt{2}\theta\psi\phi + \theta^2(2\phi F - \psi\psi), \quad \text{and} \\
\Phi^3 &= \phi^3 + 3\sqrt{2}\theta\psi\phi^2 + 3\theta^2(F\phi^2 - \psi\psi\phi).
\end{aligned}
\tag{A7}
$$

Thus,

$$W = \left(\frac{1}{2}m\phi^2 + \frac{1}{3}h\phi^3\right) + \sqrt{2}\theta\psi(m\phi + h\phi^2)$$
$$+ \theta^2\left(m\phi F - \frac{1}{2}m\psi\psi + hF\phi^2 - h\psi\psi\phi\right). \quad (A8)$$

The potential energy density term will be as follows:

$$[W + \text{H.c.}]_F = \left(m\phi F - \frac{1}{2}m\psi\psi + hF\phi^2 - h\psi\psi\phi\right) + \text{H.c.}$$
$$(A9)$$

Kinetic and potential terms averaged over θ and $\bar{\theta}$ and emergence of soft SUSY breaking terms:

Averaging over the Grassmannian coordinates, we compute $L = \langle\mathcal{L}\rangle = \int d^2\theta d^2\bar{\theta} \, \tilde{\mathcal{P}}(\theta, \bar{\theta})\mathcal{L}$. Here $\mathcal{L}$ is the usual super-Lagrangian density: $\mathcal{L} = \Phi^\dagger\Phi + W\delta^{(2)}(\bar{\theta}) + W^\dagger\delta^{(2)}(\theta)$. Then, using Eqs. (A2) and (A5) $\langle\mathcal{L}_{\text{kinetic}}\rangle$, namely, the kinetic part of L is found as

$$\langle\mathcal{L}_{\text{kinetic}}\rangle = [\Phi^\dagger\Phi]_D + \overbrace{|\xi|^2|\phi|^2 + \xi^*\phi F^* + \xi\phi^*F}^{\cancel{S}\text{USY}}.$$
$$(A10)$$

Similarly, the potential energy density averaged over θ and $\bar{\theta}$ is given by

$$\langle W + \text{H.c.}\rangle = \int d^2\theta d^2\bar{\theta} \, \tilde{\mathcal{P}}(\theta, \bar{\theta})(W\delta^{(2)}(\bar{\theta}) + W^\dagger\delta^{(2)}(\theta)).$$
$$(A11)$$

Using Eqs. (A2) and (A8) we find

$$\langle W + \text{H.c.}\rangle = \left[\overbrace{\xi^*\left(\frac{1}{2}m\phi^2 + \frac{1}{3}h\phi^3\right)}^{\cancel{S}\text{USY}}\right.$$
$$\left. + \left(m\phi F - \frac{1}{2}m\psi\psi + hF\phi^2 - h\psi\psi\phi\right)\right]$$
$$+ \text{H.c.} \quad (A12)$$

The total Lagrangian is then

$$L = \langle\mathcal{L}\rangle = \langle\mathcal{L}_{\text{kinetic}}\rangle + \langle W + \text{H.c.}\rangle. \quad (A13)$$

Using Eqs. (A9), (A10), and (A12) we break L into SUSY invariant and SUSY breaking parts as follows:

$$L = \langle\mathcal{L}\rangle_{\text{SUSY}} + \langle\mathcal{L}\rangle_{\cancel{S}\text{USY}}, \quad (A14)$$

where

$$\langle\mathcal{L}\rangle_{\cancel{S}\text{USY}} = [\Phi^\dagger\Phi]_D + [W + \text{H.c.}]_F, \quad (A15)$$

and

$$\langle\mathcal{L}\rangle_{\cancel{S}\text{USY}} = |\xi|^2|\phi|^2 + \xi^*\phi F^* + \xi\phi^*F$$
$$+ \left[\xi^*\left(\frac{1}{2}m\phi^2 + \frac{1}{3}h\phi^3\right) + \text{H.c.}\right]. \quad (A16)$$

The equations of motion of auxiliary fields are then

$$F = -(\xi^*\phi + m\phi^* + h\phi^{*2}), \quad \text{and}$$
$$F^* = -(\xi\phi^* + m\phi + h\phi^2). \quad (A17)$$

Substituting F and F^* in L, one finds

$$L = L_{\text{On-shell-SUSY}} + L_{\text{soft}}, \quad (A18)$$

where $L_{\text{On-shell-SUSY}}$ is the usual on shell SUSY invariant Lagrangian for the interacting Wess-Zumino model and is given by

$$L_{\text{On-shell-SUSY}}$$
$$= \partial_\mu\phi^*\partial^\mu\phi + \frac{i}{2}(\psi\sigma^\mu\partial_\mu\bar{\psi} - \partial_\mu\psi\sigma^\mu\bar{\psi})$$
$$- m^2|\phi|^2 - h^2(|\phi|^2)^2$$
$$- \left[\left(mh|\phi|^2\phi + \frac{1}{2}m\psi\psi + h\psi\psi\phi\right) + \text{H.c.}\right], \quad (A19)$$

and L_{soft} is given by

$$-L_{\text{soft}} = \left[\left(\frac{1}{2}\xi^*m\phi^2 + \frac{2}{3}h\xi^*\phi^3\right) + \text{H.c.}\right]. \quad (A20)$$

We remind that a negative sign in the left-hand side of Eq. (A20) appears simply because of considering a positive sign before $\langle W + \text{H.c.}\rangle$ in Eq. (A13) while writing the total Lagrangian. We note that it is only the superpotential term in the original theory that leads to soft breaking terms in the resulting Lagrangian ($m \to \xi^*m$ and $h \to 2\xi^*h$ going from W to L_{soft}). The presence of a vector field will not lead to any soft SUSY breaking term. This can easily be seen by considering a vector superfield in Wess-Zumino gauge.

MSSM: In the MSSM, as mentioned in Ref. [8] the superpotential term $\mu H_u H_d$ will lead to $\xi^*\mu\tilde{H}_u\tilde{H}_d$, and terms like $\hat{y}^{up}QU^cH_u$ will lead to $2\xi^*\hat{y}^{up}\tilde{Q}\tilde{U}_c\tilde{H}_u$ as soft SUSY breaking terms. Here, fields with tildes denote the scalar component of the corresponding chiral superfields. One further obtains a gaugino mass term $\frac{\xi^*}{2}\Sigma_i\lambda^{(i)}\lambda^{(i)}$. Thus, one finds a universal gaugino mass $m_{1/2} = \frac{1}{2}|\xi|$, a bilinear Higgs soft parameter $B_\mu = \xi^*$, and a universal trilinear soft parameter $A_0 = 2\xi^*$. With no resulting scalar mass term, one has the universal scalar mass parameter $m_0 = 0$. These are the input quantities to be given at a scale Λ. Low energy spectra are then found via RG evolutions. Reference [8] considered ξ and Λ as the input quantities and considered $M_G < \Lambda < M_P$.

[1] For reviews on supersymmetry, see, e.g., H. P. Nilles, Phys. Rep. **110**, 1 (1984); J. D. Lykken, arXiv:hep-th/9612114; J. Wess and J. Bagger, *Supersymmetry and Supergravity* (Princeton University, Princeton, NJ, 1991), 2nd ed.

[2] D. J. H. Chung, L. L. Everett, G. L. Kane, S. F. King, J. D. Lykken, and L. T. Wang, Phys. Rep. **407**, 1 (2005); H. E. Haber and G. Kane, Phys. Rep. **117**, 75 (1985); S. P. Martin, arXiv:hep-ph/9709356.

[3] M. Drees, P. Roy, and R. M. Godbole, *Theory and Phenomenology of Sparticles* (World Scientific, Singapore, 2005).

[4] H. Baer and X. Tata, *Weak Scale Supersymmetry: From Superfields to Scattering Events* (Cambridge University Press, Cambridge, England, 2006), p. 537.

[5] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012); S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).

[6] G. Aad *et al.* (ATLAS Collaboration), Phys. Rev. D **86**, 092002 (2012); , Phys. Lett. B **710**, 67 (2012); S. Chatrchyan *et al.* (CMS Collaboration), Phys. Rev. Lett. **107**, 221804 (2011).

[7] 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. J. Hall, J. Lykken, and S. Weinberg, Phys. Rev. D **27**, 2359 (1983); P. Nath, R. Arnowitt, and A. H. Chamseddine, Nucl. Phys. **B227**, 121 (1983); N. Ohta, Prog. Theor. Phys. **70**, 542 (1983); P. Nath, R. Arnowitt, and A. H. Chamseddine, *Applied $N = 1$ Supergravity* (World Scientific, Singapore, 1984).

[8] A. Kobakhidze, N. Pesor, and R. R. Volkas, Phys. Rev. D **79**, 075022 (2009).

[9] A. Kobakhidze, N. Pesor, and R. R. Volkas, Phys. Rev. D **81**, 095019 (2010).

[10] A. Kobakhidze, N. Pesor, R. R. Volkas, and M. J. White, Phys. Rev. D **85**, 075023 (2012).

[11] R. Aaij *et al.* (LHCb Collaboration), Phys. Rev. Lett. **110**, 021801 (2013).

[12] E. Aprile *et al.* (XENON100 Collaboration), Phys. Rev. Lett. **109**, 181301 (2012).

[13] R. L. Arnowitt and P. Nath, Phys. Rev. D **56**, 2833 (1997).

[14] G. F. Giudice, M. A. Luty, H. Murayama, and R. Rattazzi, J. High Energy Phys. 12 (1998) 027; L. Randall and R. Sundrum, Nucl. Phys. **B557**, 79 (1999); J. Bagger, T. Moroi, and E. Poppitz, J. High Energy Phys. 04 (2000) 009.

[15] T. Gherghetta, G. F. Giudice, and J. D. Wells, Nucl. Phys. **B559**, 27 (1999); J. L. Feng, T. Moroi, L. Randall, M. Strassler, and S. Su, Phys. Rev. Lett. **83**, 1731 (1999); J. L. Feng and T. Moroi, Phys. Rev. D **61**, 095004 (2000); U. Chattopadhyay, D. K. Ghosh, and S. Roy, Phys. Rev. D **62**, 115001 (2000).

[16] R. Arnowitt and P. Nath, Phys. Rev. D **46**, 3981 (1992).

[17] G. Gamberini, G. Ridolfi, and F. Zwirner, Nucl. Phys. **B331**, 331 (1990); V. D. Barger, M. S. Berger, and P. Ohmann, Phys. Rev. D **49**, 4908 (1994); S. P. Martin, Phys. Rev. D **66**, 096001 (2002).

[18] A. Djouadi, J.-L. Kneur, and G. Moultaka, Comput. Phys. Commun. **176**, 426 (2007).

[19] M. Drees and M. M. Nojiri, Nucl. Phys. **B369**, 54 (1992).

[20] A. Djouadi, Phys. Rep. **459**, 1 (2008).

[21] H. E. Haber, R. Hempfling, and A. H. Hoang, Z. Phys. C **75**, 539 (1997).

[22] H. E. Haber and R. Hempfling, Phys. Rev. Lett. **66**, 1815 (1991); Y. Okada, M. Yamaguchi, and T. Yanagida, Prog. Theor. Phys. **85**, 1 (1991); Phys. Lett. B **262**, 54 (1991); J. Ellis, G. Ridolf, and F. Zwirner, Phys. Lett. B **257**, 83 (1991); **262**, 477 (1991).

[23] H. E. Haber and Y. Nir, Phys. Lett. B **306**, 327 (1993); H. E. Haber, arXiv:hep-ph/9505240; A. Dobado, M. J. Herrero, and S. Penaranda, Eur. Phys. J. C **17**, 487 (2000); J. F. Gunion and H. E. Haber, Phys. Rev. D **67**, 075019 (2003); A. Djouadi and R. M. Godbole, arXiv:0901.2030.

[24] A. Arbey, M. Battaglia, A. Djouadi, and F. Mahmoudi, J. High Energy Phys. 09 (2012) 107.

[25] S. Heinemeyer, O. Stal, and G. Weiglein, Phys. Lett. B **710**, 201 (2012).

[26] B. C. Allanach, A. Djouadi, J. L. Kneur, W. Porod, and P. Slavich, J. High Energy Phys. 09 (2004) 044.

[27] G. Degrassi, S. Heinemeyer, W. Hollik, P. Slavich, and G. Weiglein, Eur. Phys. J. C **28**, 133 (2003).

[28] R. V. Harlander, P. Kant, L. Mihaila, and M. Steinhauser, Phys. Rev. Lett. **100**, 191602 (2008); **101**, 039901 (2008); S. P. Martin, Phys. Rev. D **75**, 055005 (2007).

[29] S. Bertolini, F. Borzumati, and A. Masiero, Phys. Rev. Lett. **59**, 180 (1987); N. G. Deshpande, P. Lo, J. Trampetic, G. Eilam, and P. Singer, Phys. Rev. Lett. **59**, 183 (1987); B. Grinstein and M. B. Wise, Phys. Lett. B **201**, 274 (1988); B. Grinstein, R. P. Springer, and M. B. Wise, Phys. Lett. B **202**, 138 (1988); W.-S. Hou and R. S. Willey, Phys. Lett. B **202**, 591 (1988); B. Grinstein, R. P. Springer, and M. B. Wise, Nucl. Phys. **B339**, 269 (1990).

[30] S. Bertolini, F. Borzumati, A. Masiero, and G. Ridolfi, Nucl. Phys. **B353**, 591 (1991); R. Barbieri and G. F. Giudice, Phys. Lett. B **309**, 86 (1993); R. Garisto and J. N. Ng, Phys. Lett. B **315**, 372 (1993); P. Nath and R. L. Arnowitt, Phys. Lett. B **336**, 395 (1994); M. Ciuchini, G. Degrassi, P. Gambino, and G. F. Giudice, Nucl. Phys. **B534**, 3 (1998).

[31] D. Asner *et al.* (Heavy Flavor Averaging Group Collaboration), arXiv:1010.1589.

[32] S. R. Choudhury and N. Gaur, Phys. Lett. B **451**, 86 (1999); K. S. Babu and C. Kolda, Phys. Rev. Lett. **84**, 228 (2000); A. Dedes, H. K. Dreiner, and U. Nierste, Phys. Rev. Lett. **87**, 251804 (2001); P. H. Chankowski and L. Slawianowska, Phys. Rev. D **63**, 054012 (2001); R. Arnowitt, B. Dutta, T. Kamon, and M. Tanaka, Phys. Lett. B **538**, 121 (2002); J. K. Mizukoshi, X. Tata, and Y. Wang, Phys. Rev. D **66**, 115003 (2002); S. Baek, P. Ko, and W. Y. Song, J. High Energy Phys. 03 (2003) 054; G. L. Kane, C. Kolda, and J. E. Lennon, arXiv:hep-ph/0310042; T. Ibrahim and P. Nath, Phys. Rev. D **67**, 016005 (2003); J. R. Ellis, K. A. Olive, and V. C. Spanos, Phys. Lett. B **624**, 47 (2005); S. Akula, D. Feldman, P. Nath, and G. Peim, Phys. Rev. D **84**, 115011 (2011); A. J. Buras, J. Girrbach, D. Guadagnoli, and G. Isidori, arXiv:1208.0934.

[33] A. J. Buras, J. Girrbach, D. Guadagnoli, and G. Isidori, Eur. Phys. J. C **72**, 2172 (2012).

[34] L. Roszkowski, S. Trojanowski, K. Turzynski, and K. Jedamzik, arXiv:1212.5587.

[35] G. Isidori and P. Paradisi, Phys. Lett. B **639**, 499 (2006); B. Bhattacherjee, A. Dighe, D. Ghosh, and S. Raychaudhuri, Phys. Rev. D **83**, 094026 (2011).

[36] J. P. Lees *et al.* (*BABAR* Collaboration), arXiv:1207.0698.

[37] I. Adachi *et al.* (Belle Collaboration), arXiv:1208.4678.

[38] K. Hara *et al.* (Belle Collaboration), Phys. Rev. D **82**, 071101 (2010).

[39] W. Altmannshofer, M. Carena, N. Shah, and F. Yu, arXiv:1211.1976.

[40] S. Bodenstein, C. A. Dominguez, and K. Schilcher, Phys. Rev. D **85**, 014029 (2012); G.-C. Cho, K. Hagiwara, Y. Matsumoto, and D. Nomura, J. High Energy Phys. 11 (2011) 068; D. Ghosh, M. Guchait, S. Raychaudhuri, and D. Sengupta, Phys. Rev. D **86**, 055007 (2012); S. Akula, P. Nath, and G. Peim, Phys. Lett. B **717**, 188 (2012).

[41] G. Bertone, D. Hooper, and J. Silk, Phys. Rep. **405**, 279 (2005); G. Jungman, M. Kamionkowski, and K. Griest, Phys. Rep. **267**, 195 (1996).

[42] E. Komatsu *et al.* (WMAP Collaboration), Astrophys. J. Suppl. Ser. **192**, 18 (2011).

[43] S. Alekhin, A. Djouadi, and S. Moch, Phys. Lett. B **716**, 214 (2012).

[44] Tevatron Electroweak Working Group, CDF Collaboration, and D0 Collaboration, arXiv:1107.5255.

[45] F. Boudjema and G. D. La Rochelle, Phys. Rev. D **86**, 115007 (2012); S. Caron, J. Laamanen, I. Niessen, and A. Strubig, J. High Energy Phys. 06 (2012) 008; U. Chattopadhyay and D. P. Roy, Phys. Rev. D **68**, 033010 (2003).

[46] S. Mohanty, S. Rao, and D. P. Roy, J. High Energy Phys. 11 (2012) 175; M. A. Ajaib, T. Li, and Q. Shafi, Phys. Rev. D **85**, 055021 (2012); D. Feldman, Z. Liu, P. Nath, and B. D. Nelson, Phys. Rev. D **80**, 075001 (2009); U. Chattopadhyay, D. Das, D. K. Ghosh, and M. Maity, Phys. Rev. D **82**, 075013 (2010); D. Feldman, Z. Liu, and P. Nath, Phys. Rev. D **80**, 015007 (2009); U. Chattopadhyay, D. Das, and D. P. Roy, Phys. Rev. D **79**, 095013 (2009); U. Chattopadhyay and D. Das, Phys. Rev. D **79**, 035007 (2009); R. M. Godbole, M. Guchait, and D. P. Roy, Phys. Rev. D **79**, 095015 (2009);

U. Chattopadhyay, D. Das, P. Konar, and D. P. Roy, Phys. Rev. D **75**, 073014 (2007); U. Chattopadhyay, D. Das, A. Datta, and S. Poddar, Phys. Rev. D **76**, 055008 (2007); G. Belanger, F. Boudjema, A. Cottrant, R. M. Godbole, and A. Semenov, Phys. Lett. B **519**, 93 (2001).

[47] M. Drees and M. M. Nojiri, Phys. Rev. D **48**, 3483 (1993).

[48] N. Fornengo, S. Scopel, and A. Bottino, Phys. Rev. D **83**, 015001 (2011); A. Bottino, F. Donato, N. Fornengo, and S. Scopel, Phys. Rev. D **81**, 107302 (2010); A. Bottino, V. de Alfaro, N. Fornengo, S. Mignola, and S. Scopel, Astropart. Phys. **2**, 77 (1994); T. K. Gaisser, G. Steigman, and S. Tilav, Phys. Rev. D **34**, 2206 (1986).

[49] G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, Comput. Phys. Commun. **180**, 747 (2009).

[50] M. Perelstein and B. Shakya, arXiv:1208.0833; C. Beskidt, W. de Boer, D. I. Kazakov, and F. Ratnikov, J. High Energy Phys. 05 (2012) 094; J. Giedt, A. W. Thomas, and R. D. Young, Phys. Rev. Lett. **103**, 201802 (2009); H. Ohki, H. Fukaya, S. Hashimoto, T. Kaneko, H. Matsufuru, J. Noaki, T. Onogi, E. Shintani, and N. Yamada, Phys. Rev. D **78**, 054502 (2008); J. R. Ellis, K. A. Olive, and C. Savage, Phys. Rev. D **77**, 065026 (2008).

[51] See, for example, the second reference of Ref. [50].

[52] N. Arkani-Hamed, A. Delgado, and G. F. Giudice, Nucl. Phys. **B741**, 108 (2006).

[53] K. L. Chan, U. Chattopadhyay, and P. Nath, Phys. Rev. D **58**, 096004 (1998); U. Chattopadhyay, A. Corsetti, and P. Nath, Phys. Rev. D **68**, 035005 (2003); S. Akula, M. Liu, P. Nath, and G. Peim, Phys. Lett. B **709**, 192 (2012).

[54] J. L. Feng, K. T. Matchev, and T. Moroi, Phys. Rev. D **61**, 075005 (2000); Phys. Rev. Lett. **84**, 2322 (2000); J. L. Feng, K. T. Matchev, and F. Wilczek, Phys. Lett. B **482**, 388 (2000); J. L. Feng and F. Wilczek, Phys. Lett. B **631**, 170 (2005); U. Chattopadhyay, T. Ibrahim, and D. P. Roy, Phys. Rev. D **64**, 013004 (2001); U. Chattopadhyay, A. Datta, A. Datta, A. Datta, and D. P. Roy, Phys. Lett. B **493**, 127 (2000); S. P. Das, A. Datta, M. Guchait, M. Maity, and S. Mukherjee, Eur. Phys. J. C **54**, 645 (2008).