Supersymmetric particle and Higgs boson masses from the landscape: Dynamical versus spontaneous supersymmetry breaking

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Perturbative supersymmetry breaking on the landscape of string vacua is expected to favor large soft terms as a power-law or log distribution, but tempered by an anthropic veto of inappropriate vacua or vacua leading to too large a value for the derived weak scale, which is a violation of the atomic principle. Indeed, scans of such vacua yield a statistical prediction for light Higgs boson mass $m_h \sim 125$ GeV with sparticles (save possibly light Higgsinos) typically beyond LHC reach. In contrast, models of dynamical supersymmetry (SUSY) breaking (DSB)—with a hidden sector gauge coupling g^2 scanned uniformly —lead to gaugino condensation and a uniform distribution of soft parameters on a log scale. Then soft terms are expected to be distributed as $m_{\rm soft}^{-1}$ favoring small values. A scan of DSB soft terms generally leads to $m_h \ll 125$ GeV and sparticle masses usually below LHC limits. Thus, the DSB landscape scenario seems excluded from LHC search results. An alternative is that the exponential suppression of the weak scale is set anthropically on the landscape via the atomic principle.

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I. INTRODUCTION

One of the mysteries of nature is the origin of mass scales. At least in QCD, we have an answer: the hadronic mass scale can arise when the gauge coupling evolves to large values such that the fundamental constituents, the quarks, condense to bound states. From dimensional transmutation, the proton mass can be found even in terms of the Planck mass $m_{\rm Pl}$ via $m_{\rm proton} \simeq m_{\rm Pl}\exp(-8\pi^2/g^2)$, which gives the right answer for $g^2 \sim 1.8$.

Another mass scale begging for explanation is that associated with weak interactions: $m_{\rm weak} \simeq m_{W,Z,h} \sim 100$ GeV. In the Standard Model (SM), the Higgs mass is quadratically divergent so one expects m_h to blow up to the highest mass scale Λ for which the SM is the viable low energy effective field theory (EFT). Supersymmetrization of the SM eliminates the Higgs mass quadratic divergences so any remaining divergences are merely logarithmic [1,2]: the minimal supersymmetric Standard Model, or MSSM [3], can be viable up to the GUT or even Planck scales. In addition, the weak scale emerges as a derived consequence of the visible sector SUSY breaking scale $m_{\rm soft}$. So the concern for the magnitude of the weak scale is transferred to a concern for the origin of the soft breaking scale. In gravity mediated SUSY breaking models,[1] it is popular to impose spontaneous SUSY breaking (SSB) at tree level in the hidden sector, for instance via the SUSY breaking Polonyi superpotential [9]: $W = m_{\rm hidden}^2(\hat{h}+\beta)$, where $\hat{h}$ is the lone hidden sector field. For $\beta = (2-\sqrt{3})m_P$ (with m_P the reduced Planck mass $m_P \equiv m_{\rm Pl}/\sqrt{8\pi}$ and $m_{\rm hidden} \sim 10^{11}$ GeV) then one determines $m_{\rm soft} \sim m_{3/2} \sim m_{\rm weak}$. Thus, the exponentially suppressed hidden sector mass scale must be put in by hand, so SSB can apparently only accommodate, but not explain, the magnitude of the weak scale.[2]

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[1] In days of yore, gauge-mediated SUSY breaking models [4] were associated with dynamical SUSY breaking in that they allowed much lighter gravitinos. In gauge-mediated SUSY breaking models, the trilinear soft term A_0 is expected to be tiny, leading to too light a Higgs boson mass unless soft terms are in the 10–100 TeV regime [5–7]. Such large soft terms then lead to highly unnatural third generation scalars. For this reason, we focus on DSB in a gravity-mediation context [8].

[2] A related problem is how the SUSY conserving μ parameter is *also* generated at or around the weak scale. A recent explanation augments the MSSM by a Peccei-Quinn (PQ) sector plus a $\mathbb{Z}_{24}^R$ discrete R symmetry [10], which generates a gravity-safe accidental approximate $U(1)_{PQ}$, which solves the strong CP and SUSY μ problems, and leads to an axion decay constant $f_a \sim m_{\rm hidden}$ whilst $\mu \sim m_{\rm weak}$ [11]. A recent review of 20 solutions to the SUSY μ problem is given in Ref. [12].

A more attractive mechanism follows the wisdom of QCD and seeks to generate the SUSY breaking scale from dimensional transmutation, which automatically yields an exponential suppression. This is especially attractive in string models where the Planck scale is the only mass scale available. Then one could arrange for dynamical SUSY breaking (DSB) [1,13,14] (for reviews, see [15–17]) wherein SUSY breaking arises nonperturbatively.[3] Some possibilities include hidden sector gaugino condensation [19], where a hidden sector gauge group such as $SU(N)$ becomes confining at the scale Λ_{GC} and a gaugino condensate occurs with $\langle\lambda\lambda\rangle \sim \Lambda_{\mathrm{GC}}^3$ leading to SUSY breaking with soft terms $m_{\mathrm{soft}} \sim \Lambda_{\mathrm{GC}}^3/m_P^2$. The associated hidden mass scale [20] is given by

$$m_{\mathrm{hidden}}^2 \sim m_P^2 \exp(-8\pi^2/g_{\mathrm{hidden}}^2), \tag{1}$$

where then $m_{\mathrm{hidden}}^2 \sim \Lambda_{\mathrm{GC}}^3/m_P$. Another possibility is nonperturbative SUSY breaking via instanton effects, which similarly leads to an exponential suppression of mass scales [21]. Of course, now the mass scale selection problem has been transferred to the selection of an appropriate value of g_{hidden}^2.

A solution to the origin of mass scales also arises within the string landscape picture [22,23]. This picture makes use of the vast array of string vacua found in IIB flux compactifications [24]. Some common estimates from vacuum counting [25] are $N_{vac} \sim 10^{500}$–$10^{272{,}000}$ [26,27]. The landscape then provides a setting for Weinberg's anthropic solution to the cosmological constant problem [28]: the value of Λ_{cc} is expected to be as large as possible such that the expansion rate of the early Universe allows for galaxy condensation and hence the structure formation that seems essential for the emergence of life.

Can similar reasoning be applied to the origin of the weak scale, or better yet, the origin of the SUSY breaking scale? This issue has been explored initially in Refs. [29–31]. Here, one assumes a fertile patch of the landscape of vacua where the MSSM is the visible sector low energy EFT. The differential distribution of vacua is expected to be of the form

$$dN_{\mathrm{vac}}[m_{\mathrm{hidden}}^2, m_{\mathrm{weak}}, \Lambda_{\mathrm{cc}}] = f_{\mathrm{SUSY}} \cdot f_{\mathrm{EWSB}} \cdot f_{\mathrm{cc}} \cdot dm_{\mathrm{hidden}}^2, \tag{2}$$

where $f_{\mathrm{SUSY}}(m_{\mathrm{hidden}}^2)$ contains the distribution of SUSY breaking mass scales expected on the fertile patch and f_{EWSB} contains the anthropic weak scale selection criteria. Denef and Douglas have argued that the cosmological constant selection acts independently and hence does not affect landscape selection of the SUSY breaking scale [26].

For a fertile patch of SUSY models with SSB in the landscape, then no particular values for the SUSY breaking F_i and D_α terms are favored. This means that the F_i values are expected to be uniformly distributed across the landscape, but as complex numbers. But it is the magnitude of F_i, which dictates the magnitude of the SUSY breaking scale: $m_{\mathrm{hidden}}^2 \sim |F_i|$. Such a distribution will favor large SUSY breaking scales as m_{soft}^1. SUSY breaking via the real-valued D_α terms will also be uniformly distributed—but this time as real numbers (giving rise to a uniform distribution in m_{soft}). In most string models, multiple hidden sectors occur, and several SUSY breaking F_i and D_α fields may be present. In this more general case, then the overall SUSY breaking scale is distributed as a spherical shell in a multidimensional space of SUSY breaking fields. This would lead, in the case of spontaneous SUSY breaking, to a power law distribution of soft terms [29–31].

$$f_{\mathrm{SUSY}}^{\mathrm{SSB}} \sim m_{\mathrm{soft}}^n, \tag{3}$$

where $n = 2n_F + n_D - 1$ and n_F are the number of hidden sector SUSY breaking F fields and n_D is the number of hidden sector D-breaking fields contributing to the overall SUSY breaking scale. Such a distribution would tend to favor SUSY breaking at the highest possible mass scales for $n \geq 1$. Also, Broeckel *et al.* [32] analyzed the distributions of SUSY breaking scales from vacua for Kachru-Kallosh-Linde-Trivedi (KKLT) [33] and large volume scenario (LVS) [34] flux compactifications and found for the KKLT model that $f_{\mathrm{SUSY}} \sim m_{\mathrm{soft}}^2$ while the LVS model gives $f_{\mathrm{SUSY}} \sim \log(m_{\mathrm{soft}})$ [35].

For the anthropic selection, an initial guess was to take $f_{\mathrm{EWSB}} = (m_{\mathrm{weak}}/m_{\mathrm{soft}})^2$ corresponding to a simple fine-tuning factor that invokes a penalty for soft terms which stray too far beyond the measured value of the weak scale. As emphasized in Refs. [36,37], this breaks down in a number of circumstances: 1. soft terms leading to charge-or-color-breaking (CCB) vacua must be vetoed, not just penalized, 2. soft terms for which electroweak (EW) symmetry does not even break also ought to be vetoed (we label these as no electroweak symmetry breaking (EWSB) vacua), 3. for some soft terms, the larger they get, then the smaller the derived value of the weak scale becomes. To illustrate this latter point, we write the pocket universe (PU) [38] value of the weak scale in terms of the pocket-universe Z-boson mass m_Z^{PU} and use the MSSM Higgs potential minimization conditions to find

$$(m_Z^{\mathrm{PU}})^2/2 = \frac{m_{H_d}^2 + \Sigma_d^d - (m_{H_u}^2 + \Sigma_u^u)\tan^2\beta}{\tan^2\beta - 1} - \mu^2 \simeq -m_{H_u}^2 - \Sigma_u^u - \mu^2, \tag{4}$$

where $m_{H_{u,d}}^2$ are Higgs soft breaking masses, μ is the superpotential Higgsino mass arising from whatever

[3] The DSB scenario has been made more plausible in recent years with the advent of metastable DSB [17,18].

solution to the SUSY μ problem is invoked,[4] and $\tan\beta \equiv v_u/v_d$ is the ratio of Higgs field vacuum expectation values (vevs). The Σ_u^u and Σ_d^d contain over 40 1-loop radiative corrections, listed in the Appendix of Ref. [40]. The soft term $m_{H_u}^2$ must be driven to negative values at the weak scale in order to break EW symmetry. If its high scale value is small, then it is typically driven deep negative so that compensatory fine-tuning is needed in the μ term. If $m_{H_u}^2$ is too big, then it does not even run negative and EW symmetry is unbroken. The landscape draw to large soft terms pulls $m_{H_u}^2$ big enough so EW symmetry barely breaks, corresponding to a natural value of $m_{H_u}^2$ at the weak scale (this can be considered as a landscape selection mechanism for tuning the high scale value of $m_{H_u}^2$ to such large values that its weak scale value becomes natural.) Also, for large negative values of trilinear soft term A_t, then large cancellations occur in $\Sigma_u^u(\tilde{t}_{1,2})$ leading to more natural $\Sigma_u^u(\tilde{t}_{1,2})$ values and a large $m_h \sim 125$ GeV due to large stop mixing in its radiative corrections. Also, large values of first/second generation soft scalar masses $m_0(1,2)$ cause the stop mass soft term running to small values, thus also making the spectra more natural [41,42].

The correct anthropic condition we believe was set down by Agrawal, Barr, Donoghue and Seckel (ABDS) in Ref. [43] (see their Fig. 1). In that work, they show that for variable values of the weak scale, then nuclear physics is disrupted if the pocket-universe value of the weak scale $m_{\rm weak}^{\rm PU}$ deviates from our measured value $m_{\rm weak}^{\rm OU}$ by a factor 2–5 (here, OU refers to "our Universe"). For values of $m_{\rm weak}^{\rm PU}$ outside this range, then nuclei and hence atoms as we know them would not form. In order to be in accord with this atomic principle, then to be specific, we require

$$m_{\rm max}^{\rm PU} < 4m_Z^{\rm OU}, \tag{5}$$

where $m_{\rm max}^{\rm PU}$ is the square root of the absolute value of the maximal contribution to the right-hand side of Eq. (4). In the absence of fine-tuning of μ or $m_{H_u}^2$ in Eq. (4), this requirement is then the same as requiring the electroweak fine-tuning measure [40,44] $\Delta_{\rm EW} < 30$ or $f_{\rm EWSB} = \Theta(30 - \Delta_{\rm EW})$ as the anthropic condition while also vetoing CCB and no EWSB vacua.[5]

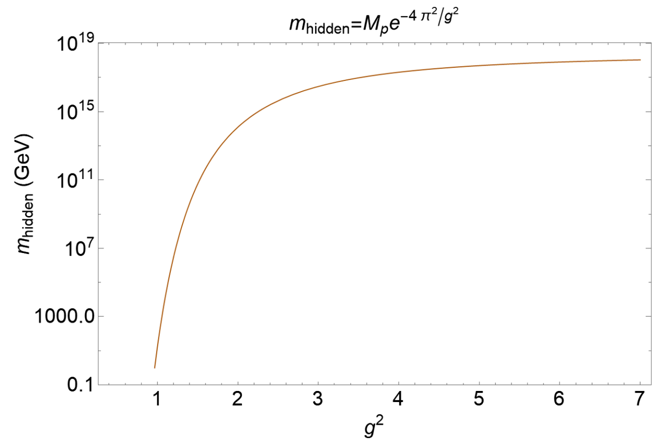

FIG. 1. Expected SUSY breaking scale $m_{\rm hidden}$ vs hidden sector coupling g^2 from dynamical SUSY breaking.

For the case of dynamical SUSY breaking, the SUSY breaking scale is expected to be of the form $m_{\rm hidden}^2 \sim m_P^2 \exp(-8\pi^2/g_{\rm hidden}^2)$ where in the case of gaugino condensation, $g_{\rm hidden}$ is the coupling constant of the confining hidden sector gauge group. It is emphasized by Dine *et al.* [45–47] and by Denef and Douglas [48] that the coupling $g_{\rm hidden}^2$ is expected to scan uniformly on the landscape. According to Fig. 1, for $g_{\rm hidden}^2$ values in the confining regime $\sim$1–2, we expect a uniform distribution of soft breaking terms on a log scale: i.e., each possible decade of values for $m_{\rm soft}$ is as likely as any other decade. Thus, with $m_{\rm soft} \sim m_{\rm hidden}^2/m_P \sim \Lambda_{\rm GC}^3/m_P^2$, we would expect

$$f_{\rm SUSY}^{\rm DSB} \sim 1/m_{\rm soft}, \tag{6}$$

which provides a uniform distribution of $m_{\rm soft}$ across the decades of possible values.[6] This simple model for $f_{\rm SUSY}^{\rm DSB}$ reflects the fact that no single value of $m_{\rm soft}$ is favored over any other across the multiple decades of possibilities.[7] For any interval of $m_{\rm soft}$ ranging from e^n to e^{n+1} GeV, the relative probability to lie within that interval is $P = \int_{e^n}^{e^{n+1}} f_{\rm SUSY}^{\rm DSB} dm_{\rm soft} = 1$ independent of n. Unlike the case of $f_{\rm SUSY}^{\rm SSB}$, the distribution $f_{\rm SUSY}^{\rm DSB}$ of course favors the lower range of soft term values.

II. RESULTS

Next, we will present the results of calculations of the string landscape probability distributions for Higgs and

[4]For the $\mathbb{Z}_{24}^R$ model for the origin of the μ parameter [11], the ensuing landscape distributions $dP/d\mu$ are plotted out in Ref. [39].

[5]Note that if large $\mu \gg 4m_{\rm weak}^{\rm OU}$ is selected, then almost never will the value of $m_{H_u}^2$ be selected such that $m_Z^{\rm PU} \sim 91.2$ GeV. This reflects the fact that fine-tuned vacua are relatively scarce on the landscape compared to non-fine-tuned vacua. Likewise, if $\mu \sim m_{\rm weak}^{\rm OU}$ is selected, then there is a substantial range of $m_{H_u}^2$ values leading to $m_Z^{\rm PU} \sim m_Z^{\rm OU}$. This is where stringy naturalness and conventional naturalness coincide (see, e.g., Fig. 3 of Ref. [41]). Notice that the value of the μ term need not scale as $f_{\rm SUSY}$ but instead depends on the particular mechanism for generating μ that is invoked. See Ref. [12] for a recent review of 20 solutions to the SUSY μ problem.

[6]Dine [45–47,49] actually finds $f_{\rm SUSY} \sim 1/[m_{\rm soft}\log(m_{\rm soft})]$, which is also highly uniform across the decades. We have checked that Dine's distribution gives even softer mass distributions than the $1/m_{\rm soft}$ which we use.

[7]Local structure in the $f_{\rm SUSY}^{\rm DSB}$ distribution might be expected since the confinement scale depends on discrete numbers like the number of colors and flavors in the confining gauge group. Thus, our distribution should reflect the overall global picture to be expected from DSB but not necessarily local substructure arising from details of the confining gauge group.

sparticle masses under the assumption of $f_{\rm SUSY}^{\rm DSB} = 1/m_{\rm soft}$ along with Eq. (5) for $f_{\rm EWSB}$. Our results will be presented within the gravity-mediated three extra parameter nonuniversal Higgs model NUHM3 with parameter space given by [50–55]

$$m_0(1,2), \quad m_0(3), \quad m_{1/2}, \quad A_0, \quad \tan\beta, \quad \mu, \quad m_A \quad \text{(NUHM3)}. \tag{7}$$

We adopt the ISAJET [56] code for calculation of the Higgs and superparticle mass spectrum [57] based on two-loop renormalization group equation (RGE) running [58] along with sparticle and Higgs masses calculated at the RG-improved one-loop level [59].

To compare our results against similar calculations which were presented in Ref. [37]—but using $f_{\rm SUSY} = m_{\rm soft}^n$—we will scan over the same parameter space

(i) $m_0(1,2)$: 0.1–60 TeV,
(ii) $m_0(3)$: 0.1–20 TeV,
(iii) $m_{1/2}$: 0.5–10 TeV,
(iv) A_0: -50–0 TeV,
(v) m_A: 0.3–10 TeV,

using the $f_{\rm SUSY}^{\rm DSB}$ distribution for soft terms[8] with $\mu = 150$ GeV. Our results hardly differ for other choices of μ: 100–360 GeV while any value of $\mu \gtrsim 360$ GeV will be vetoed due to violation of the ABDS anthropic condition of Eq. (5). A first guess for the value of $m_{H_u}^2$(weak) is gained from Eq. (4) with $m_Z^{\rm PU} \to m_Z^{\rm OU}$ and then the value of $m_{\rm weak}^{\rm PU} \simeq m_{\rm max}^{\rm PU}$ that enters Eq. (5) is determined. Then the condition Eq. (5) may be imposed as a veto on vacua with too large a value of $m_{\rm weak}^{\rm PU}$.

The parameter $\tan\beta$ is of course not a soft term so we have allowed it to scan uniformly over the range 3–60 in that not any one value of $\tan\beta$ is likely favored over any other within the phenomenologically allowed range given in the text. The ensuing distribution $dP/d\tan\beta$ is affected by the anthropic condition and is plotted out in Ref. [37] for the case of an $n = 0, 1, 2$ power-law draw of soft terms. The goal here was to choose upper limits to our scan parameters that will lie beyond the upper limits imposed by the anthropic selection from $f_{\rm EWFT}$. Lower limits are motivated by current LHC search limits, but also must stay away from the singularity in the $f_{\rm SUSY}^{\rm DSB}$ distribution. Our final results will hardly depend on the chosen value of μ so long as μ is within an factor of a few of $m_{W,Z,h} \sim 100$ GeV. We expect the different classes of soft terms to scan independently as discussed in Ref. [60]. We will compare the $f_{\rm SUSY}^{\rm DSB}$ results against the $f_{\rm SUSY}^{\rm SSB}$ results from Ref. [37] using an $n = 2$ power-law draw.

In Fig. 2, we first show probability distributions for various soft SUSY breaking terms for $f_{\rm SUSY}^{\rm DSB} = 1/m_{\rm soft}$ and also for $f_{\rm SUSY}^{\rm SSB} = m_{\rm soft}^2$. In Fig. 2(a), we show the distributions versus first/second generation soft breaking scalar masses $m_0(1,2)$. We see that the old SSB $n = 2$ result gives a peak distribution at $m_0(1,2) \sim 25$ TeV with a tail extending to over 40 TeV. This distribution reflects the mixed decoupling/quasidegeneracy landscape solution to the SUSY flavor and CP problems [42]. In contrast, the distribution from $f_{\rm SUSY}^{\rm DSB}$ peaks at the lowest allowed $m_0(1,2)$ values albeit with a tail extending out beyond 10 TeV. Thus, we would expect relatively light, LHC accessible, squarks and sleptons from gravity mediation with DSB in a hidden sector. In Fig. 2(b), we show the distribution in third generation soft mass inputs: $m_0(3)$. Here also the soft terms peak at the lowest values, but this time the tail extends only to ~4 TeV [lest $\Sigma_u^u(\tilde t_{1,2})$ becomes too large]. In contrast, the SSB $n = 2$ distribution peaks around 7 TeV. In Fig. 2(c), the distribution in unified gaugino soft term $m_{1/2}$ is shown. Here again, gaugino masses peak at the lowest allowed scales for DSB while the $n = 2$ distribution peaks just below 2 TeV. Finally, in Fig. 2(d), we show the distribution in trilinear soft term $-A_0$. Here, the DSB distribution peaks at $-A_0 \sim 0$, leading to little mixing in the stop sector and consequently lower values of m_h [61,62]. In contrast, the $n = 2$ distribution has a double peak structure with peaks at ~ -4 and -7 TeV with a tail extending to ~ -15 TeV: thus, we expect large stop mixing and higher m_h values from the SSB with $n = 2$ case.

In Fig. 3, we show distributions in light and heavy Higgs boson masses. In Fig. 3(a), we show the m_h distribution. For the DSB case, we see a peak at $m_h \sim 118$ GeV with almost no probability extending to ~125 GeV. This is in obvious contrast to the data and to the $n = 2$ distribution, which we see has a sharp peak at $m_h \sim 125$–126 GeV (as a result of large trilinear soft terms). In Fig. 3(b), we see the distribution in pseudoscalar Higgs mass m_A. In the DSB case, dP/dm_A peaks in the ~300 GeV range, leading to significant mixing in the Higgs sector and consequently possibly observable deviations in the Higgs couplings (see Ref. [63]). Alternatively, the SSB $n = 2$ distribution peaks at $m_A \sim 3.5$ TeV with a tail extending to ~8 TeV. In the latter case, we would expect a decoupled Higgs sector with a very SM-like lightest Higgs scalar h (as indeed the ATLAS/CMS data seem to suggest).

In Fig. 4, we show predictions for various sparticle masses from the DSB and SSB $n = 2$ cases. In Fig. 4(a), we show the distribution in gluino mass $m_{\tilde g}$. For the DSB case, the distribution peaks around the ~TeV range while LHC search limits typically require $m_{\tilde g} \gtrsim 2.2$ TeV. In fact, almost all parameter space of DSB is then excluded. Had we lowered the lower scan cutoff on $m_{1/2}$, the

[8]For the range of m_A values well above the measured Higgs mass (i.e., only small light-heavy Higgs mixing as is indicated by LHC Higgs coupling measurements), then m_A is very nearly m_{H_d} and should scan as m_{H_d} does. We used the NUHM3 model, which swaps m_A for m_{H_d} merely for convenience.

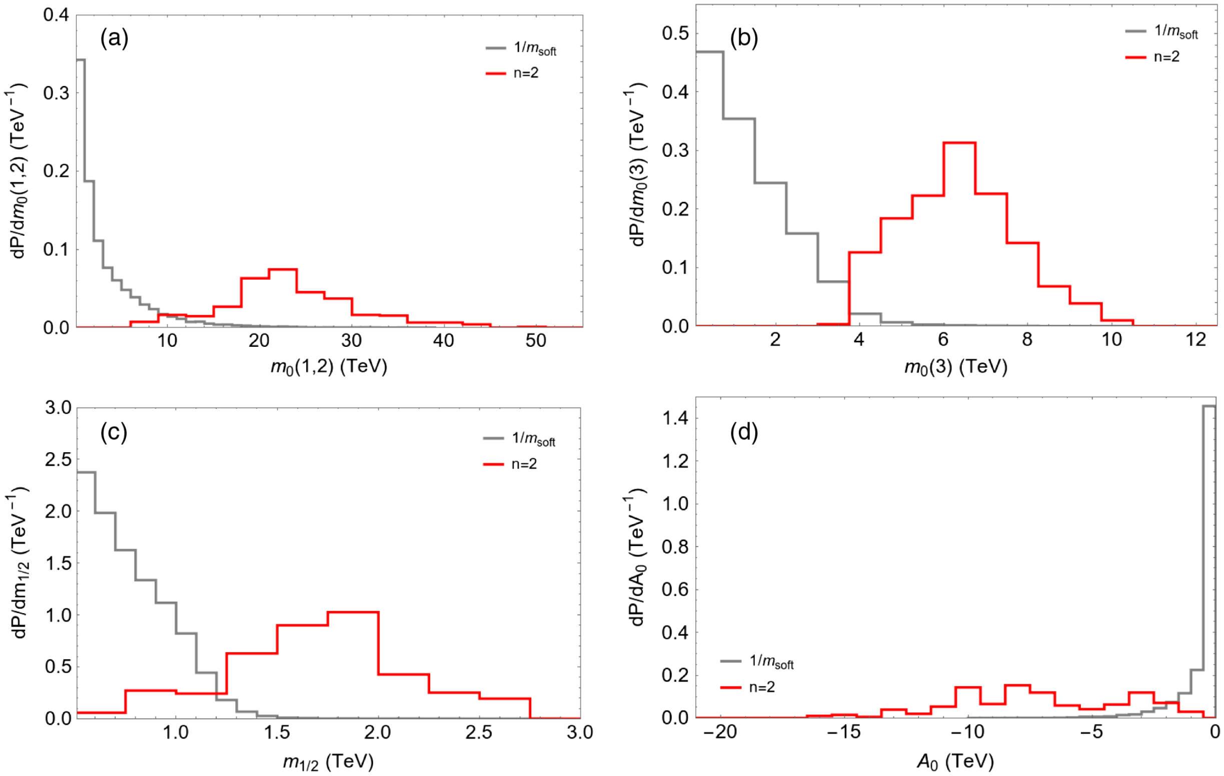

FIG. 2. Probability distributions for NUHM3 soft terms (a) $m_0(1,2)$, (b) $m_0(3)$, (c) $m_{1/2}$, and (d) A_0 from a $f_{\rm SUSY}^{\rm DSB} = 1/m_{\rm soft}$ distribution of soft terms in the string landscape with $\mu = 150$ GeV. For comparison, we also show probability distributions for $f_{\rm SUSY}^{\rm SSB} \sim m_{\rm soft}^2$.

distribution would shift lower, making matters worse. The SSB $n = 2$ distribution peaks at $m_{\tilde{g}} \sim 4$–5 TeV with a tail extending to ∼6 TeV; hardly any probability is excluded by the LHC $m_{\tilde{g}} \gtrsim 2.2$ TeV limit. In Fig. 4(b), we show the distribution in first generation squark mass $m_{\tilde{u}_L}$ (as a typical example of first/second generation matter scalars). The distribution from DSB peaks in the 0–3 TeV range with a tail extending beyond 10 TeV. Coupled with the gluino distribution, most probability space would be excluded by LHC search limits from the $m_{\tilde{g}}$ vs $m_{\tilde{q}}$ plane. The SSB $n = 2$ distribution peaks above 20 TeV with a tail extending beyond 40 TeV. In Fig. 4(c), we show the distribution in lighter top squark mass $m_{\tilde{t}_1}$. Here, we see DSB peaks around 1 TeV with a tail to ∼2.5 TeV. LHC searches require $m_{\tilde{t}_1} \gtrsim 1.1$ TeV so that about half of probability space is excluded. For the SSB $n = 2$ case, the peak shifts to $m_{\tilde{t}_1} \sim 1.6$ TeV so the bulk of p space is allowed by LHC searches. Finally, in Fig. 4(d), we show the distribution in heavier stop mass $m_{\tilde{t}_2}$. The DSB distribution peaks around ∼1.5 TeV whilst the SSB $n = 2$ distribution peaks around

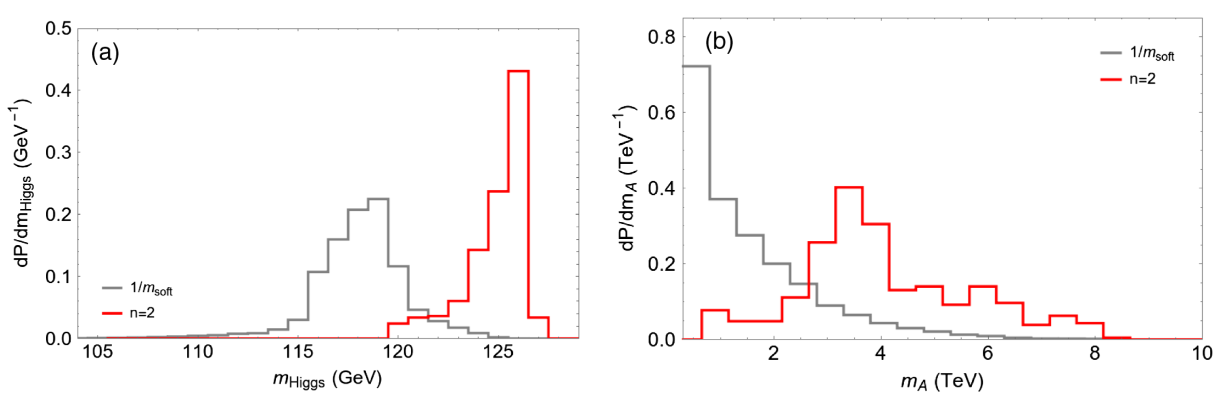

FIG. 3. Probability distributions for light Higgs scalar mass (a) m_h and pseudoscalar Higgs mass (b) m_A from a $f_{\rm SUSY}^{\rm DSB} \sim 1/m_{\rm soft}$ distribution of soft terms in the string landscape with $\mu = 150$ GeV. For comparison, we also show probability distributions for $f_{\rm SUSY}^{\rm SSB} \sim m_{\rm soft}^2$.

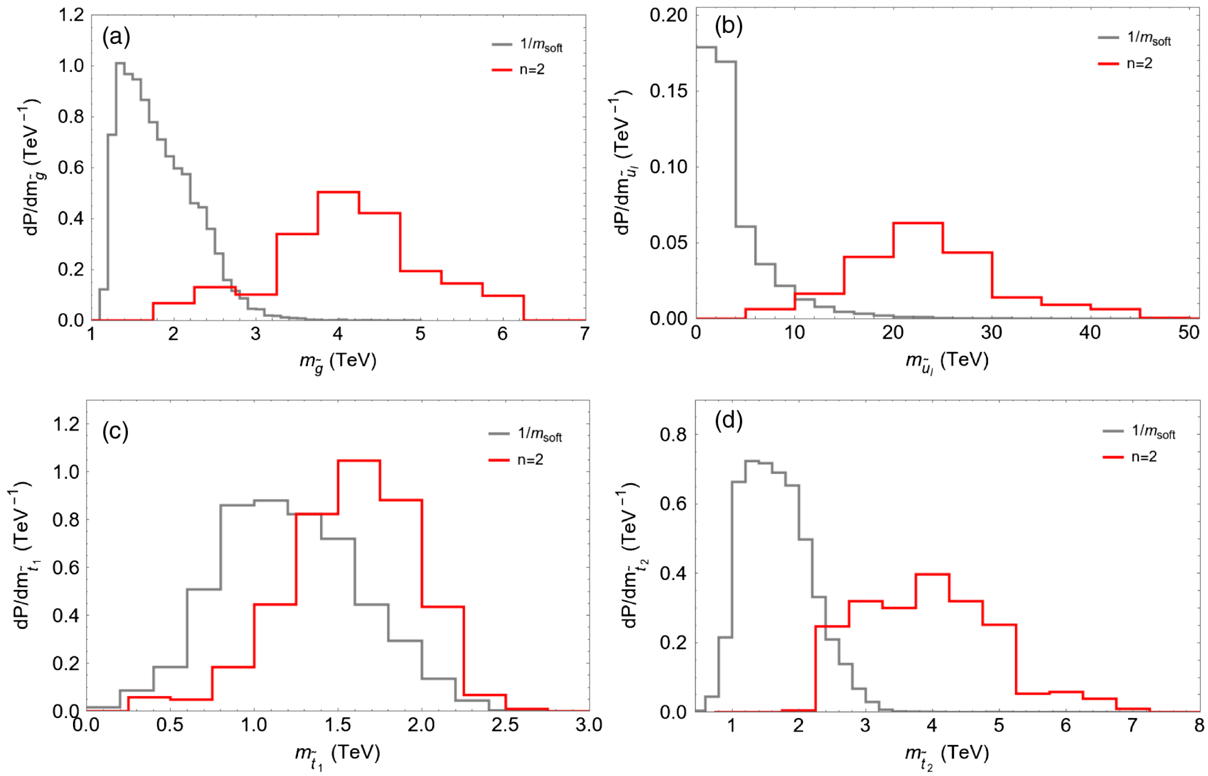

FIG. 4. Probability distributions for (a) $m_{\tilde{g}}$, (b) $m_{\tilde{u}_L}$, (c) $m_{\tilde{t}_1}$, and (d) $m_{\tilde{t}_2}$ from a $f_{\rm SUSY}^{\rm DSB} \sim 1/m_{\rm soft}$ distribution of soft terms in the string landscape with $\mu = 150$ GeV. For comparison, we also show probability distributions for $f_{\rm SUSY}^{\rm SSB} \sim m_{\rm soft}^2$.

4 TeV. Thus, substantially heavier $\tilde{t}_2$ squarks are expected from SSB as compared to DSB.

III. CONCLUSIONS

One of the mysteries of particle physics is the origin of mass scales, especially in the context of string theory where only the Planck scale m_P appears. Here, we investigated the origin of the weak scale which is presumed to arise from the scale of SUSY breaking. The general framework of dynamical SUSY breaking presents a beautiful example of the exponentially suppressed SUSY breaking scale (relative to the Planck scale) arising from nonperturbative effects such as gaugino condensation or SUSY breaking via instanton effects. The SUSY breaking scale from DSB is expected to be uniformly distributed on a log scale within a fertile patch of the string landscape with the MSSM as the low energy EFT. In this case, the probability distribution $f_{\rm SUSY}^{\rm DSB} \sim 1/m_{\rm soft}$. Such a distribution, coupled with the ABDS anthropic window, typically leads to Higgs masses m_h well below the measured 125 GeV value and many sparticles such as the gluino expected to lie below existing LHC search limits. Thus, the LHC data seem to falsify this approach. That would leave the alternative option of spontaneous SUSY breaking, where instead the soft SUSY breaking distribution is expected to occur as a power law or log distribution. These latter cases lead to landscape probability distributions for m_h that peak at $m_h \sim 125$ GeV with sparticles typically well beyond current LHC reach, but within reach of hadron colliders with $\sqrt{s} \gtrsim 30$ TeV [64]. For perturbative, or spontaneous, SUSY breaking, then apparently the magnitude of the SUSY breaking scale is set anthropically much like the cosmological constant is: those vacua with too large a SUSY breaking scale lead to either CCB or no EWSB vacua, or vacua with such a large weak scale that it lies outside the ABDS allowed window, in violation of the atomic principle.

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[1] E. Witten, Dynamical breaking of supersymmetry, Nucl. Phys. **B188**, 513 (1981).

[2] R. K. Kaul, Gauge hierarchy in a supersymmetric model, Phys. Lett. B **109**, 19 (1982).

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

[4] M. Dine, A. E. Nelson, Y. Nir, and Y. Shirman, New tools for low-energy dynamical supersymmetry breaking, Phys. Rev. D **53**, 2658 (1996).

[5] A. Arbey, M. Battaglia, A. Djouadi, F. Mahmoudi, and J. Quevillon, Implications of a 125 GeV Higgs for supersymmetric models, Phys. Lett. B **708**, 162 (2012).

[6] P. Draper, P. Meade, M. Reece, and D. Shih, Implications of a 125 GeV Higgs for the MSSM and low-scale SUSY breaking, Phys. Rev. D **85**, 095007 (2012).

[7] H. Baer, V. Barger, and A. Mustafayev, Neutralino dark matter in mSUGRA/CMSSM with a 125 GeV light Higgs scalar, J. High Energy Phys. 05 (2012) 091.

[8] M. Bose and M. Dine, Gravity mediation retrofitted, J. High Energy Phys. 03 (2013) 057.

[9] J. Polonyi, Generalization of the massive scalar multiplet coupling to the supergravity, Report No. KFKI-77-93, Ecole Normale Superieure (1977).

[10] H. M. Lee, S. Raby, M. Ratz, G. G. Ross, R. Schieren, K. Schmidt-Hoberg, and P. K. S. Vaudrevange, Discrete R symmetries for the MSSM and its singlet extensions, Nucl. Phys. **B850**, 1 (2011).

[11] H. Baer, V. Barger, and D. Sengupta, Gravity safe, electroweak natural axionic solution to strong *CP* and susy mu problems, Phys. Lett. B **790**, 58 (2019).

[12] K. J. Bae, H. Baer, V. Barger, and D. Sengupta, Revisiting the susy mu problem and its solutions in the LHC era, Phys. Rev. D **99**, 115027 (2019).

[13] S. Dimopoulos and S. Raby, Supercolor, Nucl. Phys. **B192**, 353 (1981).

[14] M. Dine, W. Fischler, and M. Srednicki, Supersymmetric technicolor, Nucl. Phys. **B189**, 575 (1981).

[15] E. Poppitz and S. P. Trivedi, Dynamical supersymmetry breaking, Annu. Rev. Nucl. Part. Sci. **48**, 307 (1998).

[16] Y. Shadmi and Y. Shirman, Dynamical supersymmetry breaking, Rev. Mod. Phys. **72**, 25 (2000).

[17] M. Dine and J. D. Mason, Supersymmetry and its dynamical breaking, Rep. Prog. Phys. **74**, 056201 (2011).

[18] K. A. Intriligator, N. Seiberg, and D. Shih, Dynamical SUSY breaking in meta-stable vacua, J. High Energy Phys. 04 (2006) 021.

[19] S. Ferrara, L. Girardello, and H. P. Nilles, Breakdown of local supersymmetry through gauge fermion condensates, Phys. Lett. **125B**, 457 (1983).

[20] I. Affleck, M. Dine, and N. Seiberg, Exponential hierarchy from dynamical supersymmetry breaking, Phys. Lett. **140B**, 59 (1984).

[21] I. Affleck, M. Dine, and N. Seiberg, Supersymmetry Breaking by Instantons, Phys. Rev. Lett. **51**, 1026 (1983).

[22] R. Bousso and J. Polchinski, Quantization of four form fluxes and dynamical neutralization of the cosmological constant, J. High Energy Phys. 06 (2000) 006.

[23] L. Susskind, *The Anthropic Landscape of String Theory* (Cambridge University Press, Cambridge, 2003), pp. 247–266.

[24] M. R. Douglas and S. Kachru, Flux compactification, Rev. Mod. Phys. **79**, 733 (2007).

[25] S. K. Ashok and M. R. Douglas, Counting flux vacua, J. High Energy Phys. 01 (2004) 060.

[26] F. Denef and M. R. Douglas, Distributions of flux vacua, J. High Energy Phys. 05 (2004) 072.

[27] W. Taylor and Y.-N. Wang, The F-theory geometry with most flux vacua, J. High Energy Phys. 12 (2015) 164.

[28] S. Weinberg, Anthropic Bound on the Cosmological Constant, Phys. Rev. Lett. **59**, 2607 (1987).

[29] L. Susskind, Supersymmetry breaking in the anthropic landscape, in *From Fields to Strings: Circumnavigating Theoretical Physics: A Conference in Tribute to Ian Kogan* (World Scientific, Singapore, 2004), pp. 1745–1749.

[30] M. R. Douglas, Statistical analysis of the supersymmetry breaking scale, arXiv:hep-th/0405279.

[31] N. Arkani-Hamed, S. Dimopoulos, and S. Kachru, Predictive landscapes and new physics at a TeV, arXiv:hep-th/0501082.

[32] I. Broeckel, M. Cicoli, A. Maharana, K. Singh, and K. Sinha, Moduli stabilisation and the statistics of SUSY breaking in the landscape, J. High Energy Phys. 10 (2020) 015.

[33] S. Kachru, R. Kallosh, A. Linde, and S. P. Trivedi, de sitter vacua in string theory, Phys. Rev. D **68**, 046005 (2003).

[34] V. Balasubramanian, P. Berglund, J. P. Conlon, and F. Quevedo, Systematics of moduli stabilisation in calabi-yau flux compactifications, J. High Energy Phys. 03 (2005) 007.

[35] H. Baer, V. Barger, S. Salam, and D. Sengupta, Landscape Higgs boson and sparticle mass predictions from a logarithmic soft term distribution, Phys. Rev. D **103**, 035031 (2021).

[36] H. Baer, V. Barger, M. Savoy, and H. Serce, The Higgs mass and natural supersymmetric spectrum from the landscape, Phys. Lett. B **758**, 113 (2016).

[37] H. Baer, V. Barger, H. Serce, and K. Sinha, Higgs and superparticle mass predictions from the landscape, J. High Energy Phys. 03 (2018) 002.

[38] A. H. Guth, Eternal inflation, Ann. N.Y. Acad. Sci. **950**, 66 (2001).

[39] H. Baer, V. Barger, D. Sengupta, and R. W. Deal, Distribution of supersymmetry μ parameter and Peccei-Quinn scale fa from the landscape, Phys. Rev. D **104**, 015037 (2021).

[40] H. Baer, V. Barger, P. Huang, D. Mickelson, A. Mustafayev, and X. Tata, Radiative natural supersymmetry: Reconciling electroweak fine-tuning and the higgs boson mass, Phys. Rev. D **87**, 115028 (2013).

[41] H. Baer, V. Barger, and S. Salam, Naturalness versus stringy naturalness with implications for collider and dark matter searches, Phys. Rev. Research **1**, 023001 (2019).

[42] H. Baer, V. Barger, and D. Sengupta, Landscape solution to the susy flavor and *CP* problems, Phys. Rev. Research **1**, 033179 (2019).

[43] V. Agrawal, S. M. Barr, J. F. Donoghue, and D. Seckel, Viable range of the mass scale of the standard model, Phys. Rev. D **57**, 5480 (1998).

[44] H. Baer, V. Barger, P. Huang, A. Mustafayev, and X. Tata, Radiative natural supersymmetry with a 125 GeV Higgs boson, Phys. Rev. Lett. **109**, 161802 (2012).

[45] T. Banks, M. Dine, and E. Gorbatov, Is there a string theory landscape?, J. High Energy Phys. 08 (2004) 058.
[46] M. Dine, E. Gorbatov, and S. D. Thomas, Low energy supersymmetry from the landscape, J. High Energy Phys. 08 (2008) 098.
[47] M. Dine, Naturalness under stress, Annu. Rev. Nucl. Part. Sci. **65**, 43 (2015).
[48] F. Denef and M. R. Douglas, Distributions of nonsupersymmetric flux vacua, J. High Energy Phys. 03 (2005) 061.
[49] M. Dine, The Intermediate scale branch of the landscape, J. High Energy Phys. 01 (2006) 162.
[50] D. Matalliotakis and H. Nilles, Implications of non-universality of soft terms in supersymmetric grand unified theories, Nucl. Phys. **B435**, 115128 (1995).
[51] M. Olechowski and S. Pokorski, Electroweak symmetry breaking with non-universal scalar soft terms and large tan solutions, Phys. Lett. **344B**, 201210 (1995).
[52] P. Nath and R. Arnowitt, Non-universal soft susy breaking and dark matter, COSMO-97, 10.1142/9789814447263_0020 (1998).
[53] J. Ellis, K. Olive, and Y. Santoso, The mssm parameter space with non-universal Higgs masses, Phys. Lett. B **539**, 107 (2002).
[54] J. Ellis, T. Falk, K. A. Olive, and Y. Santoso, Exploration of the mssm with non-universal higgs masses, Nucl. Phys. **B652**, 259347 (2003).
[55] H. Baer, A. Mustafayev, S. Profumo, A. Belyaev, and X. Tata, Direct, indirect and collider detection of neutralino dark matter in SUSY models with non-universal Higgs masses, J. High Energy Phys. 07 (2005) 065.
[56] F. E. Paige, S. D. Protopopescu, H. Baer, and X. Tata, ISAJET 7.69: A Monte Carlo event generator for pp, $\bar{p}$p, and e^+e^- reactions, arXiv:hep-ph/0312045.
[57] H. Baer, C.-H. Chen, R. B. Munroe, F. E. Paige, and X. Tata, Multichannel search for minimal supergravity at $p\bar{p}$ and e^+e^- colliders, Phys. Rev. D **51**, 1046 (1995).
[58] S. P. Martin and M. T. Vaughn, Two loop renormalization group equations for soft supersymmetry breaking couplings, Phys. Rev. D **50**, 2282 (1994).Erratum, Phys. Rev. D **78**, 039903 (2008).
[59] D. M. Pierce, J. A. Bagger, K. T. Matchev, and R.-j. Zhang, Precision corrections in the minimal supersymmetric standard model, Nucl. Phys. **B491**, 3 (1997).
[60] H. Baer, V. Barger, S. Salam, and D. Sengupta, String landscape guide to soft SUSY breaking terms, Phys. Rev. D **102**, 075012 (2020).
[61] M. Carena and H. E. Haber, Higgs boson theory and phenomenology, Prog. Part. Nucl. Phys. **50**, 63 (2003).
[62] H. Baer, V. Barger, and A. Mustafayev, Implications of a 125 GeV Higgs scalar for LHC SUSY and neutralino dark matter searches, Phys. Rev. D **85**, 075010 (2012).
[63] K. J. Bae, H. Baer, N. Nagata, and H. Serce, Prospects for Higgs coupling measurements in SUSY with radiatively-driven naturalness, Phys. Rev. D **92**, 035006 (2015).
[64] H. Baer, V. Barger, J. S. Gainer, D. Sengupta, H. Serce, and X. Tata, LHC luminosity and energy upgrades confront natural supersymmetry models, Phys. Rev. D **98**, 075010 (2018).