Implications of naturalness for the heavy Higgs bosons of supersymmetry

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Recently, it has been argued that various measures of supersymmetric naturalness—electroweak, Higgs mass and EENZ/BG—when applied consistently, concur with one another and make very specific predictions for natural supersymmetric spectra. Highly natural spectra are characterized by light Higgsinos with mass not too far from m_h and well-mixed but TeV-scale third generation squarks. We apply the unified naturalness measure to the case of heavy Higgs bosons A, H and $H^\pm$. We find that their masses are bounded from above by naturalness depending on $\tan\beta$: e.g. for 10% fine-tuning and $\tan\beta \sim 10$, we expect $m_A \lesssim$ 2.5 TeV whilst for 3% fine-tuning and $\tan\beta$ as high as 50, then $m_A \lesssim 8$ TeV. Furthermore, the presence of light Higgsinos seriously alters the heavy Higgs boson branching ratios, thus diminishing prospects for usual searches into standard model final states, while new discovery possibilities arise due to the supersymmetric decay modes. The heavy supersymmetric decay modes tend to be $H, A, H^\pm \to W, Z$, or $h + E_T + $ soft tracks so that single heavy Higgs production is characterized by the presence of high p_T W, Z or h bosons plus missing E_T. These new heavy Higgs boson signatures seem to be challenging to extract from SM backgrounds.

I. INTRODUCTION

The recent discovery of a standard model (SM) like Higgs boson with mass $m_h = 125.5 \pm 0.5$ GeV [1,2] is in accord with predictions from supersymmetric models, like the minimal supersymmetric standard model (MSSM), which require $m_h \lesssim 135$ GeV [3]. Such a large value of m_h apparently requires TeV-scale top squarks which are highly mixed, i.e. a large trilinear soft supersymmetry (SUSY) breaking parameter A_t [4]. Coupling this result with recent SUSY search limits from LHC8 [5,6] (which require $m_{\tilde{g}} \gtrsim 1.3$ TeV for $m_{\tilde{g}} \ll m_{\tilde{q}}$ and $m_{\tilde{g}} \gtrsim 1.8$ TeV for $m_{\tilde{g}} \sim m_{\tilde{q}}$) imply, within the context of gravity-mediated SUSY breaking models (SUGRA), a soft breaking scale characterized by a gravitino mass $m_{3/2} \gtrsim 2$ TeV. Indeed, a rather large SUSY breaking scale in gravity mediation models had been long anticipated via a decoupling solution to the SUSY flavor, CP, proton decay and gravitino problems [7].

In contrast, simple considerations of SUSY naturalness anticipate a SUSY breaking scale around the weak scale typified by $m_Z \sim m_h \sim 100$ GeV. Thus, the Higgs mass and sparticle mass limits combine to sharpen the "little hierarchy" [8] typified by $m_h \ll m_{3/2}$. The growing little hierarchy has prompted several authors to question whether the MSSM is overly fine-tuned, and either flat wrong [9] or at least in need of additional features which sacrifice

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parsimony/minimality [10]. Before rushing to such drastic conclusions, it is prudent to ascertain if all SUSY spectra are fine-tuned or if some spectra are indeed natural.

A. Review of SUSY naturalness

To proceed further one must adopt at least one of several quantitative naturalness measures which are available. We label these as

(i) the electroweak measure Δ_{EW} [11–15],
(ii) the Higgs mass fine-tuning measure Δ_{HS} [16,17] and
(iii) the traditional EENZ/BG measure Δ_{BG} [18,19].
Indeed, recently it has been shown that, if applied properly, then all three measures agree with one another and predict a very specific SUSY spectra with just $\sim$10% fine-tuning [20]. If applied incorrectly—by not properly combining *dependent* quantities contributing to m_Z or m_h one with another—then overestimates of fine-tuning can occur in Δ_{HS} and Δ_{BG}, often by orders of magnitude [21].

1. Δ_{EW}

The electroweak measure Δ_{EW} requires that there be no large/unnatural cancellations in deriving the value of m_Z from the weak scale scalar potential:

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

(1)

where $m_{H_u}^2$ and $m_{H_d}^2$ are the *weak scale* soft SUSY breaking Higgs masses, μ is the *supersymmetric* Higgsino mass term

and Σ_u^u and Σ_d^d contain an assortment of loop corrections to the effective potential. The Δ_{EW} measure asks for the largest contribution on the right-hand side to be comparable to $m_Z^2/2$ so that no unnatural fine-tunings are required to generate $m_Z = 91.2$ GeV. The main requirement is then that $|\mu| \sim m_Z$ and also that $m_{H_u}^2$ is driven radiatively to small, and not large, negative values [11,13]. Also, the top squark contributions $\Sigma_u^u(\tilde{t}_{1,2})$ are minimized for TeV-scale highly mixed top squarks, which also lift the Higgs mass to $m_h \sim 125$ GeV [11].

2. Δ_{HS}

The Higgs mass fine-tuning measure, Δ_{HS}, asks that the radiative correction $\delta m_{H_u}^2$ to the Higgs mass

$$m_h^2 \simeq \mu^2 + m_{H_u}^2(\Lambda) + \delta m_{H_u}^2 \qquad (2)$$

be comparable to m_h^2. This contribution is usually written as $\delta m_{H_u}^2|_{\mathrm{rad}} \sim -\frac{3f_t^2}{8\pi^2}(m_{Q_3}^2 + m_{U_3}^2 + A_t^2) \ln(\Lambda^2/m_{\mathrm{SUSY}}^2)$ which is used to claim that third generation squarks $m_{\tilde{t}_{1,2},\tilde{b}_1}$ be approximately less than 500 GeV and A_t be small for natural SUSY. However, several approximations are necessary to derive this result, the worst of which is to neglect that the value of $m_{H_u}^2$ itself contributes to $\delta m_{H_u}^2$. By combining dependent contributions, then instead one requires that the two terms on the right-hand side of

$$m_h^2 = \mu^2 + (m_{H_u}^2(\Lambda) + \delta m_{H_u}^2) \qquad (3)$$

be comparable to m_h^2.[1] The association in Eq. (3) leads back to the EW measure since $m_{H_u}^2(\Lambda) + \delta m_{H_u}^2 = m_{H_u}^2(\text{weak})$.

3. Δ_{BG}

The EENZ/BG measure [18,19] (hereafter denoted simply by BG) is given by

$$\Delta_{BG} \equiv \max_i[c_i], \quad \text{where } c_i = \left| \frac{\partial \ln m_Z^2}{\partial \ln p_i} \right| = \left| \frac{p_i}{m_Z^2} \frac{\partial m_Z^2}{\partial p_i} \right|, \qquad (4)$$

where the p_i constitute the fundamental parameters of the model. Thus, Δ_{BG} measures the fractional change in m_Z^2 due to fractional variation in (high scale) parameters p_i. The c_i are known as *sensitivity coefficients* [19]. For the pMSSM (MSSM defined only at the weak scale), then explicit evaluation gives $\Delta_{BG} \simeq \Delta_{EW}$. For models defined

in terms of high scale parameters, the BG measure can be evaluated by expanding the terms on the right-hand side of Eq. (1) using semianalytic RG solutions in terms of fundamental high scale parameters [22]: for $\tan\beta = 10$ and taking $\Lambda = m_{\mathrm{GUT}}$, then one finds [23,24]

$$m_Z^2 \simeq -2.18\mu^2 + 3.84M_3^2 - 0.65M_3 A_t - 1.27m_{H_u}^2 \\ - 0.053m_{H_d}^2 + 0.73m_{Q_3}^2 + 0.57m_{U_3}^2 + \cdots \qquad (5)$$

The BG measure picks off the coefficients of the various terms and recales by the soft term squared over the Z-mass squared: e.g. $c_{m_{Q_3}^2} = 0.73 \cdot (m_{Q_3}^2/m_Z^2)$. If one allows $m_{Q_3} \sim 3$ TeV (in accord with requirements from the measured value of m_h) then one obtains $c_{m_{Q_3}^2} \sim 800$ and so $\Delta_{BG} \geq 800$. In this case, SUSY would be electroweak fine-tuned to about 0.1%.

The problem with most applications of the BG measure is that in any sensible model of SUSY breaking, the high scale SUSY parameters are not independent. For instance, in gravity-mediation, then for any given hidden sector, the SUSY soft breaking terms are all calculated as numerical coefficients times the gravitino mass [25–27]: e.g. $M_3(\Lambda) = a_{M_3}m_{3/2}$, $A_t = a_{A_t}m_{3/2}$, $m_{Q_3}^2 = a_{Q_3}m_{3/2}^2$, etc. where the a_i are just real constants. (For example, in string theory with dilaton-dominated SUSY breaking [26,27], then we expect $m_0^2 = m_{3/2}^2$ with $m_{1/2} = -A_0 = \sqrt{3}m_{3/2}$). The reason one scans multiple SUSY model soft term parameters is to account for a wide variety of possible hidden sectors. But this does not mean each soft term is independent from the others. By writing the soft terms in Eq. (5) as suitable multiples of $m_{3/2}^2$, then large positive and negative contributions can be combined/canceled and one arrives at the simpler expression [20]:

$$m_Z^2 = -2.18\mu^2(\Lambda) + a \cdot m_{3/2}^2. \qquad (6)$$

The value of a is just some number which is the sum of all the coefficients of the terms $\propto m_{3/2}^2$.[2] Using the BG measure applied to Eq. (6), then it is found that naturalness requires $\mu^2 \sim m_Z^2$ and also that $am_{3/2}^2 \sim m_Z^2$. The first requirement is the same as in Δ_{EW}. The second requirement is fulfilled *either* by $m_{3/2} \sim m_Z$ [19] (which seems unlikely in light of LHC Higgs mass measurement and sparticle mass bounds)

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[1]It is sometimes claimed that by using this method, then the SM would not be fine-tuned for large cutoff scales $\Lambda \gg 1$ TeV. However—in contrast to the SM—for the SUSY case, EW symmetry is not even broken at tree level in models where the soft terms arise from hidden sector SUSY breaking. Further discussion of the differences is included in Refs. [20,21].

[2]If μ is also computed as $\mu = a_\mu m_{3/2}$ as in the Giudice-Masiero mechanism [28], then $m_Z^2 = \text{const} \times m_{3/2}^2$ and $\Delta_{BG} \equiv 1$ and there would be no fine-tuning [29]. In other solutions of the SUSY μ-problem, such as Kim-Nilles [30], then μ is instead related to the Peccei-Quinn breaking scale and is expected to be independent. In the former case, then the responsibility is to find a suitable hidden sector which would actually generate m_Z^2 at its measured value. We are aware of no such models which even come close to that.

or that $m_{3/2}$ is large but the coefficient a is small [20]: i.e. there are large cancellations in Eq. (5). Since $\mu(\Lambda) \simeq \mu(\text{weak})$, then also $am_{3/2}^2 \simeq m_{H_u}^2(\text{weak})$ and so a low value of Δ_{BG} also requires a low value of $m_{H_u}^2$: i.e. $m_{H_u}^2$ is driven radiatively to small negative values. This latter situation is known as *radiatively-driven natural supersymmetry*, or RNS.

B. Naturalness and heavy SUSY Higgs bosons

The natural SUSY spectra is typified by a spectra of low-lying Higgsinos $\tilde{W}_1^{\pm}$, $\tilde{Z}_{1,2}$ with mass $\sim$100–300 GeV, the closer to m_Z the better, along with TeV-scale but highly mixed top-squarks $\tilde{t}_{1,2}$ [11,13]. The gluino mass can range between current LHC8 limits and about 4 TeV, and may well lie beyond LHC14 reach [31]. First/second generation matter scalars may well lie in the 5–30 TeV range, thus supplying at least a partial decoupling solution to the SUSY flavor, *CP*, proton decay and gravitino problem.[3] In addition, it should be clear from Eq. (1) that $m_{H_d}^2/\tan^2\beta \sim m_Z^2$ (a point mentioned previously in Ref. [32]). For m_{H_d} large, then one expects $m_A \sim m_{H_d}$. Requiring the term containing $m_{H_d}^2$ in Eq. (1) to be comparable to $m_Z^2/2$ or μ^2 then implies

$$m_A \sim |m_{H_d}^2|^{\frac{1}{2}} \lesssim |\mu| \tan\beta. \qquad (7)$$

Thus, for $|\mu| < 300$ GeV, we would expect for $\tan\beta = 10$ that $m_A \lesssim 3$ TeV. But for $\tan\beta$ as high as 50, we expect $m_A \lesssim 15$ TeV without becoming too unnatural.

In this paper, we explore the implications of SUSY naturalness for the heavy Higgs bosons of the MSSM: A, H and $H^{\pm}$. This topic has also been addressed in the recent paper [33]. In Ref. [33], using several different naturalness measures along with a low mediation scale $\Lambda \sim 10$–100 TeV and hard SUSY breaking contributions to the scalar potential, the authors conclude that heavy Higgs bosons should lie around the 1 TeV scale. Furthermore, since the heavy Higgs bosons are less susceptible to having hidden decay modes, their search should be an important component of the search for natural SUSY.

In this paper, we will arrive at quite different conclusions. In Sec. II, using the unified naturalness criteria, as embodied in Δ_{EW}, we will find that SUSY models which are valid all the way up to $\Lambda = m_{\text{GUT}} \simeq 2 \times 10^{16}$ GeV can be found with fine-tuning at the $\Delta_{\text{EW}} \sim 7$–30 level, corresponding to mild fine-tunings of just 3–15%. In this case, then as suggested in Eq. (7), we find that natural SUSY models exist with $m_A \lesssim 5$ TeV for $\tan\beta \lesssim 15$ while $m_A \lesssim 8$ TeV for $\tan\beta$ values ranging as high as 50–60. While the region $m_A \lesssim 1$ TeV should be accessible

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[3]Since $m_{\tilde{q},\tilde{\ell}} \sim m_{3/2}$, then we would expect $m_{3/2}$ also at the 5–30 TeV level.

to LHC heavy Higgs searches, the bulk of this mass range is well beyond any projected LHC reach. In Sec. III, we evaluate the heavy Higgs A, H and $H^{\pm}$ branching fractions as a function of mass for a benchmark case with radiatively-driven naturalness. Since for naturalness $\mu \sim 100$–300 GeV, then the heavy Higgs decays to Higgsino pairs is almost always open. Since the Higgsinos are essentially invisible at LHC, these modes lead to quasi-invisible decays. Further, since the heavy Higgs coupling to the -ino sector (here, -ino collectively refers to both charginos and neutralinos) is a product of gaugino times Higgsino components, then once kinematically accessible, the heavy Higgs tend to decay dominantly into gaugino plus Higgsino states. Such large branching fractions reduce the heavy Higgs branching fractions into SM modes, making standard heavy Higgs searches more difficult. Alternatively, since the gauginos tend to decay to gauge/Higgs bosons W, Z or h plus Higgsinos, then the qualitatively new decay modes arise: A, H, $H^{\pm} \to W$, Z or h plus missing E_T ($\cancel{E}_T$). These new decay modes—which are quite different than those expected in non-natural SUSY models with a binolike LSP—offer new avenues for heavy Higgs searches at LHC.

II. MASS BOUNDS FROM NATURALNESS

A simple mass bound from naturalness on heavy Higgs bosons can be directly read off from Eq. (1). The contribution to Δ_{EW} from the $m_{H_d}^2$ term is given by

$$C_{H_d} = m_{H_d}^2/(\tan^2\beta - 1)/(m_Z^2/2). \qquad (8)$$

Also the tree-level value of m_A is given by

$$m_A^2 = m_{H_u}^2 + m_{H_d}^2 + 2\mu^2 \simeq m_{H_d}^2 - m_{H_u}^2 \sim m_{H_d}^2 \qquad (9)$$

where the first partial equality holds when $\mu^2 \sim -m_{H_u}^2$ and the second arises when $m_{H_d}^2 \gg -m_{H_u}^2$. Combining these equations, then one expects roughly that

$$m_A \lesssim m_Z \tan\beta \Delta_{\text{EW}}^{1/2}(\text{max}) \qquad (10)$$

where $\Delta_{\text{EW}}(\text{max})$ is the maximal fine-tuning one is willing to tolerate. For $\Delta_{\text{EW}}^{-1} = 10\%$ fine-tuning with $\tan\beta = 10$, then one expects $m_A \lesssim 3$ TeV.

However, this simple argument is not the whole story since an assortment of radiative corrections are included in Eq. (1). In particular, the radiative corrections $\Sigma_u^u(\tilde{t}_{1,2})$ and $\Sigma_u^u(\tilde{b}_{1,2})$ (complete expressions are provided in the appendix of Ref. [13]) can become large and are highly $\tan\beta$ dependent.

To evaluate the range of m_A expected by naturalness, we will generate SUSY spectra using Isajet [34,35] in the 2-parameter nonuniversal Higgs model [36] (NUHM2) which allows for very low values of $\Delta_{\text{EW}} < 10$ (numerous

other constrained models are evaluated in Ref. [20] and always give much higher EW fine-tuning). The parameter space is given by

$$m_0, m_{1/2}, A_0, \tan\beta, \mu, m_A, \quad \text{(NUHM2)}. \quad (11)$$

The NUHM2 spectra and parameter spread vs Δ_{EW} were evaluated in Ref. [13] but with the range of m_A restricted to < 1.5 TeV. Here, we improve this scan by including a much large range of m_A:

$$\begin{aligned}
m_0 &: 0\text{--}20 \text{ TeV}, \\
m_{1/2} &: 0.3\text{--}2 \text{ TeV}, \\
-3 &< A_0/m_0 < 3, \\
\mu &: 0.1\text{--}1.5 \text{ TeV}, \\
m_A &: 0.15\text{--}20 \text{ TeV}, \\
\tan\beta &: 3\text{--}60.
\end{aligned} \quad (12)$$

We require of our solutions that:[4]
   (i) electroweak symmetry be radiatively broken (REWSB),
  (ii) the neutralino $\tilde{Z}_1$ is the lightest MSSM particle,
 (iii) the light chargino mass obeys the model independent LEP2 limit, $m_{\tilde{W}_1} > 103.5$ GeV [37],
 (iv) LHC search bounds on $m_{\tilde{g}}$ and $m_{\tilde{q}}$ are respected,[5]
  (v) $m_h = 125.5 \pm 2.5$ GeV.

The Isajet calculation of the various SUSY Higgs boson masses and couplings is performed using the renormalization group (RG) improved one-loop effective potential [38]. Since the calculation of the third generation quark and lepton Yukawa couplings is performed using 2-loop RGEs including complete one-loop MSSM/SM threshold corrections, then this method includes the dominant two-loop contributions [39]. The value of m_h computed with Isajet typically agrees to within 1 GeV with other "running programs" [40] such as SOFTSUSY, SPHENO and SUSPECT, but is typically 2–3 GeV below values predicted by nonrunning programs such as FEYNHIGGS [41]. While the value of m_h is sensitive to $\tan\beta$ for low values $\tan\beta \lesssim 8$, for $\tan\beta \gtrsim 10$, the value of m_h is relatively stable.

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[4]In addition, we could impose constraints from rare B decays: $BF(B_s \to \mu^+\mu^-)$ and $BF(b \to s\gamma)$. SUSY contributions to the former are proportional to $\tan^6\beta/m_A^4$ and so one expects large deviations at large $\tan\beta \sim 50$ and small m_A, far from our upper bounds. The latter has large contributions for sub-TeV scale top squarks which are disfavored by the value of $m_h \sim 125$ GeV which instead requires highly mixed top squarks in the few TeV range. Explicit predictions for these rates in the RNS model are shown in Fig. 10 of Ref. [13].

[5]Here, we have required m_0 and $m_{1/2}$ values in accord with recent search limits from Atlas [5] (20 fb^{-1} at LHC8) for the mSUGRA model with $\tan\beta = 30$ shown at ICHEP2014. These limits should apply also very nearly for the NUHM2 model for various $\tan\beta$ values.
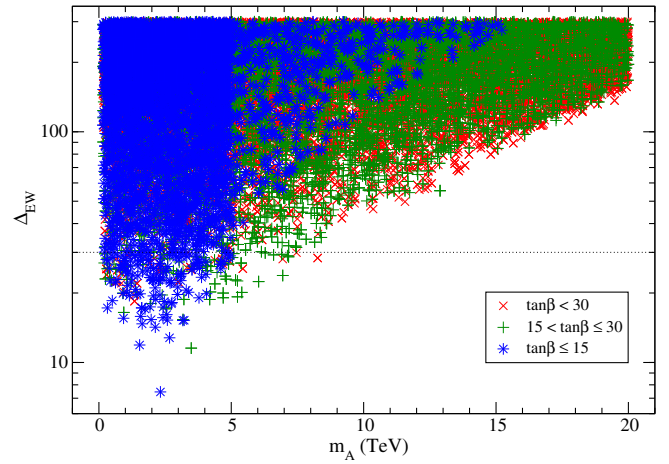

FIG. 1 (color online). Plot of Δ_{EW} vs. m_A from a scan over NUHM2 parameter space.

The results of our scan are shown in Fig. 1 where we plot Δ_{EW} vs. m_A. The dots are color-coded according to low, intermediate and high $\tan\beta$ values. From the plot, we see first that there is indeed an upper bound to m_A given by naturalness. In fact, for $\tan\beta < 15$ and $\Delta_{EW} < 10$, then indeed $m_A \lesssim 3$ TeV as suggested by the simple arguments above. For $\tan\beta > 15$, we do not generate any solutions with $\Delta_{EW} < 10$. For $\Delta_{EW} < 30$ (dotted horizontal line), then we have $m_A \lesssim 5$ TeV for $\tan\beta < 15$, and $m_A \lesssim 7(8)$ TeV for $\tan\beta < 30(60)$. While these values provide upper bounds on m_A from naturalness, we note that m_A values as low as 150–200 GeV can also be found. Since LHC14 searches for heavy Higgs are roughly sensitive to $m_A \lesssim 1$ TeV values [42], then we conclude that LHC14 searches will be able to probe a portion of natural SUSY parameter space, but perhaps the bulk of parameter space can easily lie well beyond Atlas/CMS search capabilities.

To gain more perspective on fine-tuning and how it depends on m_A and $\tan\beta$, we next adopt a proposed RNS benchmark point from Ref. [43]. This point has NUHM2 parameters given by

$$m_0 = 5 \text{ TeV}, \quad m_{1/2} = 0.7 \text{ TeV}, \quad A_0 = -8.3 \text{ TeV},$$

$$\tan\beta = 10, \quad \text{with} \quad \mu = 110 \text{ GeV} \quad \text{and} \quad m_A = 1 \text{ TeV}.$$

$$(13)$$

The value of Δ_{EW} is found to be 13.8. Here, we adopt this benchmark point, but now allow m_A and $\tan\beta$ as free parameters and plot color-coded ranges of Δ_{EW} in the m_A vs. $\tan\beta$ plane, as shown in Fig. 2.

From Fig. 2, we see that indeed the region with lowest Δ_{EW} occurs around $m_A \sim 1.2$–2.5 TeV with $\tan\beta \lesssim 10$. The yellow colored regions have $\Delta_{EW} < 50$. For these values, we find a more expansive region with $m_A \lesssim 6$ TeV and $\tan\beta \lesssim 20$. However, a second region with low $\Delta_{EW} < 50$ opens up at high $\tan\beta \sim 48$–52 with

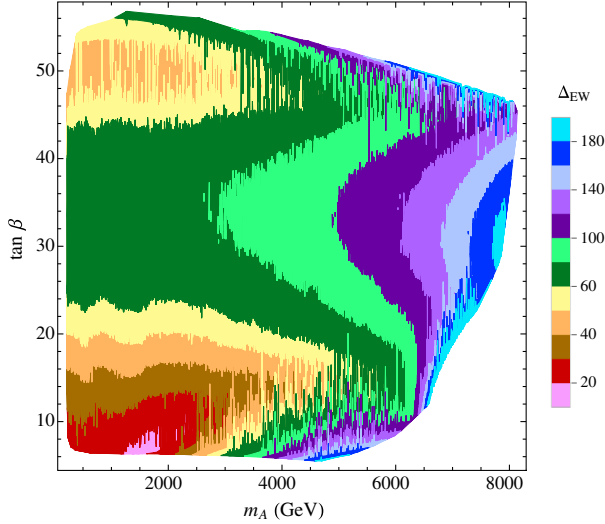FIG. 2 (color online).   Regions of SUSY naturalness, $\Delta_{\rm EW}$, in the m_A vs. $\tan\beta$ plane for the RNS benchmark point Eq. (13).

FIG. 3 (color online).   Branching fraction of A vs. m_A for the RNS benchmark point Eq. (13), but with variable m_A.

$m_A \lesssim 4$ TeV. The intermediate $\tan\beta \sim 20$–45 region has greater fine-tuning, where the maximal contributions to $\Delta_{\rm EW}$ we find arise from the radiative corrections $\Sigma_u^u(\tilde{b}_2)$.

III. IMPLICATIONS FOR HEAVY HIGGS DISCOVERY AT LHC

In many studies of the prospects for heavy Higgs boson discovery at the LHC, it is assumed that the standard model decay modes of A, H and $H^\pm$ are dominant. The prospects for discovery are usually presented in the m_A vs $\tan\beta$ plane. At next-to-leading-order in QCD, then the gluon fusion reactions $gg \to A, H$ are usually dominant out to $m_{A,H} \lesssim 1$ TeV while for higher masses then vector boson fusion (VBF) dominates [44]. The main discovery mode for $gg \to A, H$ is then the $A, H \to \tau^+\tau^-$ mode where the ditau mass can be reconstructed. Current search limits from Atlas and CMS exclude [45,46] $m_A \lesssim 950$ GeV for $\tan\beta$ as high as 50. For lower $\tan\beta$ values, the mass bounds are very much weaker (e.g. for $\tan\beta = 10$, then $m_A \gtrsim 300$ GeV is required). Production of heavy Higgs bosons in association with b-jets may aid the search [47]. In addition, the rarer decays into dimuons may also be possible [48,49], and recently dimuon signatures in association with b-jets have been explored [50,51].

The importance of heavy Higgs decay into SUSY modes was explored long ago [52] for the case where the LSP was usually taken to be a bino. If SUSY decay modes of H or A are open, then the SM branching fractions diminish while the new SUSY modes offer novel detection strategies [53].[6]

A. Heavy Higgs branching fractions in natural SUSY

The unique feature of SUSY models with radiatively-driven naturalness is the presence of light Higgsino states
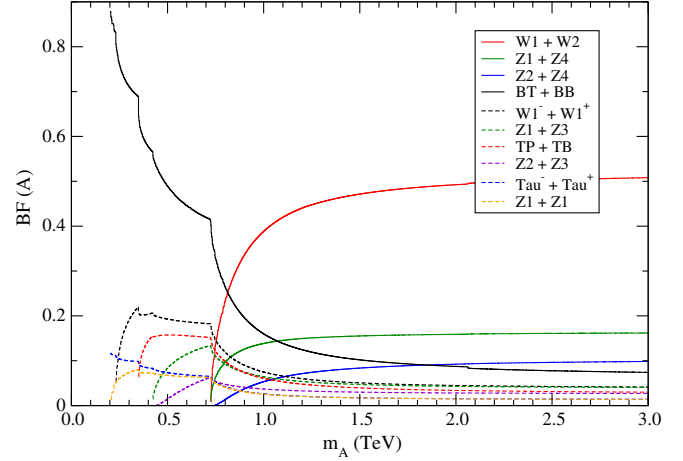
with mass $\sim$100–300 GeV, the closer to m_Z the better. This fact means that for most of the mass range of $m_{A,H}$, then SUSY decay modes should be open. Furthermore, the Higgsino-like LSP implies that the SUSY decay modes will generally be quite different than in earlier models where a binolike LSP was considered.

In Fig. 3, we show the branching fraction, as calculated by Isajet [34], of the pseudoscalar A boson vs m_A for the RNS benchmark point from Sec. II, but now with m_A taken as variable, with $\tan\beta = 10$. At low $m_A \sim 200$ GeV, SUSY decay modes are kinematically closed and $A \to b\bar{b}$ at $\sim$85% as is typical when SM decay modes are considered and the $t\bar{t}$ mode is closed. As m_A increases beyond 200 GeV, the $A \to$ Higgsino pairs mode opens up, and the SM branching fractions diminish. For $m_A \gtrsim 700$ GeV, the mixed Higgsino/wino mode $A \to \tilde{W}_1\tilde{W}_2$ turns on and rapidly dominates the branching fraction. This is because the SUSY Higgs coupling to -inos involves a product of gaugino component of one -ino times the Higgsino components of the other -ino,[7] and, in this case, $\tilde{W}_1$ is Higgsino-like and $\tilde{W}_2$ is winolike. For $m_A \gtrsim 1$ TeV, this decay mode is typically at the $\sim$50% level. For $m_A \gtrsim 1$ TeV, then the decays $A \to \tilde{Z}_1\tilde{Z}_4$ and $\tilde{Z}_2\tilde{Z}_4$ are also important. For TeV-scale values of m_A, the SM decay mode $A \to b\bar{b}$ drops to below the 10% level while $A \to \tau\bar{\tau}$ has dropped to the percent level. In this case, the search for heavy Higgs bosons utilizing SM decay modes will be much more difficult.

In Fig. 4, we show the branching fractions of the heavy scalar Higgs H vs m_H for the same RNS benchmark point. The overall behavior is similar to the case of the pseudoscalar A: at low values of m_H, the SM decay modes are dominant, but once m_H is heavy enough, the supersymmetric decay modes quickly open up and dominate the

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[6]For some recent studies, see [54,55].
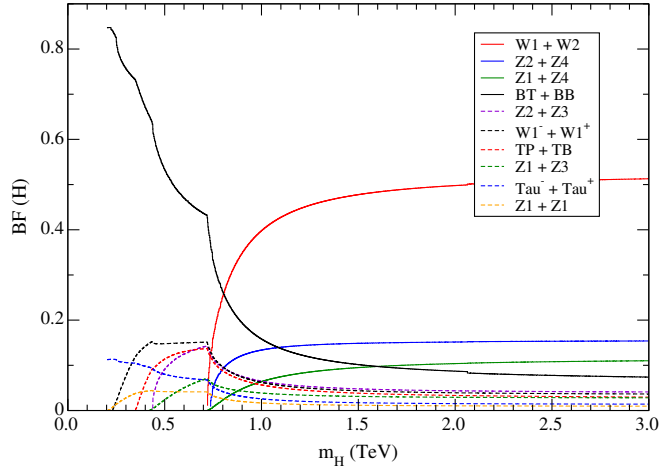
[7]See p. 178–179 of [56].

FIG. 4 (color online). Branching fraction of H vs. m_H for the RNS benchmark point Eq. (13), but with variable m_A.

branching fractions. At large m_H, the $H \to \tilde{W}_1 \tilde{W}_2$, $\tilde{Z}_2 \tilde{Z}_4$ and $\tilde{Z}_1 \tilde{Z}_4$ decay modes are dominant.

In Fig. 5, we show the branching fractions of H^+ vs m_{H^+} for the same RNS benchmark point. In this case, at low values of m_{H^+}, $H^+ \to t\bar{b}$ is dominant followed by $H^+ \to \tau^+ \nu_\tau$. As m_{H^+} increases, $H^+ \to \tilde{W}_2^+ \tilde{Z}_3$ turns on and later also $\tilde{W}_2^+ \tilde{Z}_1$, $\tilde{W}_1^+ \tilde{Z}_4$ and $\tilde{W}_2^+ \tilde{Z}_2$ turn on. At $m_{H^+} \gtrsim 1$ TeV, these latter decays into gaugino/Higgsino final states dominate. Such nonstandard decay modes will make searches for charged Higgs bosons more difficult than otherwise expected [57].

B. New SUSY Higgs signatures at LHC

1. $H, A \to W + E_T$

We have seen that for $m_{A,H} \gtrsim 1$ TeV, the dominant branching fraction is $H, A \to \tilde{W}_1^{\pm} \tilde{W}_2^{\mp}$. Since the $\tilde{W}_1$ is
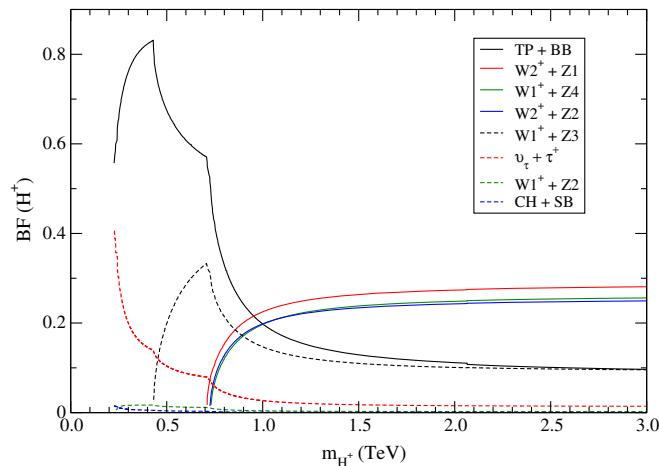
Higgsino-like, it tends to have only a small mass gap with the LSP: $m_{\tilde{W}_1} - m_{\tilde{Z}_1} \sim 10$–$20$ GeV. In this case, the visible energy from $\tilde{W}_1 \to f \bar{f}' \tilde{Z}_1$ decay (where f denotes SM fermions) is quite soft—most of the energy goes into making up the $\tilde{Z}_1$ rest mass—and so the Higgsinos are only quasivisible. On the other hand, the branching fractions for $\tilde{W}_2$ decay in the RNS model have been plotted out in Ref. [31] and found to be: $\tilde{W}_2 \to \tilde{W}_1 Z$, $\tilde{Z}_1 W$, $\tilde{Z}_2 W$ and $\tilde{W}_1 h$ each at about 25%. Thus, we expect s-channel H and A production to give rise to

$$gg \to H, \qquad A \to W^{\pm} + E_T \to \ell^{\pm} + E_T \qquad (14)$$

which is a rather unique signature for heavy Higgs boson production.

The dominant backgrounds come from direct W production followed by $W \to \ell \nu_\ell$ decay and also WZ production followed by $Z \to \nu \bar{\nu}$ and $W \to \ell \nu_\ell$. In Fig. 6, we plot the $e^+ + E_T$ transverse mass distribution from the signal using the RNS benchmark point with $m_A = 1$ TeV along with SM backgrounds. The signal from A, H production with $m_{A,H} \sim 1$ TeV and $\tan \beta = 10$ and 30 is well below background.

C. $H, A \to Z + E_T$

As mentioned above, $\tilde{W}_2 \to \tilde{W}_1 Z$ at ~ 30–35% in radiatively-driven natural SUSY. Thus, an alternative signature comes from

$$gg \to H, \qquad A \to Z + E_T \to \ell^+ \ell^- + E_T. \qquad (15)$$

The background to this process comes from ZZ production where one $Z \to \nu \bar{\nu}$ whilst the other goes as $Z \to \ell^+ \ell^-$. In Fig. 7 we plot the distribution in cluster transverse mass [58] $m_T(\ell^+ \ell^-, E_T)$ from heavy Higgs H, A production followed by their decays to $Z(\to \ell^+ \ell^-) + E_T$ from the
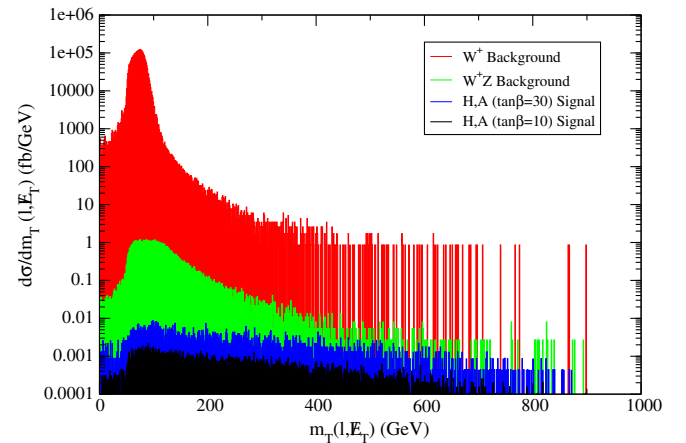

FIG. 5 (color online). Branching fraction of H^+ vs. m_{H^+} for the RNS benchmark point Eq. (13), but with variable m_A.

FIG. 6 (color online). Transverse mass distribution for e^+ plus E_T events at LHC14 from W^+, $W^+ Z$ and A, H production from benchmark point Eq. (13).
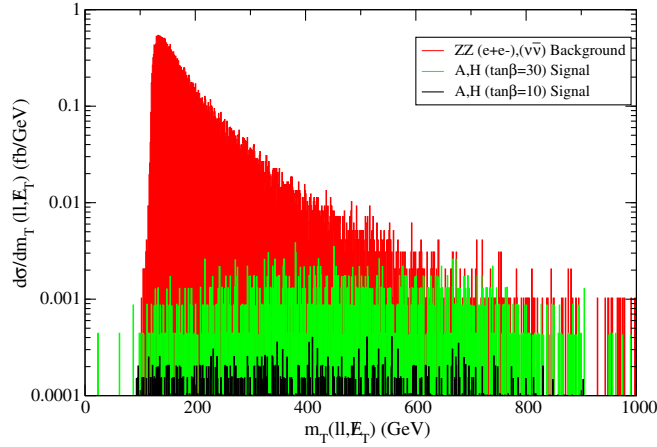
FIG. 7 (color online). Dilepton cluster transverse mass distribution for e^+e^- plus E_T events at LHC14 from ZZ and A, H production from benchmark point Eq. (13).

RNS benchmark point for $m_A = 1$ TeV along with ZZ background. Here, we see that signal from A, H production with $m_{A,H} \sim 1$ TeV lies well below the diboson background for $\tan\beta = 10$. If we increase $\tan\beta$ to 30, signal and BG become comparable at very large $m_T(\ell^+\ell^-, E_T)$, although, in this range, the event rate is quite limited.

D. $H, A \rightarrow h + E_T$

A third possible signature consists of $A, H \rightarrow \tilde{Z}_{1,2}\tilde{Z}_{3,4}$ where $\tilde{Z}_{3,4} \rightarrow \tilde{Z}_{1,2}h$ resulting in a $(h \rightarrow b\bar{b}) + E_T$ signature. We expect such a signal to lie well below backgrounds from Zh and ZZ production.

IV. CONCLUSIONS

In this paper, we have examined the implications of SUSY naturalness for the heavy Higgs boson sector. We use the $\Delta_{\rm EW}$ measure of naturalness, although we show that—properly applied—the Higgs mass fine-tuning, and also the EENZ/BG fine-tuning, would give similar results since

$$\Delta_{\rm HS} \simeq \Delta_{\rm BG} \simeq \Delta_{\rm EW}, \qquad (16)$$

so long as *dependent* terms are properly combined before evaluating naturalness.

Using the $\Delta_{\rm EW}$ measure, we find upper bounds on the heavy Higgs masses: for 10% fine-tuning and $\tan\beta \sim 10$, we expect $m_A \lesssim 2.5$ TeV whilst for 3% fine-tuning and $\tan\beta$ as high as 50, $m_A \lesssim 8$ TeV. These values are considerably larger than the range depicted in Ref. [33] and much of this range likely lies beyond LHC14 reach. This means LHC14 will be able to probe only a portion—but by no means all—of natural SUSY parameter space via heavy Higgs boson searches.

Furthermore, since light Higgsino states $\tilde{W}_1^\pm$ and $\tilde{Z}_{1,2}$ are expected to have mass $\sim$100–300 GeV (the closer to m_Z the more natural), there will almost always be supersymmetric decay modes open to the heavy SUSY Higgs states. We evaluated these branching fractions and find that they can in fact be the dominant decay modes, especially if $m_{A,H} > m_{\tilde{W}_1} + m_{\tilde{W}_2}$, in which case this decay mode tends to dominate. The supersymmetric decay modes diminish the SM decay modes of H, A and $H^\pm$ making standard search techniques more difficult for a specified heavy Higgs mass. However, qualitatively new heavy Higgs search modes appear thanks to the supersymmetric decay modes. Foremost among these are the decays $H, A \rightarrow \tilde{W}_1^\pm \tilde{W}_2^\mp$ which results in final states characterized by W, Z or h plus E_T. These new signatures seem to be rather challenging to extract from SM backgrounds which occur at much higher rates. It may well be that forward b-jet tagging in $bg \rightarrow bA$ or bH production or $gb \rightarrow tH^+$ production followed by $A, H, H^\pm \rightarrow$ SUSY decays could ameliorate the situation.

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