Dark Matter and Collider Phenomenology with two light Supersymmetric Higgs Bosons

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Recently, it has been pointed out that two different excesses of events observed at LEP could be interpreted as the CP-even Higgs bosons of the MSSM with masses of approximately 98 and 114 GeV. If this is the case, the entire MSSM Higgs sector is required to be light. In this article, we explore such a scenario in detail. We constrain the Higgs and supersymmetric spectrum using B physics constraints as well as the magnetic moment of the muon. We then point out the implications for neutralino dark matter — next generation direct detection experiments will be sensitive to all MSSM models with such a Higgs sector. Finally, we find that all models outside of a very narrow corridor of the parameter space have a charged Higgs boson which will be observed at the LHC. In those exceptional models which do not contain an observable charged Higgs, a light top squark will always be seen at the LHC, and likely at the Tevatron.

I. SETTING THE STAGE

The four experiments at LEP have searched for Higgs bosons up to a mass of approximately 115 GeV [1]. Although no strong indication of a Higgs has been detected by LEP, excesses with statistical significances of 2.3σ and 1.7σ have been reported for Higgs-like events corresponding to masses of 98 GeV and 115 GeV, respectively. Very recently, it has been pointed out by Manuel Drees that both of these excesses reported by LEP could be accommodated within the Minimal Supersymmetric Standard Model (MSSM) as signatures of the two CP-even Higgs bosons [2].

One of the reason that the LEP excess at 98 GeV has received so little attention might be that the rate observed corresponding to this mass is about a factor of ten below that which would be expected from a Standard Model Higgs boson. In the MSSM, however, the $h$-$Z$-$Z$ coupling is suppressed relative to the value in the Standard Model by a factor of $\sin(\beta - \alpha)$, where $\alpha$ is the mixing angle between the two CP-even Higgs bosons and $\tan\beta$ is the ratio of the two vacuum expectation values. To predict the number of events seen at LEP for a 98 GeV light Higgs boson, we constrain $0.056 < \sin^2(\beta - \alpha) < 0.144$ (the precise range of $\sin^2(\beta - \alpha)$, $m_h$ and $m_H$ we allow is the same as used in Ref. [2] to allow for easy comparison of results). Meanwhile, the analogous coupling of the heavy CP-even Higgs is only slightly suppressed below the Standard Model value and is consistent with the number of events reported by LEP corresponding to a 115 GeV Higgs boson. These two excesses together yield a combined statistical significance of $3.1\sigma$ [2].

The mass scale of the MSSM Higgs sector is usually fixed by the mass of the pseudoscalar Higgs boson $m_A$. In the well-known leading $m_t^4$ approximation [3, 4], we can very roughly link the pseudoscalar mass to the two scalar Higgs masses in the limit of non-mixing top squarks:

$$m_h^2 + m_H^2 = m_A^2 + m_Z^2 + \epsilon$$
$$m_{H^\pm}^2 = m_A^2 + m_W^2$$
$$\epsilon = \frac{3G_F}{\sqrt{2}\pi^2} \frac{m_t^4}{\sin^2\beta} \log \frac{m_t^2}{m_h^2}$$

In this approximation $m_t$ is the top squark mass, $m_{h,A,H}$ are the neutral Higgs masses and $m_{H^\pm}$ is the charged Higgs mass. We immediately see that the existence of two light scalar Higgs bosons leads to all of the MSSM Higgs bosons being fairly light. This can lead to interesting phenomenology in the MSSM. In this article, we consider the constraints on this scenario from the $B \to X_s\gamma$ and $B_s \to \mu^+\mu^-$ branching fractions and the magnetic moment of the muon. We discuss the implications of such a model on the phenomenology of neutralino dark matter, and find that the prospects for the direct detection of neutralinos in elastic scattering experiments to be excellent. We then study the spectra of the superpartners and the Higgs sector with two light Higgs scalars and discuss in detail the prospects for discovering the charged Higgs boson at the LHC.
In order to study the phenomenology of supersymmetric models with light Higgs scalars, we first performed a scan over the relevant parameters of the MSSM. We have varied all masses up to 6 TeV, and \( \tan \beta \) between 1 and 60. We did not assume any specific supersymmetry breaking scenario or unification scheme. To take into account the radiative corrections to the Higgs masses, we have used FeynHiggs [5], which includes all contributions up to the two-loop level. We find that this level of precision is needed to obtain an accurate representation of the MSSM phenomenology within this class of models. We give a detailed analysis of the Higgs sector in Section IV.

In Fig. 1, we confirm the conclusion of Ref. [2] that the entire MSSM Higgs sector is required to be rather light to accommodate both excesses reported by LEP. In the overwhelming majority of models found, the mass of the pseudoscalar Higgs bosons is between 95 and 130 GeV and the mass of the charged Higgs lies between 110 and 150 GeV. Although we do find some points outside of this range, in particular a trail of models extending to the upper right of the figure (left frame), these are quite rare compared to the dense collection of models near \( m_A \sim 110 \text{ GeV} \) and \( m_{H^\pm} \sim 130 \text{ GeV} \). For each point shown, we require 95 GeV \( < m_h < 101 \text{ GeV} \), 111 GeV \( < m_H < 119 \text{ GeV} \) and \( 0.056 \leq \sin^2(\beta - \alpha) \leq 0.144 \), corresponding to the range matching the observations at LEP. The black points are consistent with measurements of \( B \rightarrow X_s \gamma \) at the 3\( \sigma \) level and do not violate the Tevatron constraint on the \( B_s \rightarrow \mu^+ \mu^- \) branching fraction. Blue points violate \( B_s \rightarrow \mu^+ \mu^- \), but are consistent with \( B \rightarrow X_s \gamma \). Green points violate \( B \rightarrow X_s \gamma \). These constraints will be discussed further in section II.

Constraints from rare \( B \)-decays and the magnetic moment of the muon are some of the most useful tools we currently have to guide our studies of supersymmetric phenomenology. The \( B \) physics constraints gain their power mostly because of a possible \( \tan \beta \) enhancement of the bottom Yukawa couplings. In this section, we consider some of these constraints within the context of models with 98 and 115 GeV Higgs bosons.

The branching fraction of \( B \) decays to a strange state plus a photon has been measured by the BELLE [6], CLEO [7] and ALEPH [8] experiments. The weighted average of these experiments’ results indicate \( \text{BR}(B \rightarrow X_s \gamma) = (3.25 \pm 0.37) \times 10^{-4} \). In comparison, the Standard Model prediction for this transition rate is \( (3.70 \pm 0.30) \times 10^{-4} \) [9].

The supersymmetric processes which are most likely to contribute substantially to this branching ratio involve a charged Higgs or a chargino. These contributions are enhanced by powers of the Yukawa coupling \( m_b \tan \beta \) for large values of \( \tan \beta \). In passing we emphasize that there are additional \( \Delta \phi \) corrections [10] which can have huge effects if the MSSM spectrum is split between light gluinos or higgsinos and heavier sbottoms, but we will see that this is not the part of parameter space in which we are interested in. If the gluino mass is smaller than 300 GeV, the LHC will be swamped by gluino pair production with cross sections as large as 1000 pb and the effect of the gluino mass on the charged Higgs boson is negligible.
In Fig. 2, we plot $\text{BR}(B \to X_s \gamma)$ versus $\tan \beta$ (left frame), the lightest chargino mass (center frame) and the charged Higgs mass (right frame). Shown as horizontal dashed lines are the $2\sigma$ confidence bounds on this quantity as measured by BELLE, CLEO and ALEPH. Again, all points shown have $95 \text{GeV} < m_h < 101 \text{GeV}$, $111 \text{GeV} < m_H < 119 \text{GeV}$ and $0.056 \lesssim \sin^2(\beta - \alpha) \lesssim 0.144$. Black points do not violate the Tevatron constraint on the $B_s \to \mu^+ \mu^-$ branching fraction. Green points violate this constraint.

Figure 2: The $b \to s\gamma$ branching fraction as a function of $\tan \beta$, the lightest chargino mass, and the charged Higgs mass. Shown as horizontal dashed lines are the $2\sigma$ confidence bounds on this quantity as measured by BELLE, CLEO and ALEPH. Again, all points shown have $95 \text{GeV} < m_h < 101 \text{GeV}$, $111 \text{GeV} < m_H < 119 \text{GeV}$ and $0.056 \lesssim \sin^2(\beta - \alpha) \lesssim 0.144$. Black points do not violate the Tevatron constraint on the $B_s \to \mu^+ \mu^-$ branching fraction. Green points violate this constraint.

Limits on the branching fraction of the rare decay $B_s \to \mu^+ \mu^-$ can also be exploited to limit the allowed MSSM parameter space. The Standard Model prediction for this branching ratio is $\text{BR} = (2.4 \pm 0.5) \times 10^{-9}$. In particular, the Higgs-induced MSSM contributions to this branching fraction scales with $\tan^6 \beta/m_A^2$ and tends to be especially large in our scenario with a light Higgs sector [11]. If we require the Higgs sector to be light, models with very large values of $\tan \beta$ are likely to violate the constraints on this quantity placed at the Tevatron [12]. In Fig. 3, we plot $\text{BR}(B_s \to \mu^+ \mu^-)$ versus $\tan \beta$. Shown as a dashed horizontal line is the 90% upper confidence bound placed by the Tevatron experiments. The Standard Model value is clearly visible as the low $\tan \beta$ nose in the distribution of the MSSM parameter points — small $\tan \beta$ means small contributions from the MSSM Higgs sector. The upper limit from the Tevatron becomes an issue for $\tan \beta > 10$ but it really only develops seriously destructive power for $\tan \beta > 50$.

Combining the upper limit on the $B_s \to \mu^+ \mu^-$ decay and the measurement of $B \to s\gamma$ makes it considerably harder for us to find viable MSSM scenarios with two light Higgs scalars. On the one hand, we need an MSSM contribution to $B \to X_s \gamma$ from the chargino sector with the right sign and magnitude to compensate for the charged Higgs diagrams. This is possible by exploiting the enhancement by one power of $\tan \beta$ in the chargino amplitude. On the other hand, we do not want too large a contribution to $B_s \to \mu^+ \mu^-$ from the same sector which scales as $\tan^6 \beta$. Putting both of these constraints together, we favor fairly large, but not too large, values of $\tan \beta$ and at the same time accept some light superpartners to accommodate $B \to X_s \gamma$. We should stress, however, that this argument is not strict in the sense that these two constraints guarantee a light MSSM mass spectrum. By choosing properly tuned parameters, we can still get by with a spectrum where the lightest of the charginos and stops weighs more than 1 TeV.

Looking into the near future, we see that of course the LHC prospects of seeing these light-Higgs MSSM scenarios in $B$ physics are excellent. Simulations of CMS, ATLAS and LHCb events probing $\text{BR}(B_s \to \mu^+ \mu^-)$ predict tens of
One very attractive feature of R-parity conserving supersymmetry is that it naturally provides a stable particle which, in many models, can be a viable candidate for dark matter. The lightest neutralino is particularly appealing in this respect [20].
Neutralinos can annihilate through a variety of channels, including through the exchange of CP-even or odd Higgs bosons, charginos, neutralinos, sfermions and gauge bosons. Which annihilation channel(s) dominates varies from model to model. In the scenario we are studying here, however, the presence of light Higgs bosons suggest that s-channel Higgs exchange to fermion pairs is likely to be a particularly efficient annihilation channel, especially in those models with moderate to large values of $\tan\beta$. The formulae for these annihilation channels are collected in the appendix.

The annihilation cross section can be used to calculate the thermal relic abundance present today:

$$\Omega_{\chi^0} h^2 \approx \frac{10^9}{M_{Pl}} \frac{x_{FO}}{g_*} \frac{1}{(a + 3b/x_{FO})},$$

where $g_*$ is the number of relativistic degrees of freedom available at freeze out, $a$ and $b$ are the amplitudes given in the appendix, and $x_{FO}$ is the value at freeze out:

$$x_{FO} \approx \ln \left( \frac{45 m_{\chi^0} M_{Pl}(a + 6b/x_{FO})}{8 \pi^3 \sqrt{g_*}} \right).$$

Over the range of cross sections and masses expected for the lightest neutralino, $x_{FO} \approx 20 - 30$. To yield the density of cold dark matter measured by WMAP [23], we thus require $a + 3b/x_{FO} \approx 3 \times 10^{-26} \text{cm}^3/\text{s}$. The main contribution to $a$ and $b$ are s-channel CP-odd and CP-even Higgs exchange, respectively. From the amplitudes given in the appendix, we find $a \sim b \times \tan^2 \beta/3$, so we conclude that the $b$-term induced through CP-even Higgs plays a subdominant role in the relic density calculation to the CP-odd Higgs induced $a$-term. In the left frame of Fig. 4, we plot the relic density of the lightest neutralino versus its mass. We have calculated this quantity using the micrOMEGAs program [19].

Over the vast majority of supersymmetric parameter space, the composition of the lightest neutralino is dominated by its bino component, although a small admixture of wino or higgsino is common. For a bino-like neutralino with large or moderate $\tan\beta$, the annihilation cross section through pseudoscalar Higgs exchange scales as $\epsilon_u^2 \tan^2 \beta$. This up-type higgsino component can be approximated in most models by $\epsilon_u^2 \approx \epsilon_d^2 \sin^2 \theta_W \sin^2 \beta m_Z^2 / \mu^2$. With this in mind we plot the neutralino relic density versus the quantity $\tan^2 \beta \sin^2 \theta_W m_Z^2 / \mu^2$ in the right frame of Fig. 4. The two solid lines shown are the analytic results for $m_A = 120 \text{ GeV}$ and $m_{\chi^0} = 50$ (bottom) and 500 (top) GeV in this approximation. The majority of models we find fall within this range, suggesting that their annihilation is in fact dominated by and s-channel pseudoscalar Higgs exchange to $bb$. As expected, no models fall above this range. In the models which lie below this range, another annihilation mode or modes (such as t-channel sfermion exchange, s-channel $Z$ exchange or s-channel pseudoscalar Higgs exchange to top quark pairs) must contribute substantially, thus lowering the relic abundance accordingly. Coannihilations between the lightest neutralino and another superpartner may also reduce the relic density.
The process is roughly given by the scattering cross section dominated by the t-channel exchange of CP-even Higgs bosons. The cross section for this process is often capable of generating the measured density of dark matter within a factor of ten of the observed quantity. The dashed curve is the current limit placed by the CDMS experiment.

From the right frame of Fig. 4, we can infer that for those models which generate the measured relic abundance (those which fall along the horizontal dashed line), most appear to annihilate substantially through pseudoscalar Higgs exchange to \( b\bar{b} \). Models with a bino-like neutralino with a small higgsino admixture and moderate to high values of \( \tan \beta \) are often capable of generating the measured density of dark matter.

The prospects for the direct detection of particle dark matter are quite encouraging for a neutralino in conjunction with a light Higgs sector. In many of the models we study here, the spin-independent neutralino-nucleon elastic scattering cross section is dominated by the t-channel exchange of CP-even Higgs bosons. The cross section for this process is roughly given by

\[
\sigma_{\chi N} \approx \frac{8G_F^2m_H^2}{\pi m_H^2} \left( \frac{m_p m_N}{m_p + m_N} \right)^2 \left( -\epsilon_B \sin \theta_W + \epsilon_W \cos \theta_W \right)^2 \epsilon_H^2 \left[ \sum_q \frac{m_q}{\cos \beta} \langle N|q\bar{q}|N \rangle \right]^2,
\]

where the sum is over quark types and \( \langle N|q\bar{q}|N \rangle \) are the matrix elements over the nucleonic state. The quantity \( \cos \beta \) should be replaced with \( \sin \beta \) for up-type quarks. \( \epsilon_B, \epsilon_W \) and \( \epsilon_H \) are the bino, wino and higgsino fractions of the neutralino. In the sum over quark species, the strange quark contribution dominates with \( m_s \langle N|s\bar{s}|N \rangle \approx 0.2 \) GeV. For models with ~100 GeV Higgs bosons and a bino-like LSP, we estimate

\[
\sigma_{\chi N} \sim 10^{-7} - 10^{-6} \text{ pb} \left( \frac{\epsilon_H}{0.01} \right) \left( \frac{\tan \beta}{20} \right)^2.
\]

Of course other elastic scattering channels (squark exchange in particular) can also contribute substantially. In Fig. 5 we show the spin-independent neutralino-nucleon elastic scattering cross section for models in this scenario. Models shown in the right frame each generate an abundance of neutralino dark matter within one order of magnitude of the observed density, while those models shown in the left frame may or may not. The dashed curve is the current limit placed by the CDMS II experiment. The dotted yellow, solid red, dotted blue and solid magenta curves (from top to bottom) are the approximate projected limits of the GERDA [25], CDMS II, ZEPLIN-MAX [26] and Super-CDMS (phase III) experiments, respectively. Edelweiss [27] should also reach a sensitivity similar to that of GERDA [25], CDMS II, ZEPLIN-MAX and Super-CDMS (phase III) experiments, respectively. The black points are consistent with measurements of \( B \rightarrow X_s \gamma \) at the 3\( \sigma \) level and do not violate the Tevatron constraint on the \( B_s \rightarrow \mu^+\mu^- \) branching fraction. Blue points violate \( B_s \rightarrow \mu^+\mu^- \), but are consistent with \( B \rightarrow X_s \gamma \). Green points violate \( B \rightarrow X_s \gamma \).

From Fig. 5, it is obvious that the prospects for direct detection are excellent for this class of supersymmetric models with light Higgs sectors which often dominate the cross section. Most of the favored models are within the reach of CDMS and Edelweiss experiments and those remaining models below this level of sensitivity should be within the reach of next generation experiments such as ZEPLIN-MAX or Super-CDMS.

For models with very small elastic scattering cross sections do not appear in Fig. 5. If we limit the magnitude of \( \mu \) to be less than a few TeV for fine tuning reasons, the higgsino fraction of the
lightest neutralino is then bounded from below to be $r_a^2 \gtrsim \sin^2 \theta_W \sin^2 \beta \frac{m_Z^2}{(3 \text{ TeV})^2} \sim 0.0002$. From this, we see that $\sigma_{\chi N}$ cannot be smaller than $10^{-9}$ to $10^{-10}$ pb without $|\mu|$ being unnaturally large.

In addition to these excellent prospects for direct dark matter searches, the characteristics of the lightest neutralino in this class of models are fairly ideal for the purposes of indirect detection. Since the annihilation cross section during freeze out is dominated by the first term in the expansion, $(\sigma v) = a + bx + \mathcal{O}(x^2)$, the annihilation cross section relevant for neutralinos annihilating near the galactic center, throughout the galactic halo, or elsewhere of interest to indirect detection, is the maximum value consistent with the measured relic abundance, $(\sigma v) \approx 3 \times 10^{-26} \text{ cm}^3/\text{s}$. This makes it rather likely that such models with be within the reach of future indirect detection experiments searching for neutralino annihilation products in the form of gamma-rays [29] or anti-matter [30]. The prospects for detecting such a neutralino in future cosmic positron experiments are especially promising.

IV. LIGHT MSSM HIGGS BOSONS AT THE LHC

Before we begin discussing the prospects of observing a light MSSM Higgs sector at the LHC, we should point out that it has been known for a long time that one CP-even Higgs scalar in the MSSM is guaranteed to be seen in weak-boson-fusion production with a subsequent decay $h, H \rightarrow \tau \tau$ [31]. The search strategy is identical to the Standard Model search for low Higgs masses and has been extensively studied including detector effects [32]. In the special case of three light MSSM Higgs scalars, this discovery channel is known to face a challenge: the $\tau \tau$ mass resolution will not be sufficient to resolve the two CP-even scalars if their mass difference is less than $\mathcal{O}(5) \text{ GeV}$ — they will appear as one sightly wider Higgs resonance. Moreover, the pseudoscalar Higgs only couples to gauge bosons through a dimension-5 operator [33], thus it will not be produced in weak boson fusion. The case of three almost mass-degenerate light Higgs scalars has been specifically studied, with an emphasis on distinguishing the three mass peaks in the decay to muons [34]. The conclusion is that for Higgs masses below 140 GeV, the inclusive Higgs search is challenging and that the heavy Higgs bosons will probably not be separately observable in this channel. The more promising strategy is to look for bottom-Higgs associated production with a subsequent Higgs decay to muons and at least one $b$-tag. The Yukawa coupling for the heavy scalar and the pseudoscalar Higgs can then be $\tan \beta$ enhanced and the mass peaks might be resolved for $\tan \beta \gtrsim 30$ and $m_A \gtrsim 130 \text{ GeV}$.

The general lesson we learn from detailed studies by ATLAS and CMS is that it might well be easier to discover a charged Higgs boson than the additional neutral scalars. A charged Higgs can be searched for either in anomalous top decays, or it can be directly produced as $bg \rightarrow tH^-$. In both cases the decay to a tau lepton is most promising [35]. The coverage of models in the $(m_A - \tan \beta)$ plane from this channel has historically included a hole for charged Higgs masses just above the top mass. We emphasize that this hole does not mean that the charged Higgs will be missed at the LHC in this parameter range. It just means that we will have to combine the search strategies for (off-shell) anomalous top decays with the usual associated production [36]. For the purpose of this paper we will use a preliminary parton level analysis [37]. We include this result in Fig. 6.

A. Light charged Higgs bosons

Phenomenologically speaking, the easiest searches for charged Higgs bosons are in the mass range $m_{H^\pm} \lesssim m_t - m_b$. A light charged Higgs will almost exclusively decay to a tau lepton, because the decays to $W h^0$ and $tb$ are kinematically closed in our scenario. This means we can look for top decays to $b \ell \bar{\nu}$, where there are too many final state taus as compared to muons and electrons. In top pair production we let the second top decay into $Wb$, where the $W$ can decay leptonically and hadronically (the latter channel dominates the reach in the charged Higgs mass) [38]. The tau from the charged Higgs typically forms a jet. The partial width of a top decay to a charged Higgs contains two terms: $m_b^2 \tan^2 \beta + m_b^2 / \tan^2 \beta$. Hence, there will be large branching fractions both for small and large values of $\tan \beta$ with the weakest point around $\tan \beta \sim \sqrt{m_t/m_b} \sim 7$. Because at the LHC these searches are sensitive to anomalous branching fractions of the order of $1\%$ [40], we can expect them to probe charged Higgs masses up to $m_{H^\pm} \lesssim m_t - m_b \sim 160 \text{ GeV}$. Including off-shell effects, this reach might even be extended by a few GeV. At the Tevatron the same searches are currently in progress, but are strongly statistics limited [39]. They only cover anomalous branching fractions of the order of $60\%$, corresponding to an enhanced Yukawa coupling $\tan \beta \gtrsim 20$ and masses $m_{H^\pm} \lesssim 140 \text{ GeV}$.

In Fig. 6 we overlay our MSSM parameter points with the LHC (ATLAS) $5\sigma$ reach. The magenta line corresponds to the low-luminosity results from anomalous top decays. The slight bending comes from the mass relation between $m_A$ and $m_{H^\pm}$ as well as from the coupling suppression for intermediate $\tan \beta$. We see that most of the parameter points with two light Higgs scalars will be clearly visible in anomalous top decays at the LHC, possibly even at the Tevatron.
Higgs, we are automatically driven into the region where the charged Higgs can be found in this region of parameter space. Instead, we emphasize that closing this hole is a crucial task for the near future. We perform a dedicated scan around the one point which appears in Fig. 6 close to the hole in the LHC reach for $\beta$ just above the top threshold and small $\tan\beta$ in anomalous top decays and in the associated production with a top are challenging. It requires charged Higgs masses their discovery prospects at the LHC in more detail next, to ensure there is no hole in the LHC discovery range. One point somewhat close to this region in our scan). However, we will study these kinds of parameter points and analyses are actually performed on data. In Fig. 6 the densities of points suggests that it is very unlikely to find such this kind of study is ongoing in CMS and in ATLAS [37] and we do not expect this to be a problem once these.

160 and 200 GeV. As mentioned above, this hole is an artifact of the two search strategies meeting in this mass range.

The best-studied regime for charged Higgs searches at the LHC is the region of large $\tan\beta \gtrsim 10$ and charged Higgs masses well above the top threshold. Following the approximation in Eq.(1), we see that the required two light scalar Higgs bosons imply $\sqrt{m_A^2 + m_Z^2 + \epsilon} \sim 150$ GeV. Requiring charged Higgs masses above the top threshold translates into $m_A > 150$ GeV, or in other words $m_Z^2 \sim -\epsilon$. This could be possible if one stop were much lighter than the top. On the other hand, the $\epsilon$ approximation is based on the assumption of non-mixing stops, so we postpone this light-stop scenario to the next subsection. The other way of enhancing the mass difference between the average CP-even Higgs mass and the charged Higgs mass (beyond the simple $m_A^2$ approximation) and to escape the anomalous top decays is to create large bottom Yukawa corrections through large values of $\tan\beta$ (while keeping the $\mu$ parameter small). This leads to additional terms contributing to the charged Higgs mass squared, the leading terms of which are proportional to $m_{\tilde{t}_1}^2$. The two kinds of contributions which appear with this mixed Yukawa coupling are either proportional to $\log(m_{\tilde{t}_1}^2 m_{\tilde{b}_1}^2/m_t^2)$, or they come without this logarithm and are directly proportional to the average stop mass instead [4]. We see their effect in Fig. 6 when we compare the $\tan\beta$ dependence of $m_A$ and $m_{H^\pm}$. For all value of $\tan\beta$ the requirement of two light Higgs scalar limits the pseudoscalar mass to a narrow corridor. We find hardly any SUSY scenarios with $m_A > 110$ GeV. By considering large values of $\tan\beta$, we can increase the charged Higgs mass by 50 GeV and avoid the top decay threshold.

Because large values of $\tan\beta$ are required to at the same time have two light Higgs scalars and a heavier charged Higgs, we are automatically driven into the region where the charged Higgs can be found in $tH^-$ production at the LHC. The only issue is the hole which usually appears in the LHC coverage region for charged Higgs masses between 160 and 200 GeV. As mentioned above, this hole is an artifact of the two search strategies meeting in this mass range. The associated production process and the anomalous (off-shell) top decay have to be combined to cover this hole. This kind of study is ongoing in CMS and in ATLAS [37] and we do not expect this to be a problem once these analyses are actually performed on data. In Fig. 6 the densities of points suggests that it is very unlikely to find such a light-Higgs MSSM scenario where the LHC sees only one neutral Higgs scalar and no charged Higgs (we find only one point somewhat close to this region in our scan). However, we will study these kinds of parameter points and their discovery prospects at the LHC in more detail next, to ensure there is no hole in the LHC discovery range.

B. Heavy charged Higgs bosons with large $\tan\beta$

The best-studied regime for charged Higgs searches at the LHC is the region of large $\tan\beta \gtrsim 10$ and charged Higgs masses well above the top threshold. Following the approximation in Eq.(1), we see that the required two light scalar Higgs bosons imply $\sqrt{m_A^2 + m_Z^2 + \epsilon} \sim 150$ GeV. Requiring charged Higgs masses above the top threshold translates into $m_A > 150$ GeV, or in other words $m_Z^2 \sim -\epsilon$. This could be possible if one stop were much lighter than the top. On the other hand, the $\epsilon$ approximation is based on the assumption of non-mixing stops, so we postpone this light-stop scenario to the next subsection. The other way of enhancing the mass difference between the average CP-even Higgs mass and the charged Higgs mass (beyond the simple $m_A^2$ approximation) and to escape the anomalous top decays is to create large bottom Yukawa corrections through large values of $\tan\beta$ (while keeping the $\mu$ parameter small). This leads to additional terms contributing to the charged Higgs mass squared, the leading terms of which are proportional to $m_{\tilde{t}_1}^2$. The two kinds of contributions which appear with this mixed Yukawa coupling are either proportional to $\log(m_{\tilde{t}_1}^2 m_{\tilde{b}_1}^2/m_t^2)$, or they come without this logarithm and are directly proportional to the average stop mass instead [4]. We see their effect in Fig. 6 when we compare the $\tan\beta$ dependence of $m_A$ and $m_{H^\pm}$. For all value of $\tan\beta$ the requirement of two light Higgs scalar limits the pseudoscalar mass to a narrow corridor. We find hardly any SUSY scenarios with $m_A > 110$ GeV. By considering large values of $\tan\beta$, we can increase the charged Higgs mass by 50 GeV and avoid the top decay threshold.

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C. Heavy charged Higgs bosons and light stops

As described above, there is this remaining part of the light-scalars parameter space where charged Higgs searches in anomalous top decays and in the associated production with a top are challenging. It requires charged Higgs masses just above the top threshold and small $\tan\beta$. Note that we do not claim that LHC will not see a charged Higgs in this region of parameter space. Instead, we emphasize that covering this hole is a crucial task for the near future. We perform a dedicated scan around the one point which appears in Fig. 6 close to the hole in the LHC reach for
A choose fairly light stops and increase the trilinear coupling, These results point at another way of increasing the charged Higgs mass while keeping two light Higgs scalars: we Fig. 8 we show that this is in fact the case, with spin-independent elastic scattering cross sections as low as 10 \( 10^{-14} \) pb (see Fig. 5).

intermediate \( \tan \beta < 20 \) and \( m_{H^\pm} \sim 180 \) GeV. We vary each mass parameter up to 50\% above or below the value of the single point at \( \tan \beta \approx 8 \), \( m_A \approx 140 \) shown in Fig. 6. Only \( \tan \beta \) we vary over the entire range 1-60. At this specific point, \( M_2 \) and \( M_3 \) are both \( \sim 1 \) TeV, \( \mu \sim 2 \) TeV, \( A_t \sim -500 \) GeV, and the bino mass is somewhat light, \( M_1 \approx 140 \) GeV.

To see what kind of MSSM parameter choice can enhance the charged Higgs mass compared to the CP-even scalar masses, we extend the analytical approximation of Eq.(1) to mixing stops \[4\]. To avoid the LHC limits, we are now interested in parameter points with as small as possible \( \tan \beta \), so we can neglect the bottom Yukawa coupling. Instead, we take into account stop mixing. In two limits the expressions are particularly simple:

\[
m^2_{H^\pm} - (m^2_{h^+} + m^2_{h^-}) = m^2_W - m^2_Z + \frac{G_F m_t^2}{4\sqrt{2}\pi^2} \frac{A_t^2(A_t^2 + \mu^2)}{m^2_{t_1} m^2_{t_2}} \quad (\tan \beta \to 1)
\]
\[
m^2_{H^\pm} - (m^2_{h^+} + m^2_{h^-}) = m^2_W - m^2_Z + \frac{G_F m_t^2}{2\sqrt{2}\pi^2} \frac{(A_t - \mu)^2(A_t^2 + \mu^2)}{m^2_{t_1} m^2_{t_2}} \quad (\tan \beta \gg 1).
\]

These results point at another way of increasing the charged Higgs mass while keeping two light Higgs scalars: we choose fairly light stops and increase the trilinear coupling, \( A_t \), or the \( \mu \) parameter. We have checked explicitly that for large values of \( \tan \beta \) the same can be achieved through increasing \( A_b \). However, keeping both stops light and at the same time increasing one way or another the stop mixing parameters, \( A_t - \mu/\tan \beta \), means that at least the lighter stop and the lighter sbottom will be very light. This is precisely what we observe in Fig. 7, where we see that the lightest stop mass is always lighter than approximately 340 GeV.

Depending on the neutralino and chargino masses there are two possible stop decays we have to consider when we search for light stops at the Tevatron and at the LHC. If the stop is heavier than the lightest chargino, it will decay mostly into \( b\tilde{\chi}_1^+ \). If this decay channel is kinematically forbidden, as the LEP2 limits suggest for the Run I stop searches at the Tevatron, the stop has to decay to \( q\tilde{\chi}_1^0 \).

The production cross sections for light stops are large: at the Tevatron, a 200 GeV stop will be produced with a total cross section of approximately 300 fb, and at the LHC a 300 GeV stop will have a production rate around 10 pb. This means that at the LHC, although this kind of MSSM parameter point might be hard to find a charged Higgs boson, there will be huge numbers of top squarks establishing supersymmetry. Ongoing Tevatron searches with a projected Run II reach of well above 200 GeV (dependent on the neutralino and chargino masses) will already severely constrain models with light stops.

One might worry that in these models with lower \( \tan \beta \) direct dark matter detection might be more difficult. In Fig. 8 we show that this is in fact the case, with spin-independent elastic scattering cross sections as low as \( 10^{-10} \) pb. In each of these models, however, there is a rather light neutralino, \( m_{\tilde{\chi}_1^0} \lesssim 200 \) GeV, which makes all of these models observable by Super-CDMS and many of them observable by ZEPLIN-MAX, even for \( \sigma_{\chi N} \approx 10^{-10} \) pb (see Fig. 5).
Figure 8: The spin-independent neutralino-nucleon elastic scattering cross section is shown versus the CP-odd Higgs mass for the range of models described in section IV C of the text. For those models in which charged Higgs searches at the LHC will be difficult at the LHC \((m_A \gtrsim 140 \text{ GeV})\), this cross section is rather small, \(\sigma_{\chi N} \sim 10^{-9} - 10^{-10} \text{ pb}\). These models each contain a rather light neutralino, however, \(m_{\chi_0} \approx 70 - 200 \text{ GeV}\), making most of them accessible to ZEPLIN-MAX and all of them accessible to Super-CDMS (see Fig. 5). All models shown have \(95 \text{ GeV} < m_h < 101 \text{ GeV}\), \(111 \text{ GeV} < m_H < 119 \text{ GeV}\), and \(0.056 \lesssim \sin^2(\beta - \alpha) \lesssim 0.144\). The black points are consistent with measurements of \(B \to X_s \gamma\) at the 3\(\sigma\) level and do not violate the Tevatron constraint on the \(B_s \to \mu^+ \mu^-\) branching fraction. Blue points violate \(B_s \to \mu^+ \mu^-\), but are consistent with \(B \to X_s \gamma\). Green points violate \(B \to X_s \gamma\).

V. CONCLUSIONS

In this article, we describe the phenomenology of models within the MSSM containing CP-even Higgs bosons with masses of approximately 98 and 114 GeV, motivated by the excesses observed by LEP [2]. We study in detail indirect constraints on such light-Higgs models, coming from \(B \to X_s \gamma\), \((g - 2)_\mu\) and \(B_s \to \mu \mu\). The first two constraints require fairly large \(\tan \beta\) when compared with data, the latter tends to disfavor this regime, and the upcoming Tevatron Run II results are going to close in on our light-Higgs parameter space. However, it is possible to avoid all indirect constraints by tuning the different weak-scale SUSY breaking parameters.

We find that the kind models we are interested in typically include a light charged Higgs boson along with the other light Higgs bosons. Such a state is observable in anomalous top decays at the Tevatron or at the LHC. Scenarios with charged Higgs above the top threshold are automatically driven in the larger \(\tan \beta\) regime, where the charged Higgs can be detected in bottom–gluon fusion at the LHC. We find exceptionally models which fall between these two mass regions, so that neither of the two searches is optimized and a combination of the two search tactics would be needed. However, these challenging scenarios require one very light top squark \(m_{\tilde{t}_1} \lesssim 300 \text{ GeV}\), so a large fraction of them will already be ruled out at the Tevatron.

We also explore the phenomenology for neutralino dark matter in this class of models. The prospects for direct detection are excellent, with the majority of models within this scenario being testable in currently operating experiments, such as CDMS-II. We find that all models within this class will be testable in next generation direct dark matter detection experiments, such as ZEPLIN-MAX or Super-CDMS. The prospects for indirect are also favorable in this scenario.

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VI. APPENDIX: NEUTRALINO ANNIHILATION CROSS SECTIONS

The squared amplitudes for the processes, $\chi^0\chi^0 \to A \to f\bar{f}$ and $\chi^0\chi^0 \to H \to f\bar{f}$, averaged over the final state angle are given by [21]

$$\omega^A_{ff} = \frac{C_{fA}^2 C_{\chi^0\chi^0 A}^2}{(s-m_A^2)^2 + m_A^2 \Gamma_A^2} \frac{s^2}{16\pi} \sqrt{1 + \frac{4m_f^2}{s}}, \quad (7)$$

$$\omega^H_{ff} = \frac{C_{fH}^2 C_{\chi^0\chi^0 H}^2}{(s-m_H^2)^2 + m_H^2 \Gamma_H^2} \frac{(s-4m_0^2)(s-4m_f^2)}{16\pi} \sqrt{1 + \frac{4m_f^2}{s}}. \quad (8)$$

Here the label $H$ denotes either CP-even Higgs boson. $C_{fA}, C_{\chi^0\chi^0 A}, C_{fH}$ and $C_{\chi^0\chi^0 H}$ are the fermion-fermion-Higgs and neutralino-neutralino-Higgs couplings. $\Gamma_{A,H}$ are the widths of the respective Higgs bosons.

These squared amplitudes can be used to attain the thermally averaged annihilation cross section [22]

$$\langle \sigma v \rangle = \frac{\omega(s_0)}{m_{\chi^0}^2} \left[ \frac{\omega(s_0)}{m_{\chi^0}^2} - 2\omega'(s_0) \right] T + O(T^2)$$

$$= \frac{1}{m_{\chi^0}^2} \left[ 1 - 3\frac{T}{m_{\chi^0}} \right] \omega(s) \bigg|_{s \to 4m_{\chi^0}^2 + 6m_{\chi^0}T} + O(T^2),$$

(9)

where $T$ is the temperature. Keeping terms to zeroth and first order in $T$ is sufficient for the relic abundance calculation. Writing this as an expansion in $x = T/m_{\chi^0}$, $\langle \sigma v \rangle = a + bx + O(x^2)$, we arrive at

$$a_{\chi\chi \to A \to f\bar{f}} = \frac{g_4^2 c_f m_f^2 \tan^2 \beta}{8\pi m_W^2} \frac{m_{\chi^0}^2 \sqrt{1 - m_f^2/m_{\chi^0}^2}}{4m_{\chi^0}^2 - m_A^2)^2 + m_A^2 \Gamma_A^2}$$

$$\times \left[ \epsilon_U \epsilon_W - \epsilon_B \tan \theta_W \right] \sin \beta - \epsilon_D \left( \epsilon_W - \epsilon_B \tan \theta_W \right) \cos \beta \right]^2$$

$$b_{\chi\chi \to A \to f\bar{f}} = 0$$

$$a_{\chi\chi \to H \to f\bar{f}} = 0$$

$$b_{\chi\chi \to H \to f\bar{f}} = \frac{3g_4^2 c_f m_f^2 \sin^2 \alpha (m_{\chi^0}^2 - m_f^2)^2 \sqrt{1 - m_f^2/m_{\chi^0}^2}}{16m_W^2 \sin^2 \beta (4m_{\chi^0}^2 - m_H^2 + im_H \Gamma_H^2)^2}$$

$$\times \left[ \epsilon_U \epsilon_W - \epsilon_B \tan \theta_W \right] \cos \alpha - \epsilon_D \left( \epsilon_W - \epsilon_B \tan \theta_W \right) \sin \alpha \right]^2$$

$$a_{\chi\chi \to H \to f\bar{f}} = 0$$

$$b_{\chi\chi \to H \to f\bar{f}} = \frac{3g_4^2 c_f m_f^2 \cos^2 \alpha (m_{\chi^0}^2 - m_f^2)^2 \sqrt{1 - m_f^2/m_{\chi^0}^2}}{16m_W^2 \sin^2 \beta (4m_{\chi^0}^2 - m_H^2 + im_H \Gamma_H^2)^2}$$

$$\times \left[ \epsilon_U \epsilon_W - \epsilon_B \tan \theta_W \right] \sin \alpha - \epsilon_D \left( \epsilon_W - \epsilon_B \tan \theta_W \right) \cos \alpha \right]^2,$$

(10)

where $c_f$ is a color factor, equal to 3 for quarks and 1 for leptons. The quantities $\epsilon_{B,W,U,D}$ are the components of the lightest neutralino which are bino, wino, up-type higgsino and down-type higgsino, respectively. These expressions assume that the final state fermions are down-type. For up-type fermions the factor of $\tan^2 \beta$ should be replaced by $\cot^2 \beta$.

[1] R. Barate et al. [ALEPH Collaboration], Phys. Lett. B 565, 61 (2003); A. Heister et al. [ALEPH Collaboration], Phys. Lett. B 526, 191 (2002).
[2] M. Drees, arXiv:hep-ph/0502075; for a similar CP violating explanation see M. Carena, J. R. Ellis, A. Pilaftsis and C. E. M. Wagner, Phys. Lett. B 495, 155 (2000).
[3] see e.g.: J. F. Gunion and A. Turski, Phys. Rev. D 39, 2701 (1989) and Phys. Rev. D 40, 2333 (1989); M. S. Berger, Phys. Rev. D 41, 225 (1990); R. Hempfling and A. H. Hoang, Phys. Lett. B 331, 99 (1994); we are using the formulas as in: A. Krause, T. Plehn, M. Spira and P. M. Zerwas, Nucl. Phys. B 519, 85 (1998).
[4] M. Carena, J. R. Espinosa, M. Quiros and C. E. M. Wagner, Phys. Lett. B 355, 209 (1995).
