A Supersymmetric Explanation of the Excess of Higgs–Like Events at the LHC and at LEP

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Abstract

The LHC collaborations have recently announced evidence for the production of a “Higgs–like” boson with mass near 125 GeV. The properties of the new particle are consistent (within still quite large uncertainties) with those of the Higgs boson predicted in the Standard Model (SM). This discovery comes nearly ten years after a combined analysis of the four LEP experiments showed a mild excess of Higgs–like events with a mass near 98 GeV. I show that both groups of events can be explained simultaneously in the minimal supersymmetric extension of the SM, in terms of the production and decay of the two neutral CP–even Higgs bosons predicted by this model, and explore the phenomenological consequences of this explanation.
1 Introduction

Recently the LHC collaborations ATLAS and CMS announced the discovery of a “Higgs–like” boson with mass near 125 GeV [1]. This new boson has been detected at the LHC chiefly in the $\gamma\gamma$ and four lepton final states. In addition, there is evidence, at the $\sim 3\sigma$ level, for decays into $b\bar{b}$ pairs from the Tevatron experiments CDF and D0 [2].

Within the still quite large experimental (and theoretical [3]) uncertainties, the properties of the new boson are consistent with those of the single physical Higgs boson in the Standard Model (SM). However, as well known, the scalar sector of the SM is technically unnatural, since it suffers from quadratic divergencies. These divergencies are canceled in supersymmetric extensions of the SM [4].

Even the simplest such theory, the Minimal Supersymmetric extension of the SM (MSSM), contains two Higgs doublets, the second doublet being required both for the cancellations of anomalies (from the higgsinos), and in order to give masses to all quarks [5]. As a result, the MSSM contains three neutral physical Higgs states. In the absence of CP violation, these can be classified as two CP–even states $h, H$ (with $m_h < m_H$) and one CP–odd state $A$. Most interpretations of the new boson discovered by the LHC experiments (many of which were published after the first experimental evidence was announced in December 2011) within the MSSM focus on the possibility that it is the lighter CP–even state $h$ [6, 7, 8]. However, achieving $m_h \simeq 125$ GeV is only possible if stop squarks are very heavy. By most definitions, this requires a somewhat uncomfortable amount of finetuning.

In this scenario the heavier neutral Higgs bosons $H, A$ as well as the charged Higgs bosons $H^\pm$ can be essentially arbitrarily heavy. In fact, in simple (constrained) scenarios of supersymmetry breaking, the large lower bounds on (first generation) squark masses typically require these states to be quite heavy; within such constrained scenarios the new state therefore has to be interpreted as $h$.

The possibility that instead the heavier state $H$ has been discovered has also been entertained [10, 7, 11]. In this case the lighter CP–even state would obviously have to be lighter than 125 GeV, and would need to satisfy limits from Higgs searches both at the LHC [1] and at LEP [12]. As pointed out in refs. [7, 11] this is not difficult to achieve, if $H$ is SM–like in agreement with experimental observations of the new state at 125 GeV. In particular, sufficiently large branching ratios for $H$ into four leptons, and into two photons, require the couplings of $H$ to two massive gauge bosons to not differ very much from the corresponding SM values. In the context of the MSSM this automatically implies that $h$ has suppressed couplings to $W$ and $Z$, making it difficult to detect.

Here I wish to point out that in this scenario $h$ might be put to good use, by explaining an excess of Higgs–like events observed some ten years ago by the four LEP collaborations [12]. Actually, the combined LEP data showed two regions of reconstructed Higgs mass where some excess occurred. One was right at the kinematic limit, near 115 GeV. This excess was observed mostly by the ALEPH collaboration [14]; in the combined data, its statistical significance reached only 1.7 standard deviations. It is compatible with an SM–like Higgs with this mass. However, this interpretation is at odds with the interpretation of the new particle discovered at the LHC as an SM–like Higgs boson.

The combination of the data from all four LEP experiments also revealed [12] a somewhat more significant excess near 98 GeV, with significance of about 2.3 standard deviations. This excess is not compatible with an SM Higgs at that mass; rather, it’s compatible with an about ten times smaller
production cross section than predicted by the SM for this mass range. In the context of the MSSM this can easily be arranged by reducing the $hZZ$ (and hence also the $hWW$) coupling [15]. This implies that the other MSSM Higgs bosons have to be relatively light: if they were heavy, $h$ would become SM–like. A detailed analysis found an upper bound on $m_H$ of about 140 GeV [15]. The new particle at 125 GeV therefore falls right in the middle of the allowed range for $m_H$ in this scenario, where the lower bound (of about 114 GeV) comes from LEP Higgs searches.

The LHC discovery obviously greatly constrains the allowed parameter space of this scenario, where now the masses of both CP–odd Higgs bosons are fixed within a few GeV theoretical and experimental uncertainty. However, as well known the MSSM Higgs sector is subject to large radiative corrections. This introduces several new parameters, the most important ones being the ones appearing in the stop mass matrix. Here I present a detailed analysis of this scenario. This not only updates ref. [15] by including the constraint $m_H \simeq 125$ GeV; I also carefully compute the relevant decay widths, and resulting branching ratios and signal strengths, of the neutral Higgs bosons, where the latter are normalized to the signal strength of the SM Higgs boson. I also include constraints from null searches for neutral MSSM Higgs bosons decaying into tau pairs performed by CMS [16], and for charged Higgs bosons produced in top quark decays performed by ATLAS [17]. The former are more important, considerably limiting the allowed parameter space of this scenario. I nevertheless find that $m_A$ masses roughly between 95 and 150 GeV are allowed in this scenario. The charged Higgs boson mass can reach up to about 170 GeV. The signals for $H$ production in both the di–photon and four lepton channels can be considerably enhanced, but the ratio of these two signals cannot exceed its SM prediction by more than about 35%.

This analysis has been performed in the framework of the general MSSM, where all relevant parameters are fixed directly at the weak (or TeV) scale. In order to limit the size of the parameter space, I will not specify soft breaking parameters for the first two generations of sfermions, which have almost no impact on the masses and couplings of Higgs bosons. This approach also permits me to ignore all constraints from flavor physics, which depend very strongly on the flavor structure of the soft breaking terms.

The rest of this paper is organized as follows. In Sec. 2 I describe details of the analysis. In Sec. 3 I explore the parameter space that is compatible with this explanation; in particular, I give allowed ranges for physical quantities of interest, and explore correlations between them. Finally, Sec. 4 contains a brief summary and conclusions.

## 2 Details of the Analysis

This analysis is performed in the framework of the general MSSM, where all relevant parameters are fixed at the weak scale, and no high–scale constraints on the spectrum of superpartners are imposed.

Obviously at least the leading radiative corrections [18] to the masses and mixing angle of the MSSM Higgs bosons have to be included in any quantitative analysis. This is most easily done using the effective potential (or, equivalently, Feynman diagrammatic calculations with vanishing external momentum). Recall that the entire Higgs spectrum should be relatively light in this scenario; this

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Ref. [10] also considered the MSSM with $m_H \simeq 125$ GeV, including scenarios with $m_h \leq 110$ GeV and suppressed $ZZh$ couplings, in the context of an analysis of scenarios with a light neutralino as Dark Matter candidate, based on very early, preliminary LHC results. No bounds on the strengths of the $H$ signals were imposed, and $h$ was not required to explain the LEP excess near 98 GeV.
should increase the reliability of the effective potential method.

In order to allow an efficient sampling of the parameter space, I only include corrections from the top–stop and bottom–sbottom sectors to the Higgs boson mass matrices, which give the by far most important contributions. The relevant expressions have been taken from refs.\[19\] (for the pure Yukawa corrections to both the neutral and charged Higgs boson mass matrices); \[20\] (for the mixed electroweak–Yukawa corrections to the neutral Higgs boson mass matrix); and \[19\] \[21\] and \[21\] \[22\] (for the inclusion of leading higher order QCD and top Yukawa corrections, respectively, by using running quark masses defined at the appropriate scale in the one–loop effective potential). The leading SUSY QCD corrections are included through the gluino–stop and gluino–sbottom corrections to the top and bottom mass, respectively \[23\]; as shown in ref.\[24\], this reproduces the full SUSY QCD correction very accurately. The calculation performed here should reproduce the neutral MSSM Higgs masses with an error of about two or three GeV \[25\] \[24\]. This theoretical uncertainty will be included in the constraints imposed on \(m_h\) and \(m_H\). Note that the SUSY QCD corrections to the bottom mass are also included in the calculation of the corresponding Yukawa couplings, which affect both the partial widths of the neutral Higgs bosons into \(b\bar{b}\) pairs and the \(b\) loop contribution to the partial widths of the decays into gluon and photon pairs.

The running top mass, \(m_t(m_t)\) is fixed to 165 GeV (in the \(\overline{\text{DR}}\) scheme). This corresponds to a pole mass near 173 GeV, the current central value \[26\]. I also fix \(m_b(m_b) = 4.25\) GeV. As final simplification, I have taken the soft breaking parameters in the stop and sbottom mass matrices to be the same. This is always true for the masses of the superpartners of the left–handed squarks, due to \(SU(2)\) invariance, but the masses of the \(SU(2)\) singlet squarks as well as the two \(A\)–parameters could in principle be different. However, we will see that the CMS di–tau search requires the ratio of vacuum expectation values \(\tan \beta\) to be relatively small, below 13; as a result, sbottom loops are always subdominant, and thus need not be treated as carefully as stop loops.

The most convincing signals for the new state at 125 GeV have been found in the di–photon channel. This is obviously only accessible through loop diagrams. In addition to the diagrams involving \(W\) bosons or third generation fermions, diagrams involving charged Higgs bosons as well as all third generation sfermions are included. It had been noticed \[27\] \[11\] that loops involving stau sleptons could significantly change the di–photon widths of neutral CP–even MSSM Higgs bosons. I therefore allow the \(\tilde{\tau}_{L,R}\) soft breaking masses as well as the trilinear soft breaking parameter \(A_\tau\) to vary independently from the parameters of the stop sector. However, it turns out that in the given scenario, stau loops always make very small contributions. Similarly, the partial widths of the Higgs bosons into gluons are computed including loops of third generation quarks as well as squarks (squark loops are absent in case of the CP–odd Higgs boson). The relevant expressions are taken from \[28\].

Altogether we are thus left with ten free parameters: \(\tan \beta\), \(m_A\), \(\mu\), \(m_{\tilde{t}_L}\), \(m_{\tilde{t}_R}\), \(A_t\), \(m_{\tilde{\tau}_L}\), \(m_{\tilde{\tau}_R}\), \(A_\tau\), \(m_{\tilde{g}}\). This ten–dimensional parameter space has been scanned randomly, subject to the following con-
straints not involving Higgs bosons:

\[ |\mu|, m_{\tilde{t}_R}, m_{\tilde{t}_L}, m_{\tilde{b}_R}, m_{\tilde{b}_L} \leq 5 \text{ TeV}; \]  
\[ |\mu|, m_{\tilde{t}_1}, m_{\tilde{b}_1} \geq 100 \text{ GeV}; \]  
\[ |m_{\tilde{t}_1} - m_{\tilde{b}_1}| \leq 50 \text{ GeV or max}(m_{\tilde{t}_1}, m_{\tilde{b}_1}) > 300 \text{ GeV}; \]  
\[ m_{\tilde{g}} \geq 600 \text{ GeV}; \]  
\[ |A_t|, |\mu| \leq 1.5 (m_{\tilde{t}_R} + m_{\tilde{t}_L}); \]  
\[ |A_{\tau}|, |\mu| \leq 1.5 (m_{\tilde{\tau}_R} + m_{\tilde{\tau}_L}); \]  
\[ \delta \rho_{\tilde{t}\tilde{b}} \leq 2 \cdot 10^{-3}. \] 

The first of these constraints is a (quite conservative) naturalness criterion. Conditions (1b) ensure that higgsino–like charginos (with mass \(|\mu|\)) as well as the lighter physical stop (\(\tilde{t}_1\)), sbottom (\(\tilde{b}_1\)) and stau (\(\tilde{\tau}_1\)) states escaped detection at LEP [29]. Condition (1c) ensures that only one of the two lighter squark states can be below 300 GeV, unless they are close in mass. In the latter case they could both be close in mass to the lightest neutralino, in which case \(\tilde{t}_1\) and \(\tilde{b}_1\) pair production would lead to events with a small amount of visible energy, which are difficult to detect. Condition (1d) is a rather conservative interpretation of gluino search limits in the general MSSM. Note that loops involving gluinos affect the Higgs masses and mixing angle only at two–loop order, but modify the \(hb\bar{b}\) and \(Hb\bar{b}\) couplings already at one–loop. The upper bounds (1e,f) on the parameters determining mixing in the stop and stau sectors have been imposed to avoid situations where \(\tilde{t}\) or \(\tilde{\tau}\) fields have non–vanishing VEVs in the absolute minimum of the scalar potential [30]. Finally, (1g) requires the contribution of stop–sbottom loops to the electroweak \(\rho\) parameter [32] to be sufficiently small.

In order to be able to describe the (mild) excess of Higgs–like events at LEP, and the properties of the new boson discovered at the LHC, the Higgs sector has to simultaneously satisfy the following constraints:

\[ 95 \text{ GeV} \leq m_h \leq 101 \text{ GeV}; \]  
\[ 123 \text{ GeV} \leq m_H \leq 128 \text{ GeV}; \]  
\[ 0.056 \leq \sin^2(\alpha - \beta) \leq 0.144; \]  
\[ 0.5 \leq R_{VV}^H \leq 2.0 \quad (V = W, Z); \]  
\[ 0.5 \leq R_{\gamma\gamma}^H. \] 

The first of these constraints places \(m_h\) in the range where an excess of events had been observed at LEP [12]. Similarly, (2b) ensures that \(m_H\) agrees with the value reported by the LHC experiments [1]. In both cases, the range is a crude estimate of theoretical and experimental uncertainties. Note that the peak at the LHC is somewhat narrower than at at LEP, since the latter has been observed chiefly in multi–hadron final states.

The third constraint [15] ensures that the \(Zh\) production cross section at LEP is roughly ten times smaller than the corresponding cross section in the SM, for given mass of the Higgs boson; as noted above, the excess at LEP is compatible with Higgs production only if the \(ZZh\) coupling is suppressed.\(^5\)

\(^5\)Note that ref. [27], where the influence of light staus on the two–photon widths of MSSM Higgs bosons was first explored, only imposes the much weaker constraint on \(|\mu|\) that follows from the requirement that the zero temperature tunnel rate into the false vacuum [31] is sufficiently small.
The last two conditions ensure that the LHC signals in the $WW^*$, four lepton and di-photon channels come out roughly correct. They are described by the quantities

$$R_{H}^{XX} \equiv \frac{\Gamma(H \to gg)}{\Gamma(H_{SM} \to gg)} \cdot \frac{\Gamma(H \to XX)}{\Gamma(H_{SM} \to XX)} \cdot \frac{\Gamma(H_{SM, tot})}{\Gamma(H_{tot})}. \quad (3)$$

They describe the strength of the $H$ signal in the $XX$ channel normalized to the strength of the corresponding signal for the SM Higgs boson $H_{SM}$. Here I have assumed that Higgs production at the LHC is dominated by gluon fusion; this is true both in the SM and in the relevant parameter range of the MSSM. The strength of the four lepton signal observed at the LHC, which in the Higgs interpretation of the signal is due to the decay of the Higgs boson into a real and a virtual $Z$ boson, agrees quite well with the SM prediction; I therefore allow this signal to be at most a factor of two stronger or weaker than in the SM. In contrast, the di–photon signal appears somewhat stronger than in the SM; I therefore only impose a lower bound on the strength of the signal in this channel. Since the $\gamma\gamma$ invariant mass peak has a finite width of roughly 1 GeV, essentially given by the experimental resolution, $R_{H}^{\gamma\gamma}$ includes the contribution from $gg \to A \to \gamma\gamma$ whenever $|m_H - m_A| < 1$ GeV. However, in the parameter region of interest this contribution is always very small, due to the small branching ratio of $A \to \gamma\gamma$ decays, which in turn is due to the absence of an $AW^+W^-$ coupling.

Finally, null results of additional searches for Higgs bosons have to be imposed. In particular, for charged Higgs bosons with mass (well) below $m_t - m_b$, ATLAS searches for $t \to H^+ b$ decays, with $H^+ \to \tau^+ \nu_{\tau}$, exclude [17] both small and large values of tan $\beta$, leaving an allowed strip centered at $\tan \beta \simeq \sqrt{m_t(m_t)/m_b(m_t)} \approx 7$ where the $H^+tb$ coupling is minimal. At least in the present context, the CMS search for neutral MSSM Higgs bosons in the di–tau channel [16] is even more constraining. Here I have taken both analyses at face value. Since the ATLAS charged Higgs search is basically independent of the details of the neutral Higgs spectrum, it should indeed apply to the present scenario. CMS states its bounds on MSSM parameter space in the context of the “maximal mixing” scenario, which maximizes $m_h$ for given average stop mass. In this scenario the CP–odd state is typically quite closely degenerate with either $h$ or $H$, especially for large tan $\beta$ where this search is most sensitive. Such a degeneracy obviously increases the yield of tau pairs of a given invariant mass. In the present context the mass splittings between all three neutral Higgs bosons are often sizable; this should lead to somewhat smaller signals in the di–tau channel than in the “maximal mixing” scenario. Incorporating the CMS constraints in the $(m_A, \tan \beta)$ plane without modification therefore probably overstates their impact somewhat. However, this is not easy to quantify without a full simulation including experimental resolutions.

This concludes the description of the analysis. Let us now turn to the results.

3 Results

This Section contains a discussion of the results of the scan of parameter space, subject to the constraints discussed in the previous Section. Of course, the first and quite nontrivial result is that allowed sets of parameter sets can indeed be found, i.e. the (phenomenological) MSSM can indeed explain at the same time the (mild) excess of Higgs–like events at LEP and the detection of a Higgs–like particle by the LHC experiments.
In order to further test this scenario, one has to know what it implies for the relevant observables. To that end, I will first describe upper and/or lower bounds on quantities of interest that were found in the scan, before discussing correlations between pairs of these quantities.

### 3.1 Bounds on Observables

Let us first look at observables in the Higgs sector. Note first of all that the upper and lower limits on both the $h$ and $H$ mass can be saturated, i.e. the scenario doesn’t allow to further shrink either of these mass regions beyond the limits imposed as constraints in eqs.(2a,b).

However, not surprisingly there are nontrivial bounds on the masses of the CP–odd and charged Higgs bosons. Some bounds already follow [15] from the constraint (2c) on the $Zhh$ coupling: if $m_A$ or $m_{H^+}$ becomes very large, $h$ automatically becomes SM–like; in this “decoupling scenario” the upper bound on the $Zhh$ coupling is therefore badly violated. At the same time the constraint $m_H > 123$ GeV imposes a non–trivial lower bound on the mass of the charged Higgs. Altogether I find

$$120 \text{ GeV} \leq m_{H^+} \leq 170 \text{ GeV}.$$  \hfill (4)

The upper bound can be saturated, implying that $t \rightarrow H^+b$ decays can be closed kinematically. This is in (mild) conflict with a statement of [11], probably due to the large range of parameters I explored here. Saturating this upper bound requires very large $\mu$, a large hierarchy between the $\tilde{t}_L$ and $\tilde{t}_R$ masses, a top mixing parameter $|A_t|$ saturating its upper bound (with $A_t < 0$, $\mu > 0$), and moderate $\tan \beta \simeq 6$.

The corresponding allowed range for the mass of the CP–odd Higgs boson $A$ reads

$$96 \text{ GeV} \leq m_A \leq 152 \text{ GeV}.$$  \hfill (5)

The upper bound on $m_A$ is saturated for the same choice of parameters as the upper bound on $m_{H^+}$. The lower bound on $m_A$ together with the constraint (2b) on $m_h$ implies that limits from searches for $hA$ production at LEP are always satisfied.

The LHC searches for non–SM Higgs bosons discussed at the end of Sec. 2 considerably restrict the allowed values of $\tan \beta$, leading to

$$5.5 \leq \tan \beta \leq 12.5.$$  \hfill (6)

The lower bound is largely determined by the ATLAS search for charged Higgs bosons, while the upper limit is chiefly due to the CMS search for neutral Higgs bosons in the di–tau channel. The allowed range of $\tan \beta$ thus looks quite narrow. However, closing it entirely may not be easy. As noted above, the $H^+tb$ coupling reaches its minimum near $\tan \beta = \sqrt{m_t/m_b}$, which falls in the range $[6]$. The signal strength in the di–tau channel scales essentially like $\tan^2 \beta$, so reducing the upper bound on $\tan \beta$ by a factor of about 2.5 requires an increase of the sensitivity of the search by a factor of six. Recall also that my interpretation of the CMS bound might be overly strict, i.e. the true bound might be somewhat weaker.

As noted earlier, the constraint (2c) implies that the $HWW$ and $HZZ$ couplings have close to SM strength. However, this doesn’t imply that the $gg \rightarrow H \rightarrow ZZ^* \rightarrow 4\ell$ signal also has close to SM strength. On the one hand, loops of new strongly interacting sparticles, in particular stops, can change the $H$ production cross section significantly. On the other hand, the couplings of $H$ to SM fermions, in particular to $b$ quarks and $\tau$ leptons, can still differ considerably from their SM values,
thereby modifying the $H$ decay branching ratios. As a result, both the upper and the lower limits on $R^{ZZ}_H$ in (2d) can be saturated, if $\tilde{t}_1$ is not too heavy. The lower bound on $R^{\gamma\gamma}_H$ can also be saturated, and the upper bound is
\[ R^{\gamma\gamma}_H \leq 2.2. \tag{7} \]

A significant enhancement of the $H \to \gamma\gamma$ signal, which is hinted at by present data, is thus possible in this scenario. However, this enhancement is mostly due to the increase of the $H$ production cross section and/or decrease of its total width; both these effects also increase $R^{ZZ}_H$. In fact, when considering the ratio of signal strengths in the $\gamma\gamma$ and $4\ell$ (or, more generally, $VV^*$) channels normalized to their respective SM values, only a moderate deviation from unity is possible in this scenario:
\[ 0.66 \leq \frac{R^{\gamma\gamma}_H}{R^{ZZ}_H} \leq 1.3. \tag{8} \]

I do not find any scenarios where the branching ratio for $H \to \gamma\gamma$ decays is affected significantly by $\tilde{\tau}$ loops; this is probably due to the vacuum stability constraint $|\mu| \leq 3(m_{\tilde{\tau}_L} + m_{\tilde{\tau}_R})/2$, which has not been imposed in refs.\[27\] and \[11\]. The contribution of charged Higgs loops to this branching ratio is also always very small, although the charged Higgs boson is quite light in this scenario, as shown in \[4\].

The di–tau channel is currently poorly constrained by the data. In fact, the strength of the signal in this channel can deviate quite significantly from its SM value in the present scenario:
\[ 0.2 \leq R^{\tau\tau}_H \leq 5.7. \tag{9} \]

Note that the size of the $\tau$ Yukawa coupling exceeds its SM value for $\tan \beta > 1$. However, the size of the $H\tau^+\tau^-$ coupling also depends on the mixing angle $\alpha$ between the neutral CP–even Higgs bosons, and even vanishes if $\cos \alpha = 0$. This limit cannot be realized in the present scenario, but a substantial suppression of the di–tau signal strength is possible. On the other hand, the maximal enhancement of this signal occurs when stop and sbottom loops simultaneously enhance the $gg \to H$ production cross section and suppress the $Hb\bar{b}$ coupling, while the $H\tau^+\tau^-$ coupling is enhanced since $|\cos \alpha| > |\cos \beta|$. Moreover, if $|m_H - m_A| \leq 4$ GeV, the contribution $R^{\tau\tau}_A$ has been added to $R^{\tau\tau}_H$, since the $H$ and $A$ di–tau signals would be difficult to distinguish experimentally in this case. However, the $A \to \tau^+\tau^-$ signal is quite small in the relevant region of parameter space, which has $\tan \beta \approx 6$; here the $\tan \beta$ enhancement of the $Ab\bar{b}$ coupling cannot yet compensate for the cot $\beta$ suppression of the $At\bar{t}$ coupling, leading to an $A$ production cross section from gluon fusion which is significantly smaller than the corresponding SM value. Nevertheless the upper end of the range \[9\] is probably already disfavored by present data, given the absence of a clear signal in this channel. Note, however, that requiring $R^{\tau\tau}_H < 3$ does not appreciably alter the allowed ranges of most other observables.

The di–tau channel also seems to offer the best chance for detecting the light CP–even scalar $h$ at the LHC in this scenario. The $\gamma\gamma$ signal is very weak for this state, $R^{\gamma\gamma}_h \leq 0.035$; this is due to the

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*Such ratios have very recently also been discussed in \[8\], which however assumes that the LHC signals are due to the production of $h$, not $H$.  

†In general there is also a significant contribution to the inclusive $A$ production cross section from the tree–level process $gg \to b\bar{b}A$; however, the presence of two additional $b$–jets should allow to discriminate this process from SM Higgs production, so this contribution should not simply be added to $R^{\tau\tau}_H$ even if $m_A = m_H$.  

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reduction of the $hW^+W^-$ coupling implied by the constraint (2c). On the other hand, the $gg \to h$ production cross section need not be suppressed relative to its SM value, so that

$$0.12 \leq R_{h\tau\tau} \leq 3.4.$$  \hfill (10)

Even the upper end of this range might be difficult to probe at the LHC, since $h$ is quite close in mass to the $Z$ boson which yields a much stronger signal in the $\tau^+\tau^-$ channel. For this reason, the lower end will probably remain unobservable even for LHC upgrades.

Before concluding this Subsection, let me briefly mention some constraints on quantities related to the stop sector; these are the only MSSM parameters not directly related to the Higgs sector for which some non–trivial constraints can be derived in the present context. For example, requiring that the heavy CP–even Higgs boson $H$ explains the LHC signals makes it quite difficult to find acceptable scenarios with small values of the Higgs(ino) mass parameter $\mu$: only for $\tan \beta \gtrsim 9$ do some solutions with $|\mu| < 400$ GeV survive; for $\tan \beta = 10$, some solutions with $\mu > 0$ saturate the lower bound $|\mu| = 100$ GeV. Similarly, $|A_t|$ has to exceed 400 GeV, and the sum $|A_t| + |\mu| > 2$ TeV.

Turning to sfermion masses, the $\tilde{t}_1$, $\tilde{b}_1$ and $\tilde{\tau}_1$ masses can all saturate the lower bounds of 100 GeV; moreover, no meaningful upper bounds on these masses can be derived. On the other hand, the $\tilde{t}_2$ mass must exceed 600 GeV in this scenario, and the sum of $\tilde{t}_1$ and $\tilde{t}_2$ masses must exceed 900 GeV. For comparison: demanding $m_h > 123$ GeV, as required if the recent LHC discovery is to be interpreted in terms of the production and decay of $h$, leads to the lower bound $m_{\tilde{t}_1} + m_{\tilde{t}_2} > 950$ GeV. This indicates that the amount of finetuning required in these two MSSM explanations of the LHC signal is comparable.

3.2 Correlations between Observables

Let us now analyze correlations between observables that result from the constraints (2) as well as the upper limits on MSSM Higgs searches at the LHC, beginning with correlations between physical masses. The most obvious such correlations are shown in Figs. 1a,b, which show the correlation between the mass of the CP–odd Higgs boson and the mass of the charged Higgs boson and lighter stop eigenstate, respectively. These, and all following, scatter plots are based on scans over parameter space containing several million sets of parameters (not all of which are plotted), with special emphasis on those regions of parameter space where an observable reaches an extremum.

As shown in Fig. 1a, the masses of the charged and CP–odd Higgs bosons are strongly correlated. This is not surprising, given the tree–level relation $m_{H^+} = \sqrt{m_A^2 + M_W^2}$, which is indicated by the solid red line. We see that the radiative corrections to this relation are usually negative, but rather modest in size in the allowed region of parameter space. This is a consequence of the upper bounds (1e,f) on the parameters determining stop mixing, which also largely determine the size of trilinear couplings between Higgs bosons and stop and sbottom squarks; the upper bound (6) on $\tan \beta$ also plays a role in limiting the size of the corrections to this relation.

The right frame in Fig. 1 shows that an upper bound on the $\tilde{t}_1$ mass results in the present scenario if $m_A \gtrsim 110$ GeV; this results from the upper bound on $m_H$. Notice that for relatively light $\tilde{t}_1$, the heavier CP–even Higgs boson can be significantly lighter than the CP–odd Higgs boson; in contrast, at the tree level one has $m_H > m_A$. The magnitude of these negative corrections to $m_H$ is limited by the upper bound on $|A_t|$ and $|\mu|$ given in (1e).

Figure 2 shows correlations between a mass and a (ratio of) signal strength(s). Recall that only the gluon fusion contribution to various Higgs signals is included here; this should be the dominant
channel in general, but, depending on the cuts, there might be significant contributions also from $WW$ and $ZZ$ fusion. Associate production with a $b\bar{b}$ pair, which can become quite important at large $\tan\beta$ \cite{28}, is not expected to be very important, given the upper bound \cite{6}.

Frame (a) shows correlations between the mass of the CP–odd Higgs boson and the $h$ signal strength in the $\tau^+\tau^-$ channel. The $h \rightarrow \tau^+\tau^-$ signal strength shows a first peak at $m_A = 100$ GeV $\simeq m_h$; due to this near–degeneracy, the signal from $A \rightarrow \tau^+\tau^-$ has been added, which increases $R_{\tau\tau}^h$ by up to one unit. This signal reaches its (local) maximum for the largest allowed value of $\tan\beta$. Here the cross sections for producing an $h$ or $A$ boson are slightly larger than the corresponding cross section for producing an SM Higgs boson with equal mass; the enhancement of the bottom Yukawa coupling over–compensates the suppression of the couplings of these two lighter Higgs bosons to top quarks. Both stop squarks need to be fairly heavy in this region of parameter space, in order to obtain a sufficiently large value of $m_H$.

Recall that the $A$ and $h$ contributions to this channel are added only for $|m_A - m_h| \leq 4$ GeV. For intermediate values of $m_A$ the maximal strength of this signal therefore decreases, before reaching its absolute maximum near $m_A = 130$ GeV. At the absolute maximum of $R_{\tau\tau}^h$, the lower bound on $m_{\tilde{t}_1}$ is saturated, and sbottom loops suppress the $h\bar{b}b$ coupling. It is still significantly larger than the corresponding coupling in the SM, but the enhancement of the $h\tau^+\tau^-$ coupling is even larger, leading to an enhanced branching ratio for $h \rightarrow \tau^+\tau^-$. Moreover, the $h \rightarrow gg$ width, and hence the $h$ production cross section, is dominated by $\tilde{t}_1$ and $b$ loops, which have the same sign, while the subleading $t$ loop contribution has opposite sign. As a result, the gluonic decay width of $h$ exceeds its SM value by about a factor of two. In combination, this enhances the $h \rightarrow \tau^+\tau^-$ signal by up to a factor of 3.3 over its SM value. However, given that $m_h$ is quite close to $M_Z$, it is not clear whether this enhancement is sufficient to make $h$ detectable at the LHC in this channel. Note also that values of $R_{\tau\tau}^h$ well below 1 are possible for nearly all values of $m_A$. This is chiefly due to the suppression of the $h \rightarrow gg$ width, which in turn is caused by strong cancellations between the $t$ and $b$ loop contributions in this region of parameter space. There is a nontrivial, although phenomenologically probably not

Figure 1: Allowed region in the $(m_A, m_{H^\pm})$ (a) and the $(m_A, m_{\tilde{t}_1})$ plane (b), after the constraints \cite{2} as well as the various sparticle and Higgs search limits discussed in the text have been imposed. The solid (red) line in (a) shows the tree–level relation between $m_{H^\pm}$ and $m_A$. 


very interesting, lower bound on this quantity, as shown in (10), since the $b$ loop contribution has a sizable imaginary part, which cannot be canceled by loops involving much heavier $t$ or $\tilde{t}_1$ particles.

Fig. 2 shows that obtaining $|\mu| < 0.5$ TeV is quite difficult in this scenario. It is possible only if $\tan \beta$ is near the upper end of the allowed range [6]. Moreover, a large mass splitting between the $\tilde{t}_L$ and $\tilde{t}_R$ masses is required. The masses of the CP–even Higgs bosons $h$ and $H$ are then near the lower and upper ends of their allowed ranges, respectively. In this region of parameter space both the $\tilde{t}_1$ and the real part of the $b$ loop contributions to $H \to gg$ as well as $H \to \gamma\gamma$ have the same sign as the $t$ loop contributions. This enhances the partial width for $H \to gg$ by up to a factor of 1.6, but suppresses the partial width for $H \to \gamma\gamma$, which is dominated by $W$ loops, by up to 20%. In addition, the partial widths for $H \to b\bar{b}$ and $H \to \tau^+\tau^-$ are enhanced, further suppressing the branching ratio for $H \to \gamma\gamma$. This overcompensates the increase of the $H$ production cross section.

On the other hand, $R_H^{\gamma\gamma}$ can exceed unity for $|\mu| > 1$ TeV, and reaches its absolute upper bound near $\mu = 2$ TeV. Here both $\tilde{t}_1$ and $\tilde{b}_1$ are quite light, while $m_{\tilde{t}_2} \simeq 1$ TeV, and $\tan \beta \simeq 6$. The light $\tilde{t}_1$
increases the partial width for $H \to \gamma\gamma$ by about 5%, but suppresses the partial width for $H \to gg$ by about 30%. This suppression of the total cross section for $H$ production is over–compensated by the greatly reduced partial widths for $H \to b\bar{b}$ and, to a lesser extent, $H \to \tau^+\tau^-$ decays; the light $b_1$ significantly reduces the $Hb\bar{b}$ coupling vial SUSY QCD loop corrections in this case. The signals in the $VV^*$ channels ($V = W^\pm$ or $Z$) are therefore also enhanced by almost a factor of two.

Note finally that Fig. (2) shows more solutions with $\mu > 0$ than with $\mu < 0$; moreover, the LEP lower bound on $|\mu|$ can only be saturated for positive $\mu$. The reason for this asymmetry is that only positive values of the gluino mass parameter were considered. The relative sign (more generally, relative phase) between $\mu$ and the gluino mass parameter has physical meaning, just as the relative signs (or phases) between $\mu$ and the soft breaking $A$–parameters are significant. Since only these relative signs are physical, the gluino mass parameter can be chosen to be positive without lack of generality, so long as both signs for $\mu$ and the $A$–parameters are considered, as is done in the current analysis.

Fig. (2) shows that the $H \to \tau^+\tau^-$ signal strength can be enhanced by more than a factor of three only if $t_1$ is very light, $m_{\tilde{t}_1} \leq 200$ GeV. In this region of parameter space light $\tilde{t}_1$ loops increase the $H$ production cross section by about a factor of three over its SM value. Since $\sin^2(\beta - \alpha)$ saturates its upper bound, the partial width for $H \to WW^*$ is reduced by about 15%, while the partial widths into the $\tau^+\tau^-$ and $b\bar{b}$ final states are increased by factors of 4 and 2.4, respectively; the signals in the $VV$ channels are therefore slightly smaller than in the SM, in spite of the increased production cross section. Note also that light $b_1$ loops again play an important role in suppressing the $Hb\bar{b}$ coupling.

If all squarks are heavier than a few hundred GeV, the $H \to \tau^+\tau^-$ signal can still be enhanced by up to a factor of three, essentially by enhancing the $H\tau^+\tau^-$ couplings since $|\cos\alpha| > \cos\beta$; since $\tilde{b}_1$ is heavy, the $Hb\bar{b}$ coupling will then be enhanced by a similar amount.

Of perhaps greater interest, given current trends in the data, is that the $H \to \tau^+\tau^-$ signal can also be considerably weaker than in the SM. This signal is weakest for the smallest allowed value of $\tan\beta$, and requires $\tilde{t}_1$ and $\tilde{b}_1$ to be relatively light; the former reduces the $H$ production rate via gluon fusion, whereas the latter partly compensates for the reduction of the $Hb\bar{b}$ coupling that originates from the very small values of $|\cos\alpha|$ that can be realized in this region of parameter space. The lower bound on the strength in the $\tau^+\tau^-$ channel is then essentially set by the upper bound (21) on the strength of the signal in the $VV$ channels, which imposes an upper bound on the branching ratios for these channels. For larger squark masses the ratio of the $H\tau^+\tau^-$ and $Hb\bar{b}$ couplings is essentially fixed, independent of the parameters of the Higgs sector, leading to a slightly stronger lower bound on the $H \to \tau^+\tau^-$ signal strength. However, even this increased lower bound is still below conceivable near–future sensitivities in this channel.

Finally, Fig. (2) shows that the double ratio $R_{\gamma\gamma}^H / R_{HVV}^V$ can differ by more than 10% from unity only if $m_{\tilde{t}_1} < 300$ GeV. Note that the production cross section, i.e. the partial width for $H \to gg$, as well as the total width of $H$ cancel out in this double ratio, which is simply given by the ratio of the corresponding partial widths, $\Gamma(H \to \gamma\gamma) / \Gamma(H \to W^+W^-)$, normalized to the same ratio of partial widths of the SM Higgs boson. Since the $HW^+W^-$ coupling, which is proportional to $\cos(\alpha - \beta)$, is only slightly reduced from its SM value, the biggest contribution to radiative $H \to \gamma\gamma$ decays always comes from $W$ loops in this scenario. For $m_{\tilde{t}_1} > 300$ GeV the only other significant contribution comes from top loops, which always interfere destructively with the $W$ loops here, just as in the SM. Due to this destructive interference the reduction of the $HW^+W^-$ (and $HZZ$) coupling implied by the constraint (2) reduces the $H \to \gamma\gamma$ partial width slightly more than the $H \to VV^*$ ($V = W, Z$) partial widths. However, this reduction of the double ratio by $\sim 5\%$ will likely remain unobservable.
at the LHC.

Figure 3: Allowed region in the $(R_{\gamma\gamma}^H, R_{VV}^H)$ (a), the $(R_{TT}^T, R_{TT}^h)$ (b), the $(\mathcal{R}, R_{TT}^T)$ (c) and the $(\mathcal{R}, R_{TT}^h)$ plane (d), after the constraints (2) as well as the various sparticle and Higgs search limits discussed in the text have been imposed; here $R_{\gamma\gamma}^H = R_{VV}^H$. The solid (red) line in frame a) shows $R_{\gamma\gamma}^H = R_{VV}^H$.

For smaller $\tilde{t}_1$ mass the double ratio may differ by up to $\sim 30\%$ from unity. This is due to the effect of $\tilde{t}_1$ and, to a lesser extent, $\tilde{b}_1$ loops on $\Gamma(H\to\gamma\gamma)$; recall that equal soft breaking parameters have been used in the $\tilde{t}$ and $\tilde{b}$ sectors here. Depending on the sign of the $H\tilde{t}_1\tilde{t}_1^*$ coupling, this contribution can interfere constructively or destructively with the dominant $W$ loop contribution; this explains the bifurcation of the results for $m_{\tilde{t}_1} < 400$ GeV. In either case the effect is maximized for large mass splitting between the $\tilde{t}$ mass eigenstates and large $|\mu|$, with maximal suppression (enhancement) of the double ratio requiring positive (negative) $\mu$. A light $\tilde{\tau}_1$ can suppress the double ratio by an additional 3% or so.

The correlation between the two signal rates defining the double ratio is explored in the first frame of Fig. 3, which shows correlations between various (ratios of) signal strengths. The red line corresponds to $R_{\gamma\gamma}^H = R_{VV}^H$, leading to a unit value for the double ratio. We see that both signal strengths can saturate their lower bounds defined in (2d,e), and that $R_{TT}^h$ can also saturate its upper
bound. Evidently relaxing the bounds on $R_{HVV}^{VV}$ would also lead to an increased allowed range for $R_{VV}^{VV}$, beyond the range shown in (7). On the other hand, neither $R_{H}^{\gamma\gamma}$ nor $R_{HH}^{VV}$ is strongly correlated with the ratio between these two quantities: the allowed range of the ratio of signal strengths moves only slightly towards smaller values as $R_{H}^{\gamma\gamma}$ increases, such that $R_{H}^{\gamma\gamma} \lesssim 1.1 R_{H}^{VV}$ when $R_{H}^{VV}$ saturates its upper bound of 2; this is to be compared with the absolute upper bound of 1.3 on the ratio, see eq. (8).

Fig. 3b shows that the $h$ and $H$ signals in the di–tau channel are positively correlated. The reason is that increasing $\tan \beta$ increases the basic $\tau$ Yukawa coupling in the Lagrangian, which therefore also tends to increase the couplings of both $h$ and $H$ to $\tau$ leptons. However, for small value of $R_{H}^{\gamma\gamma}$ this correlation is not particularly strong: the $H \rightarrow \tau^+\tau^-$ signal can then be both significantly stronger and significantly weaker than in the SM. On the other hand, the $h \rightarrow \tau^+\tau^-$ signal strength can only exceed that of the SM significantly if the $H \rightarrow \tau^+\tau^-$ signal is also enhanced. Current data disfavor an enhanced signal in the di–tau channel for the new boson near 125 GeV; in the present context a significant upper bound on this signal would make it even more difficult to detect $h$ at the LHC.

The correlation between the $h \rightarrow \tau^+\tau^-$ signal strength and the double ratio $R_{H}^{\gamma\gamma}/R_{HH}^{VV}$ is explored in Fig. 3c. We see that the former can only be enhanced significantly beyond its SM value if the latter is somewhat below unity. Recall from the discussion of Fig. 2a that maximizing $R_{H}^{\gamma\gamma}$ requires small $m_{\tilde{g}}$. In the case at hand this enhances the partial width for $H \rightarrow gg$, but reduces the partial width for $H \rightarrow \gamma\gamma$, leading to a reduction of the ratio of signal strengths in the $\gamma\gamma$ and $VV^*$ channels relative to their SM value. Again the current data favor this double ratio to be enhanced; Fig. 3c shows that this would reduce the upper bound on the $h \rightarrow \tau^+\tau^-$ signal strength in this scenario.

Finally, Fig. 3d shows the correlation between the double ratio of $H \rightarrow \gamma\gamma$ and $H \rightarrow VV^*$ signal strengths and the $H \rightarrow \tau^+\tau^-$ signal. These quantities are clearly anti–correlated in the present scheme. The $H \rightarrow \tau^+\tau^-$ signal is maximized in a similar region of parameter space as the $h \rightarrow \tau^+\tau^-$ signal; we just saw that this leads to a suppression of the double ratio.

Conversely, the double ratio reaches its maximum when both the $H \rightarrow \gamma\gamma$ and $H \rightarrow VV^*$ signals are suppressed by destructive interference of top and stop loop contributions to $H \rightarrow gg$; stop loop contributions then maximally enhance $\Gamma(H \rightarrow \gamma\gamma)$. The suppression of the $H$ production cross section also reduces the strength of the signal in the di–tau channel. One can also find configurations with slightly less enhanced double ratio where both the $H \rightarrow \gamma\gamma$ and the $H \rightarrow VV^*$ signals are enhanced over their SM values. This can be achieved if $|\cos \alpha| < \cos \beta$, which suppresses the $H \rightarrow b\bar{b}$ and $H \rightarrow \tau^+\tau^-$ partial widths, and thus the total decay width of $H$. This mechanism also leads to a suppression of the $H \rightarrow \tau^+\tau^-$ signal. Finally, the branch with $R_{H}^{\gamma\gamma}/R_{HH}^{VV} \simeq 1.2$, $R_{H}^{\gamma\gamma} \simeq 0.3$ requires very large and negative $\mu$, large and positive $A_{t}$, and (as usual) large splitting between $m_{\tilde{t}\tilde{L}}$ and $m_{\tilde{e}_{\tilde{L}}}$. The light $\tilde{t}_{1}$ loops then again suppress $H \rightarrow gg$ decays and enhance $H \rightarrow \gamma\gamma$ decays, whereas $|\cos \alpha| < \cos \beta$ reduces the total width of $H$; the former effect is dominant, i.e. the $H \rightarrow \gamma\gamma$ and $H \rightarrow VV^*$ signals are both suppressed relative to their SM values. In this case the di–tau signal is further suppressed because the light $\tilde{b}_{1}$ enhances, rather than suppresses, the ratio of $Hb\bar{b}$ and $H\tau^+\tau^-$ coupling.

4 Summary and Conclusions

In this paper I have shown that one can explain both the recent discovery of a “Higgs–like particle” by the LHC experiments, and the $2.3\sigma$ excess of Higgs–like events found by the LEP collaborations.
some ten years ago, in the “phenomenological” MSSM, where all (relevant) weak–scale soft breaking parameters are treated as independent free parameters. In this interpretation the masses of the two CP–even Higgs bosons $h$ and $H$ are essentially fixed by the data, to $\sim 98$ and $\sim 125$ GeV, respectively. Radiative corrections to the Higgs sector are crucial for the viability of this scheme. As a result, the masses of the remaining Higgs bosons, the CP–odd state $A$ and the charged state $H^\pm$, can still vary considerably. Nevertheless stringent upper bounds on the masses of these states can be derived, which would be straightforward to test at an $e^+e^-$ collider operating at $\sqrt{s} \geq 350$ GeV. Much of the allowed parameter space can probably also be probed by searches for these states at the LHC; in particular, $t \to H^+b$ decays are open over almost the entire parameter space. However, this is not sufficient to guarantee that such decays, or other $A$ or $H^\pm$ production processes, can actually be detected at the LHC.

The upper bound on $m_{H^\pm}$ implies that loops involving the charged Higgs boson and the $t$ quark will give significant positive contributions to the partial width for radiative $b \to s\gamma$ decays [33]; if these were the only new contributions the predicted partial width would exceed the measured value [29], which is quite close to the SM prediction. However, it is well known that even within the MSSM with minimal flavor violation, chargino–stop loops can cancel the charged Higgs loops [34], so a portion of the parameter space (with rather light $\tilde{t}_1$) is most likely allowed even in this constrained scenario. Moreover, as argued in the Introduction, the general MSSM contains many additional parameters that can be tuned to satisfy flavor constraints. In particular, a small amount of $\tilde{b} - \tilde{s}$ mixing would lead to large gluino loop contributions to $b \to s\gamma$ [35] of either sign.

The light state $h$ is also very difficult to detect at the LHC. It has greatly reduced couplings to $Z$ and $W$ bosons, and hence also a greatly reduced branching ratio into $\gamma\gamma$ final states. Part of the allowed parameter space could perhaps be probed through $h \to \tau^+\tau^-$ decays, but the small value of $m_h$ implies that $Z \to \tau^+\tau^-$ decays will be a formidable background.

Although the couplings of $H$ to $W$ and $Z$ bosons are quite SM–like in this scenario, both the production cross section and the decay branching ratios of $H$ can still differ significantly from those of the SM Higgs. On the one hand, the couplings of $H$ to third generation fermions can be quite different from those of the SM Higgs; here light $\tilde{b}$ loop contributions to the $Hb\bar{b}$ coupling can play a significant role. A good measurement of, or upper bound on, the strength of the di–tau signal would therefore narrow down the allowed parameter space of this scenario. A reliable observation of the $H \to b\bar{b}$ signal and/or $ttH$ production would be similarly useful, but are experimentally (even) more challenging. Moreover, light $\tilde{t}$ loops can modify the partial widths for $H \to gg$ and, to a somewhat lesser extent, for $H \to \gamma\gamma$ decays significantly. In particular, in this scenario one can simultaneously reduce the di–tau signal and enhance the di–photon signal, in agreement with the (statistically not very compelling) trend of current data.

However, the di–photon signal can be enhanced relative to the $VV^* \ (V = W^\pm, Z)$ signals only if $\tilde{t}_1$ is rather light. Ongoing and future searches for light stop and sbottom squarks therefore have the potential to further constrain the parameter space of this model. Unfortunately the interpretation of such searches also depends on the chargino and neutralino sectors of the MSSM, which have not been specified here, since they hardly affect the Higgs sector. Note, however, that the LHC experiments should eventually be able to probe $\tilde{t}_1$ masses well above 200 GeV even in the experimentally most difficult case where $\tilde{t}_1$ is nearly degenerate with a stable neutralino [36]. If all squarks are heavy, $H$ can be so SM–like that it probably cannot be distinguished from the Higgs boson of the SM. In particular, the squared couplings to $W$ and $Z$ are just 5 to 15% smaller than in the SM; this follows from the normalization of the excess observed at LEP.
It should be admitted that this scenario is theoretically not especially appealing. In particular, the LEP excess cannot be explained in a constrained version of the MSSM \[37\]. Moreover, the lower bound on the sum of the stop masses, and hence the required finetuning associated with radiative corrections from the top and stop sector, is similar to that in the more common MSSM interpretation of the LHC discovery in terms of production and decay of the light CP–even state $h$. It is nevertheless amusing to note that the MSSM can *simultaneously* explain two sets of observations of excesses of Higgs–like events.

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