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**Predictions for Higgs and SUSY Higgs properties and their signatures  
at the Hadron Colliders.<sup>a</sup>**

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In this talk I shall present a discussion of the theoretical bounds on the mass of the Higgs in the Standard Model (SM) as well as in the Minimal Supersymmetric Standard Model (MSSM). Then I will point out a few facts about the couplings of scalars that are relevant for its search at hadronic colliders. After that I discuss the search possibilities at the Tevatron and the LHC, paying special attention to the issue of how well one can establish the quantum numbers and the couplings of the Higgs, when (if) it is discovered.

## 1 Introduction

I would concentrate here on the theoretical bounds on Higgs mass both in the SM and the MSSM as well as on the theoretical information about its couplings which are relevant for the Higgs search at the Tevatron and the LHC. This would be followed by a discussion of the search prospects for the Higgs at both these colliders. I will also address the issue of the feasibility of establishing the quantum numbers and the couplings of a spin zero particle when it is discovered at these colliders. This would be necessary to establish it as **the** Higgs boson which arises from the Higgs mechanism of Spontaneous Symmetry Breakdown (SSB). We will see that to achieve the latter, search at the hadron colliders needs to be complimented by that at a high energy  $e^+e^-$  collider (the next linear collider NLC) <sup>1</sup>.

## 2 Higgs Couplings and masses: Theoretical predictions

### 2.1 Predictions in the SM:

In the SM the existence of Higgs boson is necessary to bring about the SSB which gives masses to the fermions and the gauge bosons, still keeping the theory renormalisable. For the SSB to happen, the (mass)<sup>2</sup> term for the complex scalar doublet  $\Phi$  has to be negative, i.e. the potential  $V(\Phi)$  is

$$V(\Phi) = \frac{\lambda}{4!}(\Phi^\dagger \Phi)^2 - \mu^2 \Phi^\dagger \Phi \quad (1)$$

with  $\mu^2$  positive. After the SSB, out of the four scalar fields which comprise  $\Phi$ , we are left with only the physical scalar  $h$  with a mass

$$m_h^2 = \lambda v^2 \quad (2)$$

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Further, the tree level couplings of the Higgs boson  $h$  to the SM fermions and the gauge bosons are uniquely determined and are proportional to their masses. The coupling of a Higgs to a pair of gluons/photons does not exist at the tree level, but is induced at one loop level by the diagrams shown in Fig. 1. As with the other

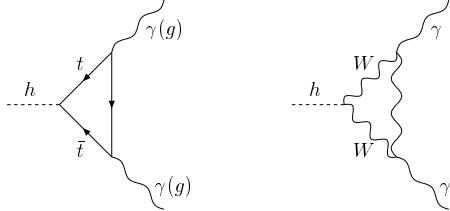


Figure 1: Loop diagrams responsible for  $h \rightarrow \gamma\gamma(gg)$ .

couplings, this coupling is also completely calculable given the particle content of the SM, to a given order in the strong and electromagnetic coupling<sup>3</sup>. The  $hgg$  coupling is dominated by the top quark contribution in the loop diagram whereas the  $h\gamma\gamma$  coupling receives dominant contribution from both, the top loop as well as the  $W$  loop. These channels have appreciable branching ratios, albeit very small, only for  $m_h \lesssim 2m_W$ . Recall also that the precision measurements at the  $Z$  indicate  $m_h \lesssim 300$  GeV<sup>2</sup>. It should be noted that the QCD corrections for  $\Gamma(h \rightarrow gg)$  are significant ( $\sim 65\%$ )<sup>3</sup>.

In the intermediate mass range (i.e.  $m_h \leq 140$  GeV) the total width of the Higgs is  $\lesssim 10$  MeV, dominant decay is into a  $b\bar{b}$  final state (e.g. for  $m_h = 120$  GeV  $\Gamma(h \rightarrow b\bar{b}) \simeq 68\%$ ); on the other hand the branching ratio into a  $\gamma\gamma$  final state is about one part per mille. The total width is  $\sim 1$  GeV around  $m_h \sim 300$  GeV and rises very fast after that, reaching  $\Gamma_h \sim m_h$  around  $m_h \gtrsim 500$  GeV. Calculations of various branching ratios, including higher order QCD effects are available<sup>3</sup>.

While the various couplings and hence the branching fractions of the Higgs are well determined once  $m_h$  and various other parameters in the SM such as  $m_t, \alpha_s$  etc. are specified,  $m_h$  itself is completely undetermined in the SM. However, as seen from Eq. 2, it is linearly related to the self coupling of the scalar field. Eventhough, the theory has nothing to say about the Higgs mass *per se* the behaviour of the self coupling  $\lambda$  is determined by field theory. This then puts bounds on  $m_h$ . These bounds can be understood as follows. The self coupling receives radiative corrections from the diagrams indicated in Fig. 2. The contributions from the diagrams involving the scalar and the gauge boson loops on one hand and the fermion loops on the other, are opposite in sign. The requirement that the self coupling  $\lambda$  stay positive, i.e. the vacuum remain stable under radiative corrections, puts a lower bound on  $m_h$  for a given value of  $m_t$ <sup>4</sup>. This bound however, depends on the  $htt$  coupling and hence can be evaded in models with more than one Higgs doublets.  $m_h$  is also bounded from above by considerations of triviality. This can be understood by considering only the contributions of the scalar loops, for simplicity, to the

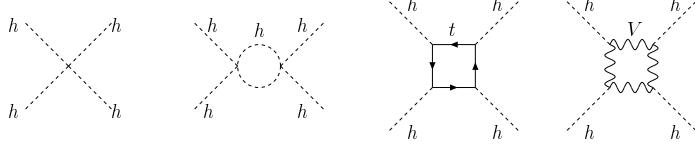


Figure 2: Diagrams responsible for radiative corrections to self coupling.

radiative corrections to  $\lambda$ . It can be shown that the self coupling then satisfies,

$$\frac{d\lambda(t)}{dt} = \frac{3}{4\pi^2} \lambda^2(t) \quad (3)$$

with  $t = \ln(\Lambda/v)$ ; where  $\Lambda$  is the momentum scale at which the coupling  $\lambda$  is evaluated. This equation needs a boundary condition to solve it, which is chosen as

$$\lambda = \lambda(\Lambda = v) = \lambda(1) = \sqrt{2}G_F m_h^2. \quad (4)$$

Then Eq. 3 above can be solved to give

$$\lambda(t) = \frac{\lambda}{1 - 3\lambda t/4\pi^2} \quad (5)$$

This shows that  $\lambda(t)$  will diverge at high scales. If we demand that the Landau pole, where  $\lambda(t)$  will blow up, lies above a scale  $\Lambda$ , we get

$$m_h \lesssim \frac{893}{\sqrt{\ln(\Lambda/v)}} \text{GeV}. \quad (6)$$

Hence the requirement that the theory be valid at large  $\Lambda$  and yet be nontrivial at a scale  $v$ , puts an upper limit on  $\lambda(v)$  and hence on  $m_h$  due to the identification in Eq. 2. The above analysis, of course, has to be improved using the renormalisation group equation<sup>5,6</sup>. As  $\lambda$  becomes large, clearly perturbative methods used above must fail. In the region of large  $\lambda$  the analysis has been done using lattice theory<sup>7</sup>.

The lower bound on  $m_h$  implied by the vacuum stability arguments<sup>4</sup> and the upper bound implied by the triviality considerations<sup>5,6</sup>, depend on the value of  $m_t$  and the uncertainties in the nonperturbative dynamics respectively. The resulting bands, taking into account these theoretical uncertainties, are shown in Fig. 3<sup>8</sup>.

The results can be summarised as follows:

1. If we demand that the Landau pole lies above  $10^{15}(10^{18})$  GeV and there exists no new physics other than the SM upto that scale, then  $m_h < 190(130)$  GeV.
2. If we assume that the SM is an effective theory only upto 1 TeV, i.e., there is some new physics at that scale then  $m_h \lesssim 800$  GeV. As an aside let us also mention here that this upper limit of 800 GeV is of the same order as the limit obtained by requiring that  $WW \rightarrow WW(ZZ \rightarrow ZZ)$  amplitude satisfies perturbative unitarity<sup>9,10</sup>.

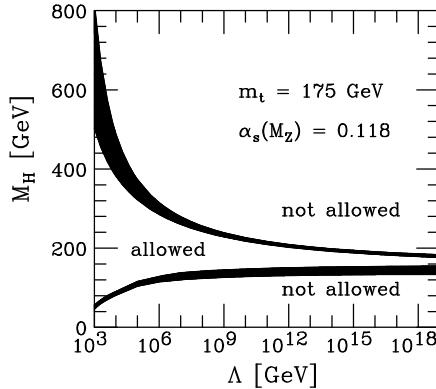


Figure 3: Theoretical bounds on  $m_h$  in the SM.

## 2.2 Masses and couplings in the MSSM

In the MSSM there exist two complex Higgs doublets  $\Phi_1, \Phi_2$  with hypercharge  $Y = \pm$  respectively. As a result there are five physical degrees of freedom left after the spontaneous symmetry breakdown. The MSSM thus contains, in all, five scalars: three neutrals out of which two are CP even states denoted by  $h_0, H_0$  and one is a CP odd state denoted by  $A$  and a pair of charged Higgs bosons  $H^\pm$ . Thus the scalar sector of the MSSM is much richer than in the SM.  $h_0$  denotes the lighter of the two CP even neutral scalars.

The most general scalar potential for a two Higgs doublet model contains a large number of free parameters, essentially analogues of the self coupling  $\lambda$  and  $\mu^2$  term in the case of the SM. However, supersymmetry either fixes all these self couplings in terms of the gauge couplings or requires them to vanish. Hence the scalar potential for the MSSM is

$$\begin{aligned}
 V = & m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - \left[ m_{12}^2 \Phi_1^\dagger \Phi_2 + \text{h.c.} \right] \\
 & + \frac{1}{8} (g^2 + g'^2) \left[ (\Phi_1^\dagger \Phi_1)^2 + (\Phi_2^\dagger \Phi_2)^2 \right] \\
 & + \frac{1}{4} (g^2 - g'^2) (\Phi_2^\dagger \Phi_2) (\Phi_1^\dagger \Phi_1) - \frac{1}{2} g^2 (\Phi_1^\dagger \Phi_2) (\Phi_2^\dagger \Phi_1),
 \end{aligned} \tag{7}$$

where  $g, g'$  are the  $SU(2), U(1)$  coupling constants respectively. After the SSB where the neutral members of the two doublets acquire vacuum expectation values  $v_1/\sqrt{2}, v_2/\sqrt{2}$  respectively, the resulting masses of the five physical scalars that follow from the abovementioned potential, can be written in terms of two parameters which can be chosen to be  $m_A, \tan \beta = v_1/v_2$  or equivalently  $m_{H^\pm}, \tan \beta$ . At the tree level the five scalar masses satisfy the following inequalities:

$$m_{h_0} \leq m_Z, \quad m_{H_0} > m_Z, \quad m_{H^\pm} > m_W, \quad m_{h_0} < m_{H_0}, m_{H^\pm}. \tag{8}$$

In the decoupling limit <sup>11</sup> ( $m_A \rightarrow \infty$ ) one finds that, independent of  $\tan \beta$ , all the four heavy scalars become degenerate and infinitely heavy and the mass of the

lightest scalar approaches the upper bound. In this limit the couplings of the  $h_0$  to matter fermions and the gauge bosons approach those of the SM higgs  $h$ . The interesting thing to note here is that all the masses  $m_{H^\pm}, m_A, m_{h_0}$  can become large without some self coupling becoming strong, unlike the case of the SM.

If we denote by  $\alpha$  the mixing in the two CP even neutral fields to yield the mass eigenstates  $h_0, H_0$ , then the couplings of all the three neutral scalars, with the fermions and gauge bosons are given in Table 1. For the MSSM scalars only

Table 1: Couplings of the three neutrals with fermions and gauge bosons

	$h$	$h_0$	$H_0$	A
$b\bar{b}$	$gm_b/2m_W$	$\sin \alpha / \cos \beta$ $\rightarrow 1$	$\cos \alpha / \sin \beta$ $\tan \beta$	$\tan \beta$ $\tan \beta$
$t\bar{t}$	$gm_t/2m_W$	$\cos \alpha / \sin \beta$ $\rightarrow 1$	$\sin \alpha / \cos \beta$ $\cot \beta$	$\cot \beta$ $\cot \beta$
$VV$	$gm_V$	$\sin(\beta - \alpha)$ $\rightarrow 1$	$\cos(\beta - \alpha)$ $0$	$0$ $0$

the additional factors relative to the SM case are written down. These couplings in the *decoupling limit* reduce to the ones denoted in the second line in each case. We notice that in the decoupling limit, where the upper limit on the mass of the lightest scalar is saturated, the couplings of the lightest scalar approach those for the SM higgs  $h$ . It should also be noted that the CP odd scalar A does not couple to a pair of gauge bosons at the tree level.

The inequalities of Eq. 8 get affected by the radiative corrections to the Higgs masses. The dominant corrections arise from loops involving top (t) and its scalar partner stop ( $\tilde{t}$ ) due to the large Yukawa coupling of the top quark. There are many different methods, involving different methods of approximations<sup>12</sup>, to calculate these corrections. These depend on  $m_t$  as well as the masses of and the mixing between  $\tilde{t}_L, \tilde{t}_R$  (the superpartners of  $t_L$  and  $t_R$ ). To a good approximation, the radiatively corrected upper bound on the mass of the lightest CP neutral scalar ( $m_{h_0}$ ), can be written as

$$m_{h_0}^2 < m_Z^2 \cos^2 2\beta + \epsilon + \epsilon_{\text{mix}}, \quad (9)$$

where

$$\begin{aligned} \epsilon &= \frac{3g^2 m_t^4}{8\pi^2 m_W^2} \ln \left( \frac{m_{\tilde{t}}^2}{m_t^2} \right) \\ \epsilon_{\text{mix}} &= \frac{3g^2 m_t^4}{8\pi^2 m_W^2} \frac{A_t^2}{m_{\tilde{t}}^2} \left( 1 - \frac{A_t^2}{12m_{\tilde{t}}^2} \right), \end{aligned} \quad (10)$$

with  $A_t$  being the coefficient of the trilinear, supersymmetry breaking term, and  $m_{\tilde{t}}$  being the common mass of the  $\tilde{t}_L, \tilde{t}_R$ . The second term  $\epsilon_{\text{mix}}$ , even though dependent

on  $A_t$ , can be shown to be bounded by  $9g^2m_t^4/8\pi^2m_W^4$ . Thus  $m_{h_0}$  is still bounded even though the bound on  $m_{h_0}$  of Eq. 8 is changed by radiative corrections. Also note that the corrections will vanish in the limit of exact supersymmetry. The limits on the radiatively corrected scalar masses for the case of maximal mixing in the stop

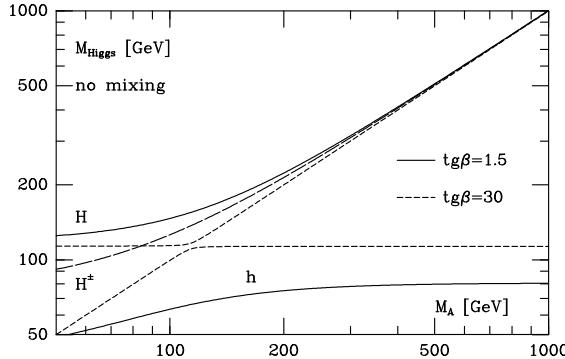


Figure 4: Bounds on the masses of the scalars in the MSSM

sector are shown in Fig. 4<sup>13</sup>. It can be seen from the results shown in the figure that the mass of the lightest scalar in MSSM is bounded by  $\sim 130$  GeV even after it is radiatively corrected. This bound does get modified in the NMSSM<sup>14,15,16</sup>. Again, it has been shown that for all reasonable values of the model parameters,  $m_{h_0}$  is bounded by  $\sim 150$  GeV. From the figure, it might seem that the current LEP bound on the mass of the  $m_{h_0}$ ,  $m_{H_0}$  and  $m_A$  rule out the low  $\tan\beta$  values for the MSSM. However, it has been shown<sup>17</sup> that a rather minor extension of the MSSM, can help avoid this conclusion. The couplings of Table 1 do get modified to some extent by the radiative corrections, but the general features discussed above remain unchanged.

In discussing the search strategies and prospects of the MSSM scalars one has to remember the following important facts:

1. Due to the reduction of the  $h_0WW$  coupling as indicated in Table 1, the  $h_0\gamma\gamma$  coupling is suppressed as compared to the corresponding SM case. Of course one also has to include the contribution of the charged sparticles in the loop<sup>18</sup>. Due to the upper limit on  $m_{h_0}$  the decay mode into  $WW(VV)$  pair is not possible for  $h_0$ , due to kinematic reasons. On the other hand, for  $H_0$  the suppression of the coupling to  $VV$  as seen from Table 1, makes the decay less probable as compared to the SM case. As a result, the MSSM scalars are expected to be much narrower resonances as compared to the SM case. For example, the maximum width of  $h_0$  is less than few MeV, for reasonable values of  $\tan\beta$  and even for the heavier scalars  $H_0, A$ , the width is not more than few tens of GeV even for masses as high as 500 GeV.
2.  $h_0$  is much narrower than the SM Higgs. However, over a wide range of  $\tan\beta, m_A$  values, the  $h_0$  has dominant decay modes into Supersymmetric par-

ticles. The most interesting ones are those involving the lightest neutralinos, which will essentially give ‘invisible’ decay modes to the  $h_0$ ,  $H_0$  and  $A$  <sup>19</sup>.

3. On the whole for the MSSM scalars the decay modes into fermion-antifermion pair are the dominant ones due to the point (1) above as well as the fact that the CP odd scalar  $A$  does not have any tree level couplings to  $VV$ . Hence, looking for the  $\tau^+\tau^-$  and  $b\bar{b}$  final state becomes very important for the search of the MSSM scalars.

It is clear from the above that the phenomenology of the MSSM scalars is much richer and more complicated than the SM case. Again, calculation of various decay widths including the higher order corrections has been done<sup>3,18</sup>.

### 3 Production and search of Higgs at Colliders

In this section we will begin by discussing the search possibilities for the SM Higgs  $h$ . As is clear from the discussions in the earlier section of the couplings of the Higgs, the most efficient way of producing the Higgs scalar at any Collider is through its coupling to gauge bosons or to a heavy fermion-antifermion ( $t\bar{t}$ ) pair. At the hadronic colliders the following processes can contribute to the production:

$$gg \rightarrow h \quad (11)$$

$$q\bar{q}' \rightarrow hW \quad (12)$$

$$q\bar{q} \rightarrow hZ \quad (13)$$

$$q\bar{q} \rightarrow hqq \quad (14)$$

$$gg, q\bar{q} \rightarrow ht\bar{t}, hbb. \quad (15)$$

At the Tevatron energies ( for Run-II/TeV33) the most efficient processes are those of Eqs. 11,12. The potential of direct Higgs search at the Tevatron is discussed elsewhere in the proceedings<sup>20</sup>. The associated  $W/Z$  can make it possible to use the dominant  $b\bar{b}$  decay mode. However, the energy of the Tevatron is just too small to give appreciable production cross-section, except for the values of  $m_h$  close to its current lower limit of 92.5 GeV<sup>2</sup>. Various strategies for using the  $WW^*$  or the  $b\bar{b}$  decay mode for  $h$  to enhance the mass reach of the Tevatron have been suggested<sup>21</sup>.

At LHC energies, due to the large available gluon fluxes and the large value of  $m_t$ , Eq. 11 is the dominant production process for all values of  $m_h$  upto the upper bounds discussed in section 2.1. The total production cross-section goes from  $\sim 20$  pb to  $\sim 0.1$  pb as  $m_h$  goes from 100 to 1000 GeV. For a superheavy Higgs ( $\gtrsim 800$  GeV i.e. above the highest upper bound of sec. 2.1), the process of Eq. 14 has significant cross-section. The higher order QCD corrections to the gluon induced higgs production are significant and have been included in the available theoretical predictions shown in Fig. 5<sup>3</sup>. These predictions, at the LHC, have a typical theoretical uncertainty of  $\sim 20\%$  due to the parton densities.

Since LEP-2 has already ruled out, from direct search,  $m_h < 92.5$  GeV<sup>2</sup>, the range of masses of interest to LHC divides neatly into two parts :  $92.5 < m_h \lesssim 140$  GeV and  $m_h > 140$  GeV. In the first mass region the dominant decay mode of

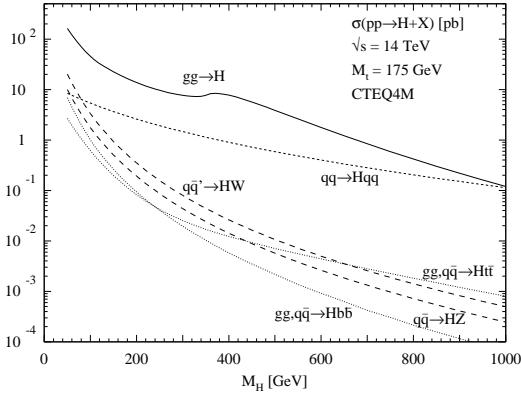


Figure 5: Production cross-section for the SM Higgs at LHC

the  $h$  is into a  $b\bar{b}$  final state which has a QCD background about  $10^3$  higher than the signal. Hence, in this mass range  $h \rightarrow \gamma\gamma$  remains the best final state, even with a branching ratio  $\sim 10^{-3}$ . Even then, the resolution required for the  $M_{\gamma\gamma}$  measurement has to be  $\lesssim 1$  GeV  $\simeq .1\% m_h$ .

The new developments in the detection aspect have been detector simulations for the  $\gamma\gamma$  and  $b\bar{b}$  mode for the planned detector designs. Fig. 6 taken from the

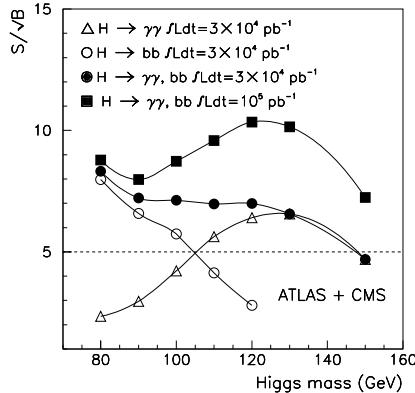


Figure 6: The expected significance level of the SM Higgs signal at LHC for the intermediate mass region

CMS/ATLAS technical proposal<sup>22</sup> shows the expected  $S/\sqrt{B}$  for the SM Higgs in the intermediate mass range, using the  $\gamma\gamma, b\bar{b}$  modes. The use of  $b\bar{b}$  mode for  $m_h < 100$  GeV, is essentially achieved by using the associated production of Eq. 12.

Once  $m_h \gtrsim 140$  GeV, the  $VV^*$  or  $VV$  decay mode is dominant. As a result, for  $140 < m_h < 600$  GeV, the 4 lepton final state following from  $h \rightarrow ZZ \rightarrow l^+l^-l^+l^-$  offers the so called ‘gold plated signal’ for the Higgs. For  $800 < m_h < 1000$ , the very forward jets produced in association with the Higgs in the process of Eq. 14

provide a much better signature than the ‘gold plated signal’. However, it should be borne in mind, that the unitarity and triviality bounds of section 2.1 imply that in the SM  $m_h \lesssim 600 - 800 \text{ GeV}$ .

Thus we see that while the detection of an intermediate mass Higgs is difficult but feasible at LHC, it will surely require the high luminosity run. It should also be kept in mind that a low value of  $m_h$  seems to be preferred by the precision measurements at LEP and SLC and further that SUSY also predicts that the lightest scalar in the theory will be in this mass range. For higher  $m_h$  values the detection is a certainty at LHC. However, to establish such a scalar as *the SM Higgs*, one needs to establish

1. The scalar is CP even and has  $J^P = 0^+$ ,
2. The couplings of the scalar with the fermions and gauge bosons are proportional to their masses.

This is also essential from the point of view of being able to distinguish this scalar from the lightest scalar expected in the MSSM. We see from Table 1 that the couplings of the scalars to the fermions and gauge bosons can be quite different in the MSSM.

As a matter of fact this issue has been a subject of much investigation of late<sup>23,24</sup>. The Snowmass Studies<sup>23</sup> indicate that for a light Higgs ( $m_h = m_Z$ ) it is possible only to an accuracy of about 30 %. It is in this respect that the planned  $e^+e^-$  colliders<sup>1</sup> can be a lot of help. At these colliders, the production processes are the same as given in Eqs. 13-15 where  $q(\bar{q})$  are replaced by  $e^-(e^+)$ . This is not an appropriate place to give a complete discussion of the search prospects for the SM (and MSSM) Higgs at these colliders. But suffice it to say that if the production is kinematically allowed, detection of the Higgs at these machines is very simple as the discovery will be signalled by very striking features of the kinematic distributions. Determination of the spin of the produced particle in this case will also be simple as the expected angular distributions will be very different for even and odd parity. Even with this machine one will need a total luminosity of  $200 \text{ fb}^{-1}$ , to be able to determine the ratio of  $BR(h \rightarrow c\bar{c})/BR(h \rightarrow b\bar{b})$ , to about 7%<sup>24</sup>. The simplest way to determine the CP character of the scalar will be to produce  $h$  in a  $\gamma\gamma$  collider, which are being discussed. There are also interesting investigations<sup>24</sup> which try to device methods to determine the CP character of the scalar using hadron colliders.

For MSSM Higgs the discussion of the actual search possibilities is much more involved and has been covered in other talks at this conference<sup>25</sup>. For the lightest scalar  $h_0$  in the MSSM, the general discussions of the intermediate mass Higgs apply, with the proviso that the  $\gamma\gamma$  branching ratios are smaller for  $h_0$  and hence the search that much more difficult. However, since there exist many more scalars in the spectrum now, one can cover the different regions in the parameter space by looking for  $A, H$  and  $H^\pm$ . At low values of  $m_A$  these other scalars are kinematically accessible at LHC and also at the NLC. However, even after combining the information from various colliders (LEP-II, Tevatron (for the charged higgs search) and of course LHC), a certain region in the  $m_A - \tan\beta$  plane remains inaccessible. This hole can be filled up only after combining the data from the CMS and ATLAS

detector for 3 years of high luminosity run of LHC. Even in this case there exist large region where one will see only the single light scalar.

At large  $m_A$  (which seem to be the values preferred by the current data on  $b \rightarrow s\gamma$ ), the SM higgs and  $h_0$  are indistinguishable as far as their couplings are concerned, as can be seen from Table 1. A recent study, gives the contours of constant values for the ratio

$$\frac{BR(c\bar{c})/BR(b\bar{b})|_{h_0}}{BR(c\bar{c})/BR(b\bar{b})|_h},$$

as well as a similar ratio for the  $WW^*$  and  $b\bar{b}$  widths as a function of  $\tan\beta$  and  $m_A$ . As we can see from Fig. 7<sup>23</sup>, a measurement of this ratio to an accuracy of

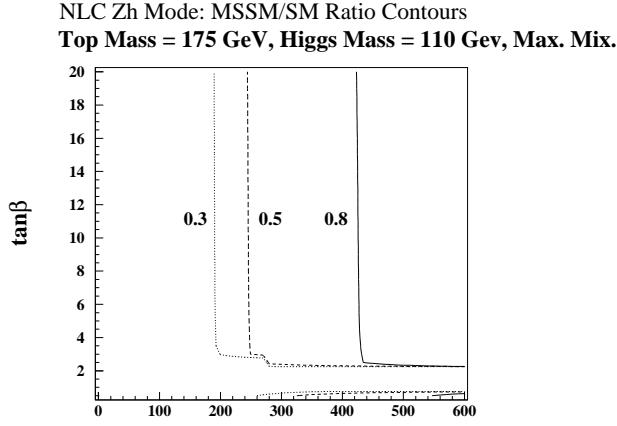


Figure 7: The ratios of relative branching fractions for the MSSM and SM for maximal mixing in squark sector. The specific value of the Higgs mass used is theoretically disallowed at large  $m_A$  and around  $\tan\beta \sim 2$ .

about 10% will allow distinction between the SM Higgs  $h$  and MSSM Higgs  $h_0$  for  $m_A$  as large as  $\sim 600$  GeV. As stated above, NLC should therefore be able to do such a job. Certainly, the issue of being able to determine the quantum numbers and the couplings of the scalar accurately, forms the subject of a large number of investigations currently.

#### 4 Conclusions

1. The experiments at the upcoming hadronic colliders will be able to ‘discover’ scalar in the entire mass range from 100 GeV to 1000 GeV, with varying degree of ease. The lower end is difficult but feasible.
2. If the experiments at the Tevatron (Run-II and TeV33, should it happen) and the LHC do not find a light scalar upto  $m_h = 130(160)$  GeV, it will certainly rule out a class of SUSY models; viz. MSSM and models with minimal extension of the MSSM: the (N)MSSM. The only way this upper

bound can be relaxed is if we enlarge the gauge group and introduce additional scales in the problem.

3. If the scalar that is ‘found’ has a mass  $\gtrsim 800$  GeV, it is an indication that the EW symmetry breakdown happens via strongly interacting sector.
4. If we *do* find a **light** Higgs, then all we can conclude is that the SM works as an effective theory upto large scales, as discussed in section 2.1. It will also be consistent with the indirect mass limits obtained from precision measurements. This limit is more or less insensitive to existence of SUSY (or otherwise) due to large mass scales to which sparticle masses have already been pushed by the lack of direct observation of the SUSY particles
5. Since in Supersymmetric models the masses and the decay modes of the various scalars in the theory are correlated, almost all the region in the parameter space of the MSSM can be explored at the experiments at the LHC<sup>25</sup>.
6. Already with the current data (particularly the information on the  $b \rightarrow s\gamma$ ) the limits on the MSSM parameter space are such that the lightest scalar  $h_0$  will be very similar in its properties to the SM Higgs  $h$ . Hence, it is important to devise strategies to determine the quantum numbers such as Spin, Parity, CP etc., of this observed scalar. To that end a TeV ( $\gtrsim 300$  GeV) linear  $e^+e^-$  collider seems indispensable. Also, the possibilities of achieving this at the LHC/Run-II/TeV-33 need to be investigated vigorously.

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