HIGGS BOSONS AT AN $e^+e^-$ LINEAR COLLIDER: THEORY AND PHENOMENOLOGY

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Abstract

I review recent developments in the theory and phenomenology of Higgs bosons at an $e^+e^-$ linear collider with $\sqrt{s}$ of order 500 GeV.

1. Introduction

Perhaps the most fundamental mission of a future high energy $e^+e^-$ collider will be to reveal the nature and source of electroweak symmetry breaking. This could prove to be a relatively straightforward task at an $e^+e^-$ collider in the case of the minimal Standard Model (MSM). If the single Higgs boson of the model (the $\phi^0$) is sufficiently light, a narrow Higgs resonance will be found, and the interactions of (longitudinally polarized) $W$ ($W \equiv W^\pm, Z$) bosons will be perturbative at all energies. Adequate machine energy will be a necessity, but $\sqrt{s}$ in the range from 500 GeV to 1 TeV should suffice. If the $W$ boson sector is strongly interacting, a thorough investigation of all $WW$ scattering channels will be required before one can hope to fully understand electroweak symmetry breaking, and much higher machine energy (2 – 4 TeV) will be required. This review will focus entirely on the perturbative scenario and on $\sqrt{s}$ below 1 TeV.\textsuperscript{*} The summary talk on Higgs physics at the previous meeting in this series, Ref. 1, is a convenient source of background material.

Even in the context of perturbative theories containing elementary Higgs bosons, the MSM need not be nature’s choice. Many generalizations have been discussed,\textsuperscript{[2]} including extensions of the Higgs sector only, extensions of both the gauge and Higgs sectors, and supersymmetric generalizations of all these types of models. Supersymmetric generalizations are particularly attractive in the perturbative context in that they require the presence of elementary spin-zero Higgs fields and solve the well-known naturalness and hierarchy problems. Thus, they provide an enormously attractive theoretical framework in which elementary Higgs bosons must exist. Further,

\textsuperscript{*} The strongly-interacting $W$ scenario is reviewed by T. Han in these proceedings. Technicolor and related scenarios will also not be considered here.
in supersymmetric models, there is always one (or more) light Higgs boson(s) with coupling(s) to the WW channels such that WW scattering remains perturbative at all energies. The most thoroughly investigated model is the Minimal Supersymmetric Model (MSSM) in which the Higgs sector contains two Higgs-doublet fields (the minimum number required in the supersymmetric context), but there is no extension of the gauge or matter sectors other than the introduction of the supersymmetric partner states. In contrast to the single Higgs boson of the MSM, there are five physical Higgs bosons in a two-Higgs-doublet model. Assuming that there is no CP violation in the Higgs sector, they are: the $h^0$, the lightest CP-even mass eigenstate; the $H^0$ the heavier of the two CP-even mass eigenstates; the $A^0$, the single CP-odd state; and a charged Higgs pair, $H^\pm$. The resulting phenomenology is much richer than that of the MSM.

In this review, I will focus on recent developments in understanding how to probe the Higgs sector in three representative cases: the MSM $\phi^0$; the SM with an extended Higgs sector containing two doublet fields, including the possibility of CP violation; and the MSSM. A mix of theoretical and phenomenological issues will prove relevant. By demonstrating that we can thoroughly explore these three quite different cases at a future $e^+e^-$ collider, we will have considerable confidence in our ability to probe any perturbative theory with elementary Higgs bosons that nature may have chosen.

2. The Standard Model Higgs Boson

Let me first review some well-known theoretical ‘facts’. In assessing where we should look for the SM $\phi^0$, the first constraint beyond perturbative unitarity of the WW sector derives from triviality. All current lattice and related investigations appear to require that $m_{\phi^0} \lesssim 650$ GeV if the scale of new physics, $\Lambda$, is to lie above $m_{\phi^0}$.[2] If $\Lambda$ is as large as $\sim 10^{15}$ GeV, then $m_{\phi^0} \lesssim 175$ GeV is required (by the renormalization group equations†) in order that the theory remain perturbative up to the scale $\Lambda$. Finally, in order that the quartic coupling of the Higgs sector not be driven to negative values (implying instability of the potential) by the large Yukawa coupling associated with the heavy top quark, it is necessary that $m_{\phi^0}$ lie above an $m_t^-$ and $\Lambda$–dependent lower bound. For $m_t = 150$ and $\Lambda \sim 10^{15}$ GeV, for instance, $m_{\phi^0} \gtrsim 100$ GeV is required in the context of perturbatively computed renormalization group equations. The lower bound decreases with decreasing $\Lambda$ and/or $m_t$. Nonetheless, it is entirely reasonable that $m_{\phi^0}$ should lie in a range that is somewhat above the current upper limit of $\sim 60$ GeV set by LEP-I, and quite possibly the $\phi^0$ will turn out to be too heavy to be found at LEP-II (which will probe up to $m_{\phi^0} \sim 80 - 90$ GeV for $\sqrt{s} \sim 190 - 200$ GeV). In such a case, an $e^+e^-$ collider of moderate energy would be the ideal machine for detecting the $\phi^0$.

† Of course, these same renormalization group equations have difficulty reproducing the low-energy value of $\sin^2 \theta_W$ in the simplest SU(5) grand unification scheme.
Although the mass of the SM Higgs boson is not known, its couplings to gauge bosons and fermions are completely determined. Thus, the branching ratios and production rates for the $\phi^0$ can be computed as a function of $m_{\phi^0}$. For our immediate purposes, it is only necessary to recall that for $m_{\phi^0} \lesssim 150$ GeV the $\phi^0$ decays primarily to $b\bar{b}$, while for $m_{\phi^0} \gtrsim 150$ GeV the $W^+W^-$ and $ZZ$ decay modes become dominant. (Below the $W^+W^-$ or $ZZ$ threshold, one of the two $W$’s or $Z$’s must be off-shell, but the branching fraction can still be substantial.) At $\sqrt{s} = 500$ GeV, the two main production mechanisms are $e^+e^- \rightarrow Z\phi^0$ and $e^+e^- \rightarrow \nu\overline{\nu}\phi^0$ (where the $\phi^0$ arises from the fusion of $W$’s emitted from the $e^+$ and $e^-$). The former is dominant for $m_{\phi^0} \gtrsim 160$ GeV, while the latter dominates for lower masses. Both detection modes have been thoroughly studied. In the case of the $Z\phi^0$ mode, the most recent studies are those of Refs. 3-5. A review of these results and further refinements are given in the contribution by K. Kawagoe. The conclusion is that at $\sqrt{s} = 500$ GeV, the $\phi^0$ can be discovered up to $m_{\phi^0} \sim 350$ GeV using the recoil missing mass technique for events in which the $Z$ decays to $l^+l^-$ or a tagged $b\bar{b}$ pair. (At $m_{\phi^0} = 350$ GeV, Ref. 3 obtains a net signal event rate of $S = 23$ compared to a net background rate of $B = 15$ for an integrated luminosity of $L = 50$ fb$^{-1}$.) The $\nu\overline{\nu}\phi^0$ mode has also been recently reexamined in Ref. 6, with the conclusion that it is possible to discover the $\phi^0$ via this production process for $m_{\phi^0} \lesssim 300$ GeV using the decay mode $\phi^0 \rightarrow W^+W^- \rightarrow 4$ jets. (At $m_{\phi^0} = 300$ GeV, $S = 30$ and $B = 50$ for $L = 50$ fb$^{-1}$.) This conclusion is substantially consistent with earlier studies.$^{7-9}$ A rough summary and extrapolation to other energies is that $\phi^0$ discovery will be possible up to $m_{\phi^0} \sim 0.7\sqrt{s}$.

It should be noted that rates for both of the above production modes are determined only by the $\phi^0WW$ couplings. Nonetheless, sensitivity to the $\phi^0f\overline{f}$ (where $f$ is a fermion) couplings through the $\phi^0$ branching ratios is significant if $b$ tagging is available, except in the case of $f = t$. This will be discussed shortly. Regarding the $t\overline{t}\phi^0$ coupling, Refs. 10 and 11 claim that the $\phi^0$ will be visible in $e^+e^- \rightarrow t\overline{t}\phi^0$ production for $m_{\phi^0} \lesssim 120$ GeV for $L = 20$ fb$^{-1}$, and probably to somewhat higher mass at $L = 50$ fb$^{-1}$. This production process would then allow a first determination of the $t\overline{t}\phi^0$ coupling for a light $\phi^0$.

In the last few years the possibility of employing collisions of back-scattered laser beams to discover the SM Higgs boson at a linear $e^+e^-$ collider has been explored.$^{12-14}$ The event rate is directly proportional to $\Gamma(\phi^0 \rightarrow \gamma\gamma)$. The interest in this mode derives primarily from two facts. First, observation of the $\phi^0$ in this production mode provides probably the only access to the $\phi^0\gamma\gamma$ coupling at an $e^+e^-$ linear collider. The $\phi^0 \rightarrow \gamma\gamma$ decay channel has (at best) a branching ratio of order $2 \times 10^{-3}$; too few events will be available in direct $e^+e^-$ collisions to allow detection of such decays. Second, in principle it might be possible to detect the $\phi^0$ for $m_{\phi^0}$ somewhat nearer to $\sqrt{s}$ than the $0.7\sqrt{s}$ that appears to be feasible via direct $e^+e^-$ collisions. Indeed, the full $\gamma\gamma$ center-of-mass energy, $W_{\gamma\gamma}$, goes into creating the $\phi^0$, and the back-scattered laser beam facility can be configured so that the $W_{\gamma\gamma}$ spectrum peaks slightly above
Figure 1: The ratio of $\Gamma(\text{Higgs} \to \gamma\gamma)$ computed for two different model choices for a number of cases. In the case of the $\phi^0$, the ratio of the width predicted in the presence of an extra heavy generation to that obtained in the MSM is shown. For the $h^0$, the ratio $\Gamma(h^0 \to \gamma\gamma)/\Gamma(\phi^0 \to \gamma\gamma)$ as a function of $m_{h^0} = m_{\phi^0}$ with $m_{A^0} = 400$ GeV is plotted. Squarks and charginos have been taken to be as light as possible without being observable at the $\sqrt{s} = 500$ GeV collider. For the $H^0$ two curves are shown. The dot-dashed curve is $\Gamma(H^0 \to \gamma\gamma)$ in a model with light charginos ($M = -\mu = 150$ GeV in the notation of Ref. 2) divided by the corresponding width with heavy charginos ($M = -\mu = 1$ TeV), keeping the squarks and sleptons heavy (with masses of order 1 TeV). The dashed curve is $\Gamma(H^0 \to \gamma\gamma)$ in a model with light squarks and sleptons (given by a common soft-SUSY breaking diagonal mass of 150 GeV for all squarks and sleptons, with all off-diagonal mass terms set to zero) divided by the corresponding width computed with heavy squarks and sleptons, keeping the charginos heavy (as specified above). For the latter two curves, the ratio of widths is plotted as a function of $m_{H^0}$ for $\tan \beta = 2$. The top quark mass is taken equal to 150 GeV for all calculations. This figure is taken from Ref. 12.

The importance of determining the $\phi^0\gamma\gamma$ coupling derives from the fact that it is determined by the sum over all 1-loop diagrams containing any charged particle whose mass arises from the Higgs field vacuum expectation value. In particular, the
1-loop contribution of a charged particle with mass \( \gtrsim \frac{m_{\phi^0}}{2} \), approaches a constant value that depends upon whether it is spin-0, spin-1/2, or spin-1. (The contributions are in the ratio \(-1/3 : -4/3 : 7\), respectively.) For a light Higgs boson, in the MSM the dominant contribution is the \( W \)-loop diagram. The next most important contribution is that from the top quark loop, which tends to cancel part of the \( W \)-loop contribution. A fourth fermion generation with both a heavy lepton, \( L \), and a heavy (\( U, D \)) quark doublet would lead to still further cancellation. For \( m_{\phi^0} \gtrsim 2m_W \), the \( W \)-loop contribution decreases, and the heavy family ultimately dominates. To illustrate, we show in Fig. 1 the ratio of \( \Gamma(\phi^0 \to \gamma\gamma) \) as computed in the MSM (with \( m_t = 150 \text{ GeV} \)) to that computed in the presence of an extra generation with \( m_L = 300 \text{ GeV} \) and \( m_U = m_D = 500 \text{ GeV} \). Except for \( m_{\phi^0} \) in the vicinity of 300 GeV, where the full set (mainly the heavy generation) of contributions accidentally matches the MSM result, even a rough measurement (or bound) on the \( \phi^0 \gamma\gamma \) coupling would reveal the presence of the otherwise unobservable heavy generation. It is especially interesting to note that a heavy generation would greatly enhance the event rate (and hence prospects) for detecting a Higgs boson with mass up near \( \sqrt{s} = 500 \text{ GeV} \).

Because of the dominance of the \( W \) loop contribution in the three family case, the \( \phi^0 \gamma\gamma \) coupling is also very sensitive to any deviations of the \( WW\gamma \) and \( WW\phi^0 \) couplings from SM values.\(^{[15,16]} \) The sensitivity to anomalies in these couplings can be substantially greater than that provided by LEP-I data.

How high in mass can the \( \phi^0 \) be detected in \( \gamma\gamma \) collisions in the case of the MSM? For \( \sqrt{s} = 500 \text{ GeV} \), the range of interest is that which cannot be accessed by direct \( e^+e^- \) collisions, \( i.e. \ m_{\phi^0} \gtrsim 350 \text{ GeV} \). In this mass region, \( \phi^0 \to ZZ \) decays provide the best signal. Certainly, the tree-level \( \gamma\gamma \to W^+W^- \) continuum background completely overwhelms the \( \phi^0 \to W^+W^- \) mode. As summarized in Ref. 12, if there were no continuum \( ZZ \) background, and if one of the \( Z \)'s is required to decay to \( l^+l^- \), the event rate would be adequate for \( \phi^0 \) detection up to \( m_{\phi^0} \sim 400 \text{ GeV} \), \( i.e. \ m_{\phi^0} \sim 0.8\sqrt{s} \). Unfortunately, even though there is no tree-level \( ZZ \) continuum background, such a background does arise at one-loop. A full calculation of this background was performed in Ref. 17. The \( W^\pm \) loop is dominant, and leads to a large rate for \( ZZ \) pairs with large mass, when one or both of the \( Z \)'s is transversely polarized. This background is such that \( \phi^0 \) observation in the \( ZZ \) mode is probably not possible for \( m_{\phi^0} \gtrsim 350 \text{ GeV} \), \( i.e. \) no better than what can be achieved in direct \( e^+e^- \) collisions. This result has been confirmed in the recent independent calculation of Ref. 18.

In general, although the \( \gamma\gamma \) mode may not extend the discovery reach of an \( e^+e^- \) collider, it will allow a first measurement of the \( \phi^0 \gamma\gamma \) coupling of any Higgs boson that is found in direct \( e^+e^- \) collisions. The accuracy that can be expected has been studied in Ref. 19. They consider two final states: the \( \phi^0 \to b\bar{b} \) channel with \( b \)-tagging, and the \( \phi^0 \to ZZ \) channel with one \( Z \) required to decay to \( l^+l^- \). In the former case, it is important, as noted in Ref. 12, to polarize the laser beams so that the colliding photons have \( < \lambda_1 \lambda_2 > \) near 1. This suppresses the \( \gamma\gamma \to b\bar{b} \)
background which is proportional to $1 - \langle \lambda_1 \lambda_2 \rangle$. They find that if $35 \lesssim m_{\phi^0} \lesssim 150$ GeV, then the $b\bar{b}$ mode will allow a 5-10\% determination of $\Gamma(\phi^0 \rightarrow \gamma\gamma)$, while for $185 \lesssim m_{\phi^0} \lesssim 300$ GeV the $ZZ$ mode will provide a 8-11\% determination. In the $150 - 185$ GeV window, the $WW$ and $b\bar{b}$ decays are in competition, and the accuracy of the measurement might not be better than 20\%.

Other recent developments in $\phi^0$ physics reported at this conference are in five areas. 1) Measuring details of the $\phi^0$ couplings and verifying its quantum numbers, once the $\phi^0$ has been discovered. 2) Computation of radiative corrections to various production processes and decays. 3) Complications (one of which is mentioned above) in the $\gamma\gamma \rightarrow \phi^0$ discovery mode. 4) Associated production of the $\phi^0$ and other particles in $\gamma\gamma$ collisions. 5) $e\gamma$ collision mode possibilities. I shall make only a few remarks on each. More details can be found in the plenary talks by P. Janot and D. Borden, and in the various parallel contributions to these proceedings.

The value of $b$ tagging for separating different decay channels, $X$, of the $\phi^0$, and thereby determining the product $\sigma_{tot} \times BR(\phi^0 \rightarrow X)$ was demonstrated in Ref. 20. The channels $X = b\bar{b}$, $WW^*$, $c\bar{c} + gg$ and $\tau^+\tau^-$ yield different numbers of displaced vertices on average. By fitting the distribution in the number of displaced vertices, approximate determinations of $\sigma \times BR$ are possible. For $m_{\phi^0} \sim 140$ GeV, the accuracies achieved in the four channels listed above are roughly 12\%, 24\%, 116\% and 22\%, respectively. Whether or not this is a useful level, is somewhat model and situation dependent. For instance, in the model where the Higgs sector of the SM is extended to two-doublets, the couplings of any one of the three neutral Higgs bosons to the various channels can be very different from couplings for the $\phi^0$. If only a single neutral Higgs boson is observed, an experimental demonstration that the couplings to $b\bar{b}$, $\tau^+\tau^-$ and $WW^*$ are within 10-20\% of the values expected for the $\phi^0$ would strongly suggest that the only two-doublet extensions that should be considered are ones in which one of the Higgs bosons automatically has couplings that are close to MSM values. The MSSM supersymmetric two-doublet extension is a case in point. There, the $h^0$ is predicted to have couplings that are not dramatically different from those predicted for the $\phi^0$. Relatively precise determinations could be required to discriminate between the $\phi^0$ and the $h^0$.

Determination of the quantum numbers of a neutral Higgs boson has been another active area. I will defer this discussion to the section on extended Higgs sectors, since it is only by comparison to a model in which there are Higgs bosons with quantum numbers other than those of the $\phi^0$ that one can assess how useful the proposed techniques are. However, I will state my conclusions here. While various production and decay distributions are sensitive to whether the Higgs boson is CP-even or CP-odd, with two exceptions the production or decay processes considered to date will only be visible if the Higgs boson is CP-even, or has a significant CP-even component. In this case, the distributions will be entirely dominated by the CP-even prediction and there will be very little sensitivity to any CP-odd component that might be present. An especially interesting exception to this generality occurs in considering
polarization asymmetries in $\gamma\gamma \rightarrow$ Higgs production using the fact that back-scattered laser beam photons can be given large polarizations.

I will not dwell on the recent progress made in the area of radiative corrections. Details can be found in some of the parallel contributions.\cite{21} Large corrections to $\phi^0 \rightarrow b\bar{b}$ (due primarily to the running QCD mass of the bottom quark) were found early on and lead to $\Gamma(\phi^0 \rightarrow b\bar{b})^{1\text{-loop}} \sim 0.55 \Gamma(\phi^0 \rightarrow b\bar{b})^{\text{tree}}$ once $m_{\phi^0} \gtrsim 60$ GeV. (See Ref. 2, and references therein.) QCD corrections enhance the one-loop width for $\phi^0 \rightarrow gg$ by about 60%,\cite{22,23} whereas corrections to $\Gamma(\phi^0 \rightarrow \gamma\gamma)$ are small if $m_{\phi^0} < 2m_t$\cite{24} and also generally small for $m_{\phi^0} > 2m_t$ except for $m_{\phi^0}$ such that there is a large cancellation between the $W$ and $t$ loops.\cite{25,26} Corrections to the $Z\phi^0$ production process are generally small at a $\sqrt{s} \sim 500$ GeV machine, and, in any case, are exhaustively studied.\cite{21,27} (See also Ref. 28.) Electroweak corrections to $\phi^0$ decays are generally small.\cite{21}

A number of contributions to this conference have explored further backgrounds to detecting the $\phi^0$ in the $\gamma\gamma$ collision mode. In general, I believe the importance of these additional backgrounds has been over emphasized, with the exception of the $\gamma\gamma \rightarrow ZZ$ continuum background from the $W$-loop graphs discussed earlier. The processes $\gamma\gamma \rightarrow Zl^+l^-$ and $Zq\bar{q}$ yield a reducible background to the $ZZ$ mode to the extent that the $q\bar{q}$ or $l^+l^-$ have mass near $m_Z$.\cite{29} The magnitude of this background depends upon the detector resolution. If $\lambda_1\lambda_2 > 1$, then this background can be significant (though not as large as the $ZZ$ continuum background). However, these processes are proportional to $1 - \lambda_1\lambda_2$ and can be suppressed substantially by appropriate polarization choices for the incoming back-scattered laser beams.

The $b\bar{b}$ channel receives a background contribution from “resolved” photon processes.\cite{30} The most important example is that where one incoming $\gamma$ fragments to a spectator jet and a gluon. The subprocess $\gamma g \rightarrow b\bar{b}$ then yields a large $b\bar{b}$ rate. However, this background will not be a problem in practice for two reasons. First, it will be possible to veto against the spectator jet that accompanies the $g$. This probably already reduces the background to a level below the true $\gamma\gamma \rightarrow b\bar{b}$ continuum. Second, for the range of $m_{\phi^0}$ such that the $b\bar{b}$ mode is appropriate ($m_{\phi^0} \lesssim 150$ GeV), the $\phi^0$ will already have been discovered in direct $e^+e^-$ collisions, i.e. $m_{\phi^0}$ will be known. To study the $\phi^0$ in $\gamma\gamma$ collisions it will be easy to adjust the machine energy and laser beam polarizations so that the $\gamma\gamma$ spectrum is peaked at $W_{\gamma\gamma} \sim m_{\phi^0} \sim 0.8\sqrt{s}$. (See, for instance, Ref. 19.) In this case, the secondary gluon in the “resolved”-photon process is quite unlikely to have sufficient energy to create a $b\bar{b}$ pair with mass as large as $m_{\phi^0}$.

A variety of other final states containing the $\phi^0$ can be produced in $\gamma\gamma$ collisions. For instance, the $\gamma\gamma \rightarrow t\bar{t}\phi^0$ analogue of $e^+e^- \rightarrow t\bar{t}\phi^0$ could provide another measure of the $t\bar{t}\phi^0$ coupling.\cite{31,32} However, because of phase space suppression, the rate for this reaction is quite small for $\sqrt{s} = 500$ GeV, and only becomes competitive with
$e^+e^- \to t\bar{t}\phi^0$ when $\sqrt{s} \gtrsim 1$ TeV.\

Turning now to $e\gamma$ collisions, I merely note here that $e\gamma \to W\phi^0\nu \to jjb\bar{b}\nu$ may be viable for $\phi^0$ searches for $\sqrt{s} \gtrsim 1$ TeV. The last reference includes some background studies.) This process is interesting in that it probes the $\gamma W \to \phi^0W$ subprocess which is determined by a combination of graphs with different basic SM couplings. Should the couplings deviate from SM predictions, the large cancellations among the graphs might be reduced and the event rate significantly enhanced. Another mode of interest is $e\gamma \to e\gamma\gamma \to e\phi^0$, in which a secondary $\gamma$ collides with the primary $\gamma$ to create the $\phi^0$. The cross section for this process is bigger than that for $e\gamma \to W\phi^0\nu$ and might allow detection of the $\phi^0$ at $\sqrt{s} = 500$ GeV in the $b\bar{b}$ mode. (Resolved photon backgrounds would have to be suppressed by spectator jet vetoing.)

In summary, it seems clear that the standard $e^+e^- \to Z^* \to Z\phi^0$ production mode has as much, or more, potential for $\phi^0$ discovery as any other process studied to date. Determination of expected couplings will be possible to a reasonable level, but (as discussed in the next section) verification through production or decay distributions that the $\phi^0$ is a pure CP-even state will not be easy. We will only know that if we see it in this production mode, it cannot be pure CP-odd. Of course, if the cross section level is that predicted for the maximally coupled $\phi^0$ of the MSM, then we can be rather certain that the observed Higgs boson is either the $\phi^0$ or a very close approximation thereto. (This latter possibility arises naturally in the MSSM.) $\gamma\gamma$ and $e\gamma$ collisions will primarily be of interest for determining the $\phi^0 \to \gamma\gamma$ coupling, which potentially probes new charged particles at high mass scales. Finally, for $m_{\phi^0} \lesssim 300$ GeV, not only can the $\phi^0$ be detected in $\gamma\gamma$ collisions, but its $\gamma\gamma$ coupling can be determined to rather good accuracy.

Before closing this MSM section, it is useful to make some comparisons to hadron supercolliders. I have no space to justify the statements made regarding hadron colliders. Please refer to Refs. 2 and 37 for sample studies.

1. Discovery: In $e^+e^-$ collisions, discovery will be possible up to roughly $0.7\sqrt{s}$, whereas at the SSC/LHC $80 \lesssim m_{\phi^0} \lesssim 1$ TeV can be probed.

2. $b\bar{b}$ and $t\bar{t}$ couplings: In $e^+e^-$ collisions these couplings can be determined provided the $b\bar{b}$ or $t\bar{t}$ branching ratio of the $\phi^0$ is significant. At the SSC/LHC, only the $t\bar{t}\phi^0 \to t\bar{t}b\bar{b}$ production/decay mode allows direct observation of the $\phi^0$ in its $b\bar{b}$ decay channel. This allows sensitivity to the combination $g_{t\bar{t}\phi^0}^2 BR(\phi^0 \to b\bar{b})$. Large deviations from expectations would be apparent for $m_{\phi^0} \lesssim 120$ GeV, but precise determinations of the couplings would be difficult.

* Radiative corrections to $\gamma\gamma \to t\bar{t}$ and $ZZ$ due to 1-loop Higgs exchange graphs are also sensitive to the $t\bar{t}\phi^0$ coupling. Sufficiently precise measurements of these processes at high luminosity and energy might allow a determination of the coupling over a significant range of $m_t$ and $m_{\phi^0}$ values, assuming no other new physics in the 1-loop graphs.
3. **WW couplings:** The $e^+e^- \to Z^* \to Z\phi^0$ recoil mass discovery mode allows a determination of the $ZZ\phi^0$ coupling that is independent of the Higgs decays. At hadron colliders, if $\sigma(gg \to \phi^0)$ can be accurately computed, and if acceptances etc. of the detector are well-known, then the $gg \to \phi^0 \to ZZ^*$ or $ZZ$ events allow a determination of $BR(\phi^0 \to ZZ^{(*)})$ to reasonable accuracy. This, however, does not directly determine the $ZZ\phi^0$ coupling. Extraction of the coupling would have very limited accuracy except in the $m_{\phi^0}$ region between about 130 and 150 GeV, where the $\phi^0$ could also be observed in the $t\bar{t}b\bar{b}$ production/decay mode mentioned above.

4. **$\tau\tau$ coupling:** At an $e^+e^-$ machine this coupling can be determined wherever the $\tau\tau$ branching ratio is significant. At the SSC/LHC, a means for determining the $\phi^0 \to \tau\tau$ branching ratio has not been convincingly established.

5. **$\gamma\gamma$ coupling:** At an $e^+e^-$ collider, a back-scattered laser beam facility will allow determination of this coupling over essentially the same mass range that observation of the $\phi^0$ in direct $e^+e^-$ collisions is possible. At the SSC/LHC, sensitivity to this coupling will only be significant in the $80 \lesssim m_{\phi^0} \lesssim 150$ GeV mass range, where the $\phi^0$ can be detected in the $gg \to \phi^0 \to \gamma\gamma$ or $gg \to t\bar{t}\phi^0 \to t\bar{t}\gamma\gamma$ production/decay modes.

In short, so long as the $e^+e^-$ machine energy is sufficient for detection of the $\phi^0$, it will be easier to study its couplings than at the SSC or LHC. The primary advantage of the latter machines is their large discovery reach, which greatly exceeds that available at, for instance, a $\sqrt{s} = 500$ GeV $e^+e^-$ collider.

Finally, although we have focused on discovery of the $\phi^0$ and measurement of its basic couplings in the simplest reactions, more complicated final states may ultimately also be of interest. These include such reactions as $e^+e^- \to W^+W^-\phi^0$, $\gamma Z\phi^0$, $ZZ\phi^0$, $\nu eW^\pm\phi^0$, $\nu\nu Z\phi^0$ and $\nu\nu\gamma\phi^0$. These were studied in Refs. 39 and 40. Aside from probing a complicated combination of SM couplings, these processes could also be sensitive to anomalous couplings of the $\phi^0$ to three gauge bosons. However, given the current lower bound on $m_{\phi^0}$ from LEP, adequate event rates for such processes will require either large integrated luminosity or $\sqrt{s} \gtrsim 1$ TeV.

### 3. Extended Higgs Sectors

In this section, I will make some remarks on models in which only the Higgs sector of the Standard Model is extended. The most attractive such extension is to two Higgs doublet fields. For such an extension $\rho = 1$ remains automatic at tree level. The general two-doublet model is also a convenient model for exploring the phenomenology of CP violation in the Higgs sector. If CP is conserved, the two-doublet model yields the five physical Higgs CP-eigenstates enumerated in the introduction: $h^0$, $H^0$, $A^0$, and $H^\pm$. If the Higgs potential violates CP, then in general all the neutral states mix, and there are simply three mass eigenstates of mixed CP character.
Of course higher representations can also be considered, the next most complicated being a Higgs triplet. As is well-known, \( \rho = 1 \) is not automatic in such a case. Various possibilities for obtaining \( \rho = 1 \) at tree level can be entertained. First, it could be that \( m_t \) is very large. In this case the \( t - b \) doublet yields a large positive contribution to \( \Delta \rho \) which could by cancelled by the negative \( \Delta \rho \) that would arise from a \( T = 1, Y = \pm 2 \) complex triplet representation. Obviously, this would require fine tuning the triplet vacuum expectation value. A second possibility is to combine one doublet Higgs field with one real \( \{ T = 1, Y = 0 \} \) and one complex \( \{ T = 1, Y = 2 \} \) field. If the neutral members of the two triplet representations have the same vacuum expectation value, then \( \rho = 1 \) is maintained at tree level. In either case, \( \rho = 1 \) is not maintained at 1-loop. Indeed, unlike the case of doublet models \( \rho \) is infinitely renormalized in triplet models (due to the fact that the interactions of the \( B \) gauge field with the Higgs bosons violate custodial SU(2)). Thus, fine-tuning would be required to maintain \( \rho = 1 \) after renormalization. For this reason, these models are generally not in favor with theorists. However, this should not stop the experimental community from searching for the many new signatures that would arise. At an \( e^+ e^- \) collider, the most spectacular and characteristic signal for a triplet model would be the detection of the doubly charged Higgs boson(s) \( \left( e^+ e^- \rightarrow H^{++} H^{--} \right) \) contained in complex Higgs triplet representation(s). All that is required is adequate machine energy, \( \sqrt{s} \gtrsim 2 m_{H^{++}} \).

Returning to the two-doublet model, let me begin by making a few general remarks. First, the new CLEO limit on the \( b \rightarrow s \gamma \) branching ratio restricts \( m_{H^\pm} \) to be large, \(^{[41-42]} \) in particular, outside the range (roughly \( 2 m_{H^\pm} \lesssim 0.8 \sqrt{s} \)) accessible for \( H^\pm \) detection in \( e^+ e^- \rightarrow H^+ H^- \) at a \( \sqrt{s} = 500 \) GeV collider. \(^{[43]} \) (It is important to emphasize that the limits on \( m_{H^\pm} \) from \( b \rightarrow s \gamma \) decays obtained in Refs. 41 and 42 assume no new physics in the loop graphs other than a single charged Higgs boson.) Second, let us recall that in the CP conserving case the \( Z^* \rightarrow Z h^0 (Z H^0) \) cross section is proportional to \( \sin^2(\beta - \alpha) \) \( \cos^2(\beta - \alpha) \), whereas the \( Z^* \rightarrow h^0 A^0 \) cross section is proportional to \( \cos^2(\beta - \alpha) \). Here, \( \alpha \) is the mixing angle determine by the \( 2 \times 2 \) mass matrix for the CP-even Higgs sector, and \( \tan \beta = v_2 / v_1 \) with \( v_2 \) \( (v_1) \) being the vacuum expectation value of the neutral member of the Higgs doublet that couples to up \((\text{down})\) quarks. The complementarity of the \( Z h^0 \) and \( h^0 A^0 \) reactions has been pointed out many times. \(^{[2]} \) Not both reactions can be suppressed by a weak coupling. However, it is not impossible that \( Z^* \rightarrow Z h^0 \) could be kinematically allowed but coupling suppressed, while \( Z^* \rightarrow h^0 A^0 \) and \( Z^* \rightarrow Z H^0 \) could be kinematically forbidden. Then, no Higgs boson would be seen without raising the machine energy. In a general two-doublet model the masses of the Higgs bosons are all completely independent of one another, the only constraint arising from the requirement that the Higgs sector contribution to \( \Delta \rho \) be small.

In this situation, \( \gamma \gamma \) collisions could play a very vital role. Any neutral Higgs boson can be produced singly in \( \gamma \gamma \) collisions, including the \( A^0 \), whose \( \gamma \gamma \) coupling is determined by fermion loops (with large mass asymptotic limit of 2, compared to the
The cross section for any given Higgs boson depends upon the precise weighting of the different loop diagrams, as determined by the appropriate $\beta$- and $\alpha$-dependent coupling constant factors.\cite{2} Generally speaking, the $\gamma\gamma$ width of all the neutral Higgs bosons of the general two-doublet model can be substantial, and their detection in $\gamma\gamma$ collisions would be possible over a large range of parameter space.

At this point it is useful to return to the issue of whether one can determine the quantum numbers of a neutral Higgs boson (which we generically denote by $h$) that has been detected. Here we focus on the CP quantum number. At least three distinctly different approaches can be identified. First, there are distributions related to the production process $e^+e^- \rightarrow Z^* \rightarrow Zh$.$^{[3,44,45]}$ Second, certain distributions of decay products, e.g. in $h \rightarrow WW \rightarrow f\bar{f}f\bar{f}$, are sensitive to the quantum numbers of the decaying Higgs boson.$^{[46-52]}$ Finally, the dependence of the production rate of a neutral $h$ in $\gamma\gamma$ collisions on the polarization of the colliding photons can be a powerful tool in revealing the $h$ quantum numbers.$^{[53]}$ In order to discuss these techniques, it will be useful to consider a two-Higgs doublet model in a bit more detail. In the most general case, where the Higgs potential includes CP-violating terms, one must diagonalize a $3 \times 3$ mass-squared matrix. There will simply be three mass eigenstates of mixed CP nature. The coupling of the $h$ to the $ZZ$ or $W^+W^-$ channel relative to the coupling of the $\phi^0$ to these channels is given at tree-level by:

$$\frac{g_{WWh}}{g_{WW\phi^0}} = u_2 \sin \beta + u_1 \cos \beta$$

where $u_i$ ($\sum_{i=1,3} |u_i|^2 = 1$) specifies the $h$ mass eigenstate in a certain basis where $i = 1, 2$ refer to CP-even components and $i = 3$ corresponds to a CP-odd component. Note that at tree-level the CP-odd component does not give rise to any $WW$ coupling, and that the net CP-even component coupling of the $h$ will in general be reduced compared to the SM $\phi^0$ value. The coupling of the CP-odd component of the $h$ to the $WW$ channels is very small, arising from one-loop diagrams.

To illustrate the potential, but also the difficulties, of the first two approaches mentioned above, let me consider examples of each. In $e^+e^- \rightarrow Z^* \rightarrow Zh$ production, consider the distribution $d\sigma/d\cos\theta$, where $\theta$ is the angle of the produced $Z$ in the center of mass with respect to the direction of collision of the initial $e^+$ and $e^-$. The distribution takes the form

$$\frac{d\sigma}{d\cos\theta} \propto \begin{cases} \frac{8m_Z^2}{s} + \beta^2 \sin^2 \theta : & \text{CP-even} \\ 1 + \cos^2 \theta : & \text{CP-odd} \end{cases}$$

where $\beta$ is the center-of-mass velocity of the final $Z$. Thus, in principle one can measure this distribution and determine the quantum numbers of the $h$ being produced.
The difficulty with this conclusion is best illustrated by considering a $h$ which is a mixture of CP-even and CP-odd components, as described above. The crucial point is that only the CP-even portion of the $h$ couples at tree-level to $ZZ$, whereas the CP-odd component of the $h$ couples weakly to $ZZ$ via one-loop diagrams. Consequently, the $d\sigma/d\cos\theta$ distribution will reflect only the CP-even component of the $h$, even if the $h$ has a fairly large CP-odd component. The $h$ would have to be almost entirely CP-odd in order for the $\cos\theta$ distribution to deviate significantly towards the CP-odd prediction. However, in this case, the $e^+e^-\rightarrow Z^* \rightarrow Zh$ production rate would be very small, and the $h$ would probably not be detectable in the $Zh$ associated production mode in any case. In summary, any $h$ which is not difficult to detect in the $Zh$ mode will automatically have a $\cos\theta$ distribution that matches the CP-even prediction, even if there is a significant CP-odd component in the $h$. Unfortunately, the $Zh$ production rate itself cannot be used as a measure of the CP-even vs. CP-odd component of the $h$; as noted earlier, even a purely CP-even $h$ can have a $ZZ$ coupling that is suppressed relative to SM strength.

Turning to decay distributions, one encounters a similar problem. In $h \rightarrow WW \rightarrow 4$ fermions, one can determine the angle $\phi$ between the decay planes of the two $W$’s. One finds:

$$\frac{d\sigma}{d\phi} \propto \begin{cases} 1 + \alpha \cos \phi + \beta \cos 2\phi : & \text{CP-even} \\ 1 - (1/4) \cos 2\phi : & \text{CP-odd} \end{cases},$$

where $\alpha$ and $\beta$ depend upon the types of fermions observed and the kinematics of the final state. In general these two distributions are distinguishable. However, in analogy to the previous case, it is almost entirely the CP-even component of the $h$ which will be responsible for its decays to $WW$, and the $\phi$ distribution will thus closely match the CP-even prediction, even if the $h$ has a substantial CP-odd component. This is explicitly verified in the calculations of Ref. 51. In order for the $\phi$ distribution to deviate significantly towards the CP-odd prediction, the CP-even component(s) of the $h$ must be small. Consequently, decays of the $h$ to $WW$ channels will be substantially suppressed, and either the $b\bar{b}$ or $t\bar{t}$ channel will dominate. In the $b\bar{b}$ channel, the CP-even and CP-odd components of the $h$ cannot be separated. The distribution $d\sigma/d\cos\theta^*$ (where $\theta^*$ is the $b$ angle relative to the boost direction in the $h$ rest frame) is predicted to be flat, independent of the CP nature of the $h$.

If $h \rightarrow t\bar{t}$ decays are kinematically allowed, in a general two-Higgs-doublet model they can be important or even dominant relative to the $WW$ decay modes. Further, the CP-even and CP-odd components of the $h$ would couple to the $t\bar{t}$ channel with similar strength. Finally, the $t$ decays are perturbative and the lepton spectra from the $W^\pm$ decay products reflect the spin of the $t$’s. CP-odd asymmetry observables can be constructed using the spin or lepton momenta; a few thousand Higgs boson events might allow observation of an effect if $BR(h \rightarrow t\bar{t})$ is large.\[^{54}\]

The $\gamma\gamma$ coupling of the $h$ is another experimental observable that ‘couples’ with similar strength to CP-even and CP-odd components of the $h$, while simultaneously
revealing how much of each is present. The CP-even component of the $h$ will couple to $\gamma\gamma$ in the fashion of the $\phi^0$, although the relative weights of $W$ and fermion loops can easily be different. In terms of the polarization vectors $\vec{e}_{1,2}$ of the two photons in the photon-photon center of mass, the coupling is proportional to $\vec{e}_1 \cdot \vec{e}_2$. The CP-odd component of the $h$ will also develop a $\gamma\gamma$ coupling at one-loop. As noted earlier, only fermion loops contribute. The coupling is proportional to $(\vec{e}_1 \times \vec{e}_2)_z$ (assuming the photons collide along the $z$ axis). Writing the net coupling as $\vec{e}_1 \cdot \vec{e}_2 E + (\vec{e}_1 \times \vec{e}_2)_z O$, one finds that $E$ and $O$ are naturally of similar size if the CP-odd and CP-even components of the $h$ are comparable ($i.e. |u_1|^2 + |u_2|^2 \sim |u_3|^2$).

In order to reveal the appropriate experimental observables, let us convert to a helicity basis. One finds:

\[
\begin{align*}
|M_{++}|^2 + |M_{--}|^2 &= 2(|E|^2 + |O|^2) \\
|M_{++}|^2 - |M_{--}|^2 &= -4\text{Im}(EO^*) \\
2\text{Re}(M^*_{--}M_{++}) &= 2(|E|^2 - |O|^2) \\
2\text{Im}(M^*_{--}M_{++}) &= -4\text{Re}(EO^*). 
\end{align*}
\] (4)

Each of these helicity amplitude combinations is directly observable via production rate differences obtained by flipping or rotating the polarizations of the colliding photons. If the $h$ has both CP-even and CP-odd components, so that $E$ and $O$ are comparable, then the best asymmetry observable is $A_1 \equiv (|M_{++}|^2 - |M_{--}|^2)/(|M_{++}|^2 + |M_{--}|^2)$. The term in the cross section proportional to $A_1$ changes sign when both $\lambda_1$ and $\lambda_2$ are simultaneously flipped, and is thus best measured by taking $<\lambda_1\lambda_2>$ as near 1 as possible (which suppresses backgrounds anyway), and then comparing the event rate for $<\lambda_1>$ and $<\lambda_2>$ both positive, to that obtained when both are flipped to negative values. Experimentally this is achieved by simultaneously flipping the helicities of both of the initiating back-scattered laser beams. One finds$^{[53]}$ that this asymmetry is typically larger than 10% and is observable for a large range of two-doublet parameter space if CP violation is present in the Higgs potential.

If CP violation is absent, then either $E$ or $O$ will be zero. To distinguish a CP-even from a CP-odd $h$ will then require measurement of the asymmetry $A_3 \equiv 2\text{Re}(M^*_{--}M_{++})/(|M_{++}|^2 + |M_{--}|^2)$. $A_3$ is +1 for a CP-even Higgs boson or −1 for a CP-odd Higgs boson. To extract it experimentally requires that the colliding photons have transverse polarization. For back-scattered laser beams, maximal achievable transverse polarizations are smaller than maximal achievable helicities, and $A_3$ is more difficult to measure than $A_1$. Preliminary results$^{[55]}$ indicate that large (but achievable) luminosities will be required in order to determine the sign of $A_3$.

4. The Higgs Bosons of the Minimal Supersymmetric Model
It is not possible for me to review all the attractive features of the MSSM model here. This subject is reviewed by G. Kane in these proceedings. A particularly noteworthy feature of the model is the fact that, for simple boundary conditions, grand unification in the context of perturbative renormalization/evolution equations yields highly satisfactory values for $\sin^2 \theta_W$ and other precisely measured electroweak parameters. In addition, a common GUT-scale Yukawa coupling yields fermion mass ratios, e.g. $m_b/m_\tau$ that tend to be in close agreement with experiment. Supersymmetry also provides a good candidate for dark matter (the lightest neutralino), and the predicted GUT scale adequately suppresses proton decay.

The Higgs sector of the MSSM is a highly constrained two-doublet Higgs model. Two doublets are required by the basic structure of supersymmetry which makes it impossible to use a single Higgs superfield and its complex conjugate simultaneously in the superpotential construction that is responsible for fermion masses. (Recall that in the MSM the Higgs field yields the down quark masses, while its complex conjugate appears in the Lagrangian term responsible for up quark masses.) Thus, one Higgs superfield has a spin-0 component field that gives mass to down quarks, while the spin-0 component of the other Higgs superfield yields up quark masses. Alternatively, one can also verify that two Higgs superfields are required in order to complete the anomaly cancellations in the supersymmetric context, where superpartners of the Higgs bosons must be present.

In the supersymmetric structure, the quartic couplings of the Higgs fields become related to gauge couplings, and are no longer free parameters. The quadratic mass terms are strongly constrained by minimization conditions, and in the end only two parameters are required to fully specify the Higgs potential. These are normally taken to be $\tan \beta$ and $m_{A^0}$, the mass of the CP-odd scalar. At tree-level, the masses and couplings of all the Higgs bosons can be computed in terms of $m_{A^0}$ and $\tan \beta$. (In particular, the mixing angle $\alpha$ is determined.) Additional parameters are required to determine the one-loop corrections to the Higgs masses, which can be large if $m_t$ and $m_\tilde{t}$ (the stop squark mass) are both large. The basic results are well-known. The two most important points are:

1. $m_{h^0} \leq m_Z + f(\tan \beta, m_t, m_\tilde{t}, \ldots)$, where $\ldots$ refers to generally less important parameters. This upper limit for $m_{h^0}$, which is approached at large $m_{A^0}$, increases with increasing $\tan \beta$, $m_t$ and/or $m_\tilde{t}$.

2. For large $m_{A^0}$, $m_{H^0} \sim m_{H^\pm} \sim m_{A^0}$, $m_{h^0}$ approaches an upper limit, and the couplings of the $h^0$ become rather SM-like. The approach of the $h^0$ couplings to SM-like values is, however, slow enough that important and possibly measurable deviations will be present even for $m_{A^0}$ values above several hundred GeV.

* In the common nomenclature this means that the two-doublet model will be type-II, the same type as implicitly assumed in the previous section.
† The MSSM Higgs potential is such that CP violation in the Higgs sector is not possible at tree-level.
Constraints from existing data on the MSSM Higgs sector are few. Data from LEP-I implies that $m_{A^0} \gtrsim 20$ GeV and $m_{h^0} \gtrsim 40$ GeV. No constraint is currently placed on $\tan \beta$. The $b \to s \gamma$ decay branching ratio limit which is such a powerful constraint in the non-SUSY two-doublet context is considerably weaker. In the strict supersymmetric limit, $BR(b \to s\gamma) = 0$. Not surprisingly, there is then a large region of SUSY parameter space such that there is no inconsistency with current upper limits.\cite{57-60}

The first issue of importance is to specify the range of parameter space for which LEP-II will be able to detect the $h^0$. This has received much attention in the literature, and it is not possible to cite all references here. The graph I present, Fig. 2, is taken from Ref. 61. A more experimentally oriented discussion is given in Ref. 4. Figure 2 shows the discovery boundaries in $m_t - m_{\tilde{t}}$ parameter space beyond which $h^0$ detection (at large $m_{A^0}$) would be impossible, as a function of the ($\sqrt{s}, \tan \beta$) values chosen. The following observations are perhaps useful:

1. If $m_t = 180$ GeV, $\sqrt{s} \gtrsim 240$ GeV is required for detection to be possible at large $m_{\tilde{t}}$ and large $\tan \beta$.

2. If $m_t \lesssim 140$ GeV, $\sqrt{s} \gtrsim 200$ GeV yields good coverage, but $\sqrt{s} \lesssim 190$ GeV would not. For instance, if $m_{\tilde{t}} \sim 350$ GeV, a value representative of that found in some grand unification scenarios, $\sqrt{s} \sim 200$ GeV would cover all $m_t$ and $\tan \beta$ possibilities (within reason), whereas $\sqrt{s} \sim 190$ GeV would leave a region at large $\tan \beta$ in which $h^0$ detection would not be possible.

If the $h^0$ is not found at LEP-II because the machine energy is not pushed to the required level, then one must rely on future colliders to probe the MSSM Higgs sector. An $e^+e^-$ linear collider could prove to be an ideal machine for this purpose. A number of recent studies have appeared from both a theoretical\cite{62-64} and an experimental perspective.\cite{4} See also Ref. 61. Certainly, any $e^+e^-$ collider with $\sqrt{s} \gtrsim 300$ GeV is guaranteed to find the $h^0$. At $\sqrt{s} = 500$ GeV, if $m_{A^0}$ is large enough that the $h^0$ has SM-like couplings, then the $WW \to h^0$ fusion process yields the largest rate for $h^0$ production. But $Z^* \to Zh^0$ also yields an entirely adequate event rate for $h^0$ discovery. In Fig. 3, the 30 and 100 event number contours for these and four other basic production processes are displayed in $m_{A^0}$–$\tan \beta$ parameter space. Typically, 100 events (before cuts and branching ratios) should be adequate for observation of any given process. We see that it is only in the large $\tan \beta$, $m_{A^0} \lesssim 100$ GeV corner of parameter space (where the $WWh^0$ coupling – proportional to $\sin^2(\beta - \alpha)$ – is small) that one must turn to the alternative, but perfectly adequate $Z^* \to h^0 A^0$ production mode. Detection of the $e^+e^- \to h^0 t\bar{t}$ process is likely to be possible whenever the $h^0$ has roughly SM-like couplings, i.e. when $m_{h^0}$ is near its upper limit.\cite{11}

The most important limitation of a $e^+e^-$ collider is also apparent from Fig. 3. We see that if $m_{A^0} \gtrsim 100$ GeV, then the only possible means for detecting the $H^0$ and $A^0$ is via the pair production process, $Z^* \to H^0 A^0$. However, the parameter range for which this process has adequate event rate is limited by the machine energy...
Figure 2: The boundaries in $m_t$ and $m_{\tau}$ parameter space, beyond which the $h^0$ cannot be detected in $Z^* \rightarrow Zh^0$ at LEP-II, are plotted. We assume that $L = 500 \text{ pb}^{-1}$ and require 100 total $h^0Z$ events (before efficiencies). Each curve corresponds to a different choice for the machine energy and $\tan \beta$, as indicated by the $(\sqrt{s}, \tan \beta)$ labellings. $m_{A^0}$ has been chosen to be 1 TeV, so that $m_{h^0}$ is near its upper limit. We have neglected squark mixing. The region where discovery is possible lies to the left and below the boundary curves.

to $m_{A^0} \sim m_{H^0} \lesssim \sqrt{s}/2 - 30 \text{ GeV}$ (recall that $m_{H^0} \sim m_{A^0}$ at large $m_{A^0}$). At $\sqrt{s} = 500 \text{ GeV}$, this means $m_{A^0} \lesssim 210 \text{ GeV}$. Meanwhile, $e^+ e^- \rightarrow H^+ H^-$ is also limited to $m_{H^\pm} \sim m_{A^0} \lesssim 210 - 230 \text{ GeV}$.\[43\] Thus, it could happen that only a rather SM-like $h^0$ is detected in $e^+ e^-$ collisions at the linear collider, and none of the other Higgs bosons are observed.
Figure 3: Event number contours for an NLC with $\sqrt{s} = 500$ GeV and $L = 10$ fb$^{-1}$. We have taken $m_t = 150$ GeV, $m_{\tilde{t}} = 1$ TeV and have neglected squark mixing. We display contours for 30 and 100 total events (no detection efficiencies included) for the six most important processes at an NLC: $Z^* \rightarrow Zh^0, ZH^0, h^0A^0, H^0A^0$, and $WW \rightarrow h^0, H^0$. Generally speaking, 100 events ensure that the reaction can be observed.

However, $\gamma\gamma$ collisions using back-scattered laser beams might allow discovery of the $H^0$ and/or $A^0$ up to higher masses.$^{[12]}$ And, detection of the $h^0$ in $\gamma\gamma$ collisions
is relatively certain to be possible. Observation of any of the three neutral Higgs bosons would constitute a measurement of the $\gamma\gamma$ Higgs coupling, which in principle is sensitive to loops involving other charged supersymmetric particles such as squarks and charginos. Fig. 4 illustrates the discovery potential for the $H^0$ and $A^0$ at $m_t = 150$ GeV in various final state channels. (Superpartner masses have been taken to be large and machine energy is assumed to be about 20% higher than the Higgs mass.) Particularly interesting channels at moderate $\tan\beta$ and below $t\bar{t}$ threshold are $H^0 \rightarrow h^0h^0$ (leading to a final state containing 4 $b$ quarks) and $A^0 \rightarrow Zh^0$. These channels are virtually background free unless $m_{h^0} \sim m_W$, in which case $b$ tagging would be needed to eliminate the large $\gamma\gamma \rightarrow W^+W^-$ continuum background process. These plots are machine energy independent; $\sqrt{s}$ would have to be about 20% higher than the Higgs mass.

* Of course, since charged Higgs bosons can only be pair produced in $\gamma\gamma$ collisions, $e^+e^-$ collisions will yield the greatest kinematical reach in $m_{H^\pm}$. For a study of the $\gamma\gamma \rightarrow H^+H^-$ process see Ref. 65.
$\gamma\gamma \to W^+W^-$ continuum background would have to be eliminated by $b$-tagging. Above $t\bar{t}$ threshold, $H^0, A^0 \to t\bar{t}$ decays dominate (at moderate $\tan\beta$). We see that the event rate is high and that the $\gamma\gamma \to t\bar{t}$ continuum background is of the same general size as the signal rate. Discovery of the $A^0$ and $H^0$ up to roughly $0.8\sqrt{s}$ would be possible.

For large $\tan\beta$, it is necessary to look for the $A^0$ and $H^0$ in the $b\bar{b}$ final state. For the effective integrated luminosity chosen, $L = 20$ fb$^{-1}$, Fig. 4 shows that detection will be difficult except at low masses, in particular masses such that $Z^* \to H^0A^0$ would be observable in $e^+e^-$ collisions. However, it is technically feasible (although quite power intensive) to run the $\gamma\gamma$ collider at very high instantaneous luminosity such that accumulated effective luminosities as high as 200 fb$^{-1}$ can be considered. In this case, detection of the $A^0$ and $H^0$ in the $b\bar{b}$ channel should be possible for masses $\lesssim 0.8\sqrt{s}$.

A particularly interesting question is the extent to which the $\gamma\gamma$ widths (that would be measured by detection of the $h^0, H^0, A^0$ in $\gamma\gamma$ collisions) depend upon the the SUSY context and/or superpartner masses. Some exploration of this issue has appeared in Refs. 12 and 67. Potentially, these widths are sensitive to loops containing heavy charged particles. However, it must be recalled that supersymmetry decouples when the SUSY scale is large. (In particular, superpartner masses come primarily from soft SUSY-breaking terms in the Lagrangian and not from the Higgs field vacuum expectation value(s).) Several cases are illustrated in Fig. 1. First, suppose the lightest CP-even Higgs boson has been discovered, but that no experimental evidence for either the heavier Higgs bosons or any supersymmetric particles has been found. Could a measurement of the $h^0\gamma\gamma$ coupling provide indirect evidence for physics beyond the SM? Figure 1 shows that if the MSSM parameters are chosen such that all new particles beyond the SM are too heavy to be produced (technically we take $M = -\mu = 300$ GeV for the charginos and a common squark/slepton diagonal mass of 300 GeV), then the deviation of $\Gamma(h^0 \to \gamma\gamma)$ from the corresponding SM value for the $\phi^0$ is less than 15%. This is because of decoupling; as the SUSY breaking scale and the scale of the heavier Higgs bosons become large, all couplings of the $h^0$ approach their SM values and the squark and chargino loops become negligible. Even with the MSSM parameters chosen such that the supersymmetric partners lie only just beyond the reach of a $\sqrt{s} = 500$ GeV $e^+e^-$ collider, it will not be easy to distinguish the $h^0 \to \gamma\gamma$ decay width from that of the $\phi^0$. Ref. 19 claims measurement accuracies for the $\gamma\gamma$ width of a SM-like Higgs boson of order 10%, i.e. just on the borderline of what is required.

On the other hand, suppose that the $H^0$ or $A^0$ is light enough to be seen in $\gamma\gamma$ collisions. In this case, a measurement of its $\gamma\gamma$ coupling can provide useful information on the spectrum of charged supersymmetric particles (even if the latter are too heavy to be directly produced). Fig. 1 provides two examples in the case of the $H^0$. Large alterations in the $\gamma\gamma$ width occur if either (or both) the chargino mass scale or the squark mass scale is taken to be significantly below 1 TeV. In the
case of the $A^0$, squark loop contributions to the $A^0\gamma\gamma$ coupling are absent. However, the sensitivity of this coupling to the chargino loops is similar to that of the $H^0\gamma\gamma$ coupling.

Of course, if only one light Higgs boson is discovered, its $\gamma\gamma$ coupling is only one among many that might reveal that the Higgs boson is the $h^0$ of MSSM rather than the $\phi^0$. Hildreth\cite{20} shows that $b$ tagging can separate the various $h^0$ decay channels sufficiently that the $h^0-\phi^0$ distinction can be made if $\tan\beta \gtrsim 6$. The coupling that deviates most from the SM value when $\tan\beta$ is large is the $WW$ coupling, but the (harder to measure) $c\bar{c}+gg$ channel ($c\bar{c}$ and $gg$ are combined since they have similar displaced vertex rates) also has large deviations from SM expectations.

In another contribution to this conference, Kurihara\cite{68} has claimed that at high integrated luminosity ($L \sim 150 \text{ fb}^{-1}$) the measurement of $\sigma(e^+e^- \to Z h^0) \times BR(h^0 \to b\bar{b})$ can be made with such precision (2%) that the distinction between the $h^0$ and the $\phi^0$ would be possible over most of parameter space (for $m_{A^0} \lesssim 600 \text{ GeV}$). However, there are systematic uncertainties, such as our inability to precisely determine the $b$ quark mass. Further study of the influence of such systematic problems appears to be necessary before drawing any firm conclusions.

Once again, it is appropriate to compare the ability of an $e^+e^-$ collider to that of the SSC/LHC hadron colliders to probe the MSSM Higgs sector. For reviews of the results, see Ref. 61 and the talk by A. Rubbia in these proceedings. If only Higgs$\rightarrow\gamma\gamma$ (possibly in association with $t\bar{t}$ production), Higgs$\rightarrow ZZ(*) \rightarrow 4\ell$ and $t \rightarrow H^+b$ channels are deemed sufficiently clean to be viable at a hadron collider, then it is possible for there to be a window in $m_{A^0}-\tan\beta$ parameter space for which no MSSM Higgs boson is observed at either LEP-200 or at the SSC/LHC. This parameter space hole, located in the range $110 \lesssim m_{A^0} \lesssim 170 \text{ GeV}$ and $\tan\beta \gtrsim 7-10$, only arises if $m_t$ is in the vicinity of $130-170 \text{ GeV}$, if $m_{\tilde{t}} \sim 1 \text{ TeV}$, and if no other detection modes are viable. If $m_t$ is either smaller ($\sim 100 \text{ GeV}$) or larger ($\sim 200 \text{ GeV}$), or if $m_{\tilde{t}}$ is smaller ($\lesssim 400 \text{ GeV}$) then the hole disappears even without considering additional modes.

One additional mode that has been investigated is $H^0, A^0 \rightarrow \tau^+\tau^-$. It becomes viable at large $\tan\beta$, where the $gg \rightarrow b\bar{b}+A^0, H^0$ production rates are greatly enhanced and $b\bar{b}$ decays have $\sim 90\%$ branching ratio. In the L3P simulation study reviewed by Rubbia, $A^0$ and $H^0$ detection in this mode is viable for all $m_{A^0} \gtrsim 100 \text{ GeV}$ and $\tan\beta \gtrsim 7$. In particular, this mode is unique in that rather heavy $A^0$ and $H^0$ Higgs bosons can be seen if $\tan\beta$ is larger enough.

A second mode that has just recently received attention is associated Higgs$+t\bar{t}$ production with Higgs$\rightarrow b\bar{b}$. Based on the analysis of the SM Higgs boson,\cite{38} after correcting for branching ratio and production rate differences, Ref. 69 concludes that if $m_t \sim 150 \text{ GeV}$, then detection of the $h^0$ in this mode will be possible for any $m_{A^0} \gtrsim 110 \text{ GeV}$; $H^0$ detection in this mode is possible for $m_{A^0} \gtrsim 50 \text{ GeV}$ up to
the lower limit in $m_{A^0}$ for $h^0$ detection. That is, either the $H^0$ or the $h^0$ can be detected in this way for $m_{A^0} > 50$ GeV. This completely closes the parameter space hole. In fact, by combining just the $t \rightarrow H^+b$ detection mode and this $t\bar{t}b\bar{b}$ final state mode, detection of at least one MSSM Higgs boson is guaranteed to be possible at the SSC/LHC alone. At $m_t \sim 200$ GeV, the parameter space region for which the $t\bar{t}b\bar{b}$ mode is viable (for either $h^0$ or $H^0$) is smaller; but this is simply correlated with the fact that other decay modes, most notably the $ZZ^{(*)} \rightarrow 4\ell$ mode, of the $H^0$ and (now rather heavy) $h^0$ acquire larger branching ratios, and become viable over a large range of parameter space. Of course, it should be noted that this $t\bar{t}b\bar{b}$ mode is not viable for actually observing a heavy $H^0$ or heavy $A^0$; at large $\tan\beta$ the $t\tilde{t}A^0$ and $t\bar{t}H^0$ production processes are suppressed, while at small $\tan\beta$, the $A^0$ and $H^0$ decays would be dominated by the $t\bar{t}$ final state.

To reiterate, combining all modes, the SSC/LHC alone will detect at least one and most probably several of the MSSM Higgs bosons. In addition, it is not unlikely that the $h^0$ can be detected in all three of its most crucial decay channels, $ZZ^*$, $b\bar{b}$, and $\gamma\gamma$, simultaneously. Further, detection of rather heavy $H^0$ and $A^0$ will be possible if $\tan\beta$ is large. An $e^+e^-$ collider would undoubtedly do a better job of determining all the $h^0$ couplings, but the $A^0$, $H^0$ and $H^\pm$ could be beyond its kinematical reach.

The above discussion does not take into account possible supersymmetric decay channels for the MSSM Higgs bosons. When allowed, $\tilde{\chi}\tilde{\chi}$ (where $\tilde{\chi}$ represents a chargino or neutralino) decay modes of the MSSM Higgs are substantial, and often dominant. Some work on this subject at the SSC/LHC has appeared in Refs. 70 and 61. The kinematic reach of a $\sqrt{s} \sim 500$ GeV $e^+e^-$ collider is such that these modes are not likely to be relevant unless the lightest chargino is seen at LEP-200.\[^{63}\]

Of course, in more specific grand unification, renormalization group scenarios, rather restricted predictions for the SUSY parameters emerge. This very active field is reviewed in these proceedings in the contributions by G. Kane, V. Barger, S. Pokorski, and L. Roszkowski. Very often the resulting $m_t$ value is in the $130-180$ GeV range, and $\tan\beta$ lies between 2 and 10. Squark masses are often of rather moderate size, implying that the $h^0$ is relatively light, $m_{h^0} \lesssim 110$ GeV.\[^{\dagger}\]$ But, the $H^0$ and $A^0$ are generally in the mass range above 220 GeV (i.e. just beyond the reach of $e^+e^-$ collisions at a $\sqrt{s} = 500$ GeV collider). Neutralinos and charginos are light enough that SUSY decay modes of the $A^0$ and $H^0$ would be important. Generically speaking, the lightest chargino is typically light enough to be seen at LEP-200. If SUSY really lives at this low a mass scale, we shall soon know it! However, it should be emphasized that all these renormalization group/GUT investigations make rather

\[^{\star}\]$ There is a crossover in the vicinity of $m_{A^0} \sim 110$ GeV where the $h^0$ and $H^0$ interchange roles as being SM-like.

\[^{\dagger}\]$ Of course, $m_{\tilde{t}}$ cannot be light if $\tan\beta$ is near 1 without violating current limits from LEP-I on a light Higgs boson.
specific assumptions (e.g. limits on fine-tuning, bottom-τ Yukawa unification, etc.). It could still be that nature chooses a SUSY scale nearer 1 TeV.

What if one goes beyond the MSSM? As reviewed by G. Kane for these proceedings, in supersymmetric models there will always be a light Higgs boson with mass below roughly 150 GeV, assuming the absence of additional new physics between the scale of supersymmetry breaking and \( \sim 10^{16} \text{ GeV} \). (In fact, G. Kane argues for a much lower upper limit in most cases.) Thus, the lightest Higgs boson of the model will always be kinematically accessible at a \( \sqrt{s} \sim 500 \text{ GeV} \) \( e^+e^- \) machine. However, there are often quite a few CP-even Higgs eigenstates, and it could happen that these share the \( WW \) coupling in such a way that none have a large production event rate. This was investigated recently in Ref. 71, using the model in which there is one additional Higgs singlet field.\[^{71}\] For a significant range of acceptable renormalization group solutions, all the CP-even Higgs bosons are light (in agreement with Ref. 72) but none would have adequate \( WW \) coupling for discovery.

5. Conclusions

Overall, it is clear that a \( \sqrt{s} \gtrsim 500 \text{ GeV} \) \( e^+e^- \) collider has great potential for Higgs boson discovery and, especially, detailed study. In both the SM and the MSSM model, light Higgs boson(s) should be regarded as likely, and any such light Higgs is almost certain to be discoverable at an \( e^+e^- \) machine. The principle limitation of an \( e^+e^- \) collider relative to the SSC/LHC hadron machines is kinematic. But, the couplings and branching ratios of any Higgs boson that can be detected at the \( e^+e^- \) collider can be more accurately investigated in detail than can those of a Higgs boson found at the SSC/LHC. The importance of implementing the back-scattered laser beam technique for probing the Higgs sector cannot be overemphasized. Not only could it lead to increased kinematic reach for the \( e^+e^- \) collider, but it will allow a measurement of the very interesting (as a probe of new physics) \( \gamma\gamma \) coupling of any Higgs boson seen, and may allow a determination of its CP properties.

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