PROGRESS IN SSC HIGGS PHYSICS:
REPORT OF THE HIGGS WORKING GROUP*

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Abstract
I review new developments in Higgs physics and electroweak symmetry breaking that have resulted from the Madison–Argonne workshops on SSC physics.

1. Introduction

The most fundamental mission of the SSC will be to reveal the nature and source of electroweak symmetry breaking (EWSB). Among the many models that have been developed for the EWSB sector, the minimal Standard Model (MSM) and the minimal supersymmetric extension of the Standard Model (MSSM) are certainly the simplest examples of two quite different classes of theory, and it is upon these that our working group focused. In any case, the issues that arise in developing techniques for detecting the single Higgs boson (\(\phi^0\)) of the MSM or the family of Higgs bosons that emerge in the MSSM are representative of those that must be faced in any model.

In the case of the SM \(\phi^0\), it is well known that discovery is most challenging if the Higgs is either light or very heavy. Much of the working group focus was on the so-called ‘intermediate mass’ region, \(80 \lesssim m_{\phi^0} \lesssim 2m_Z\). Perhaps it is worth reviewing the motivation for such a focus. First, all current lattice and related investigations

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appear to require that $m_{\phi^0} \lesssim 650 \text{ GeV}$ if the scale of new physics $\Lambda_{\text{NewPhysics}}$ is to lie above $m_{\phi^0}$. If $\Lambda_{\text{NewPhysics}}$ is as heavy as $\sim 10^{15} \text{ GeV}$, then $m_{\phi^0} \lesssim 200 \text{ GeV}$ is required (by the renormalization group equations) in order that the theory remain perturbative up to the scale $\Lambda_{\text{NewPhysics}}$. 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 $\Lambda_{\text{NewPhysics}}$ dependent lower bound. For $m_t = 150$ and $\Lambda_{\text{NewPhysics}} \sim 10^{15} \text{ GeV}$, for instance, $m_{\phi^0} > \sim 200 \text{ GeV}$ is required (by the renormalization group equations) in order that the theory remain perturbative up to the scale $\Lambda_{\text{NewPhysics}}$.

Several techniques for detecting a light $\phi^0$ have been shown to be very promising. For $80 \lesssim m_{\phi^0} \lesssim 140 \text{ GeV}$, the $\phi^0$ should be visible in $gg \to \phi^0 \to \gamma\gamma$ at those detectors with excellent ($\sim 1\%$) $\gamma\gamma$ mass resolution, and in the $W\phi^0 + t\bar{t}\phi^0 \to \ell\gamma\gamma X$ mode at all detectors (with at least $\sim 3\%$ mass resolution). These channels were critically examined at the workshops, and technical improvements and uncertainties/concerns will be reviewed here and in associated contributed papers. For $130 \lesssim m_{\phi^0} \lesssim 700 - 800 \text{ GeV}$, an extremely clean signal will be available in $gg \to \phi^0 \to ZZ (\ast) \to 4\ell$. The viability of the $4\ell$ channel seems unquestionable in the indicated mass range. For $m_{\phi^0}$ in the 500-800 GeV range, $gg \to \phi^0 \to ZZ \to \ell^+\ell^-\nu\bar{\nu}$ also provides a viable signal; a possible means for extending this signal to $m_{\phi^0} \gtrsim 800 \text{ GeV}$ was examined during the workshops, and will be outlined. Missing from the above list is any means for detecting a light $\phi^0$ in its primary $b\bar{b}$ decay mode. During the workshops, it was shown that expected $b$-tagging efficiency and purity should be adequate to isolate an $\phi^0$ signal for $m_{\phi^0} \lesssim 110 - 120 \text{ GeV}$ at the SSC/LHC in $t\bar{t}\phi^0$ production followed by $\phi^0 \to b\bar{b}$ decay. In addition, the possibility of detecting $W\phi^0$ production followed by $\phi^0 \to b\bar{b}$ at the Tevatron (for $m_{\phi^0}$ below about $90 \text{ GeV}$) was explored. These possibilities will also be reviewed.

At the very high end of the mass scale, there has been recent work showing that we should be able to explore EWSB even if the $\phi^0$ is heavier than 1 TeV (or non-existent) and the interactions of (longitudinally polarized) $W$ ($W \equiv W^\pm$, $Z$) bosons become non-perturbative at high energies. This work will not be covered in detail in this report. It is reviewed in these proceedings in the contributions by J. Bagger, K. Cheung, and D. Morris. A few brief remarks will appear later.

Of course, even in the context of perturbative theories containing elementary Higgs bosons, the MSM need not be nature’s choice. Many generalizations have been

† 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.
discussed, 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, 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. For simple boundary conditions, grand unification in the MSSM 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. The MSSM 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 various gauge bosons are present. Thus, there are five physical Higgs bosons in the MSSM. 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 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 $\tan \beta$ (where $\tan \beta = v_2/v_1$ is the ratio of vacuum expectation values for the

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* In the common nomenclature, this means that the two-doublet model will be type-II.
† The MSSM Higgs potential is such that CP violation in the Higgs sector is not possible.
two doublets) 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 $\alpha$, arising in the diagonalization of the neutral CP-even mass matrix, 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. $m_{h^0} \leq m_Z + f(m_{A^0}, \tan \beta, m_t, m_{\tilde{t}}, \ldots)$, where $\ldots$ refers to generally less important parameters.

2. For large $m_{A^0}$, $m_{H^0} \sim m_{H^\pm} \sim m_{A^0}$, $m_{h^0}$ approaches an upper limit (which increases with increasing $\tan \beta$, $m_t$ and/or $m_{\tilde{t}}$), 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.

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, even when $m_{A^0}$ is relatively small (leading to $m_{H^+}$ near its lower bound of $\sim m_W$).

In the last few years, it has been demonstrated that detection of at least one of the MSSM Higgs bosons should be possible at the SSC, for almost all choices of model parameters. This is often referred to as the ‘no-lose’ theorem for the MSSM Higgs sector. For a review see, for example, Refs. 8, 9 and 10. To briefly summarize, we first note that this statement relies only upon employing the $gg \to h \to ZZ^{(s)} \to 4\ell$, $gg \to h \to \gamma \gamma$ (and/or $Wh + \tilde{t} h \to \ell \gamma \gamma X$) and $t \to H^+ b$ production/decay modes (here $h$ refers to $H^0$ or $h^0$). A region of parameter space that is not covered by these modes arises if $m_t \sim 150$ GeV and $m_{\tilde{t}}$ is large (e.g. $m_{\tilde{t}} \sim 1$ TeV); it comprises $110 \lesssim m_{A^0} \lesssim 160$ GeV and $\tan \beta \gtrsim 4$. It should also be remarked that the $A^0$ is generally not observable at the SSC in these modes (assuming $\tan \beta \gtrsim 1$), and that the $H^0$ can only be detected in small regions of parameter space. For instance, $H^0 \to 4\ell$ is detectable (at $m_t = 150$ GeV and $m_{\tilde{t}} = 1$ TeV) in a tear-drop shaped region roughly located in the region $50 \lesssim m_{A^0} \lesssim 2m_t$ and $\tan \beta \lesssim 5$. For $m_t \gtrsim 180$ GeV and $m_{\tilde{t}} \sim 1$ TeV, the region over which $H^0 \to 4\ell$ is viable expands greatly, and the SSC alone should be able to discover one of the MSSM Higgs bosons.

‡ If LEP-200 is considered as well, then the region over which no MSSM Higgs boson is found at either LEP-200 or the SSC is restricted to the same $m_{A^0}$ range, but $\tan \beta \gtrsim 10$. 

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Clearly, our MSM discussion of the technical improvements, criticisms and uncertainties for the $\gamma\gamma$ and $\ell\gamma\gamma$ modes will be highly relevant to the MSSM Higgs discovery possibilities. However, additional considerations enter in the MSSM extension, and these will be reviewed. The impact of the $gg \rightarrow t\bar{t}h \rightarrow t\bar{t}b\bar{b}$ detection channel upon MSSM Higgs boson discovery will also be outlined. In particular, expected $b$-tagging efficiency and purity should allow us to detect $h^0 \rightarrow b\bar{b}$ in the above-noted window of parameter space where the $\gamma\gamma$ (or $\ell\gamma\gamma$) and $4\ell$ modes are not adequate to detect one of the neutral MSSM Higgs bosons.

It should be noted that the no-lose theorem relies primarily on the detection of relatively light Higgs bosons. Over much of parameter space, only the $h^0$ is observed in the above-mentioned modes. An implicit assumption is that the $h^0$ does not decay to supersymmetric particle final states. While this is the most likely situation, there do exist scenarios in which the lightest neutralino ($\tilde{\chi}^0_1$) is sufficiently light that $h^0 \rightarrow \tilde{\chi}^0_1\tilde{\chi}^0_1$ is kinematically allowed, in which case this could be the dominant $h^0$ decay mode. In such an instance, $h^0$ decays would be invisible. Thus, means for detecting an invisibly decaying Higgs boson were developed, and will be summarized.

But what about the heavier Higgs bosons of the MSSM? If $m_{A^0} \gtrsim 2m_Z$, the above detection modes may not allow us to observe any of the other (approximately degenerate) Higgs bosons: the $A^0$, $H^0$ and $H^+$. Since the most characteristic signature for a two-doublet model (in general) is the existence of a charged Higgs boson, the possibility of $H^+$ detection in the $gg \rightarrow H^+b\bar{t} \rightarrow t\bar{t}b\bar{b}$ final state was explored. For a significant range of parameter space, it appears that $b$-tagging efficiency and purity should prove adequate to detect the $H^+$ in this mode. Preliminary results will be summarized.

Of course, still more exotic models for the Higgs sector can be constructed. Several popular extensions contain at least one neutral Higgs boson that could have extremely weak coupling to fermion, but SM-like couplings to gauge bosons. The signals for such a Higgs boson and its associated partners will be outlined.

More detailed reviews of the progress that has been made prior to the Madison–Argonne workshops appears in Ref. 1, and in the various Snowmass SSC workshop proceedings and LHC proceedings. Where relevant, we will reference the appropriate material. Here we will give an overview of what has been accomplished during the workshops, referring the reader for details to various individual reports and recent preprints/publications. In particular, we hope to establish the motivation for the projects that were pursued, and put into proper context the progress made.

2. The Standard Model

As noted earlier, substantial emphasis during these workshops was placed on improving our ability to probe a relatively light $\phi^0$. Since it is not at all unlikely that $m_{\phi^0}$ could fall in the range where the most established detection modes rely on $\phi^0 \rightarrow \gamma\gamma$ decays, it is the $\gamma\gamma$ and $\ell\gamma\gamma$ modes upon which we first focus.
2.1. The $\gamma\gamma$ and $\ell\gamma\gamma$ Detection Modes

The crucial $\phi^0\gamma\gamma$ coupling derives from the sum over all 1-loop diagrams containing any charged particle whose mass arises from the Higgs field vacuum expectation value, and is thus sensitive to many types of new physics. In particular, the 1-loop
contribution of a charged particle with mass \( \gtrsim m_{\phi^0}/2 \), 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 \( m_U = m_D = 500 \text{ GeV} \). For \( m_{\phi^0} \) below 150 GeV, we see that the \( \phi^0\gamma\gamma \) coupling would be substantially below the MSM result. However, at the same time the cross section for \( gg \to \phi^0 \) via heavy quark loops would be greatly enhanced. To first approximation, the net event rate for the inclusive \( \gamma\gamma \) mode remains approximately unaltered, and we
would still be able to detect the $\phi^0$ in the inclusive $\gamma\gamma$ mode.\textsuperscript{12} This is illustrated in Fig. 2. There we compare the value obtained for $\Gamma(\phi^0 \rightarrow gg) \times BR(\phi^0 \rightarrow \gamma\gamma)$ (which determines the number of inclusive $\gamma\gamma$ events seen when the Higgs is produced via $gg$ fusion, followed by decay to $\gamma\gamma$) in the presence of an extra generation to that obtained in the MSM. We see that over the $50 \lesssim m_{\phi^0} \lesssim 130$ GeV range, where the inclusive $\gamma\gamma$ mode is potentially viable, the extra generation result is between 1.1 and 1.5 times as large as the MSSM result.

On the other hand, cross sections for the associated production $W\phi^0$ and $t\bar{t}\phi^0$ processes, responsible for the $\ell\gamma\gamma$ channel signal, would not be enhanced due to the presence of an extra generation. The reduced $\phi^0 \rightarrow \gamma\gamma$ coupling and branching ratio would then cause the Higgs signal in this channel to be substantially reduced, and perhaps not even visible. Thus, the observation of a Higgs signal in the inclusive $\gamma\gamma$ mode, in combination with a reduced or absent signal in the $\ell\gamma\gamma$ mode, would imply the presence of an otherwise unobservable heavy generation.

Of course, we should remark that the hypothesis of a light $\phi^0$ in the presence of a heavy fourth generation could easily be inconsistent with vacuum stability in the context of the renormalization group equations (see the discussion in the introduction) unless there is new physics at a fairly low mass scale.

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.\textsuperscript{13,14} The sensitivity to anomalies in these couplings is potentially substantially greater than that provided by LEP-I data, but only if the $\phi^0$ can detected in the $\gamma\gamma$ or $\ell\gamma\gamma$ modes and the production cross sections reliably normalized.

Thus, the $\gamma\gamma$ and $\ell\gamma\gamma$ detection modes are of special importance. Discussion and contributions to the workshop focused on two important aspects of these modes. 1) The impact of next-to-leading order (NLO) corrections. 2) The impact of experimental ‘reality’.

Let us discuss the latter first. The implications of the CDF diphoton measurements (performed at the Tevatron) for SSC backgrounds to the $\gamma\gamma$ and $\ell\gamma\gamma$ modes were explored by R. Blair. He considered two issues: a) to what extent can Monte Carlo estimates of the backgrounds be relied upon; and b) using current experimental data, what techniques can be employed to better improve our estimates of such backgrounds, especially backgrounds arising from jets that mimic an isolated photon. His conclusions are summarized below.

With regard to a), there are several sources of uncertainty that will remain right up to the turn on of either the LHC or SSC. First, our estimates of the overall rate of the irreducible background from prompt diphotons at LHC and SSC energies may not be reliable. Currently the rate measured at the Tevatron by CDF appears to exceed the expected rate by more than a factor of two, even after including NLO corrections.\textsuperscript{15} This may indicate an even larger uncertainty for SSC/LHC experiments.
since the region of $x$ of interest at CDF is roughly $2E_T/\sqrt{s} \approx .01$ and the corresponding region for a 100 GeV $\phi^0$ at the SSC would be .005. Thus, background estimates that use the conventional distribution functions for computing the QCD prompt diphoton rate may be more than a factor of two low. HERA measurements of the relevant quark and gluon distribution functions should help to clarify this issue. Another possible area where the background may be underestimated is in the hadronic background. At low $E_T$ the prompt photon cross section as measured by CDF is also higher than that calculated from QCD. Consequently, if the background that comes from a single photon plus a jet that fragments to a photon candidate is computed based on QCD, it would be underestimated.

While CDF data raises the uncertainties outlined above, it can also be used to pin down other ingredients in the background calculations for the SSC/LHC. In particular, using existing measurements (by CDF and D0) it is possible to determine how often a jet will fragment to what appears to be an isolated photon candidate. If this rate is coupled with the best estimates of jet-jet and photon plus jet production at future colliders (including the uncertainties listed above) it should be possible to quantify the corresponding level of background that will be observed in a sample of diphoton candidates for any given choice of isolation cut. Additional detector dependent information can then be used to reject or evaluate the non-prompt portion of the background. This last step depends on the intimate details of the detector and ultimately must be evaluated once the detector is up and running, using in situ sources of single photons or background (like $\gamma$'s from $\eta$ decays or $\pi^0$'s) as CDF has.\[16\] The initial cut on isolation typically reduces the hadronic background rates by more than two orders of magnitude, as we shall shortly see, with isolation-only rejection factors of order $5 \times 10^{-4}$ being typical. CDF experience suggests that achieving an additional detector dependent rejection factor of more than one order of magnitude will not be easy. Since, adequate suppression of backgrounds in the $\gamma\gamma$ ($\ell\gamma\gamma$) modes requires a total rejection factor of $10^{-4}$ ($\sim 3 \times 10^{-4}$) for each jet, we see that the required rejection is fairly certain in the $\ell\gamma\gamma$ case, but on the borderline for the inclusive $\gamma\gamma$ mode.

Let us now return to the rejection factor that can be achieved purely by an isolation cut. At noted above, this can be directly measured by CDF. The procedure is to determine by other means (mainly shower profiles) the non-prompt component of the event rate for a single isolated electromagnetic cluster (i.e. isolated photon candidate). By comparing this rate to the inclusive jet rate (at the same $E_T$), the fraction of jets that fake an isolated photon candidate can be computed directly as the ratio of these two experimentally measured rates. CDF results for the isolation cut rejection factor appear in Fig. 3, for a variety of $E_T$ values. The isolation cut is defined by the fraction of additional energy observed in a cone of size $\Delta R = 0.7$ centered on the candidate. As noted above, quite decent isolation-only rejection factors of order $5 \times 10^{-4}$ are achievable even for an additional energy fraction as large as 0.1. The rejection factors of Fig. 3 can be checked directly by CDF by using them to determine the background to two-photon production coming from one-photon plus
Figure 3: Rate at which jets turn into photon candidates for various values of isolation cut. The candidates correspond to more than one photon, presumably originating from hadronic decays. Only background candidates are included here; the prompt single photon contribution has been subtracted off using a method similar to that described in Ref. 16. Figure prepared by R. Blair.
jet production; simply multiply the latter by the appropriate $E_T$-dependent rejection factor. The resulting diphoton background can be compared to that directly measured using shower profile information. For a common choice of isolation-cone criteria, the agreement is quite good.

Thus, the CDF results tend to provide a warning that backgrounds could be larger than predicted using purely Monte Carlo computations, even after NLO corrections are included. But, they also indicate that the jet rejection factors required to eliminate non-prompt backgrounds are achievable. In regard to the latter, it should be noted that SSC/LHC analyses typically employ an isolation criteria of $\lesssim 5$ GeV of extra activity inside a cone centered on the photon candidate. This corresponds to an extra energy fraction of $\lesssim 0.25$ for photons with $E_T$ of order 20 GeV, the lowest $E_T$ generally allowed in the SSC/LHC cuts. This is a larger fraction than appears in Fig. 3, and typically employed in the above CDF analysis. Fig. 3 suggests that the jet-photon discrimination that can be achieved on the basis of isolation alone, for an extra energy fraction of $\lesssim 0.25$, might be worse than $10^{-3}$. An energy fraction of $\lesssim 0.1$ would correspond to rejecting extra energy that exceeds only about 2 GeV. While this can be imposed at the Tevatron in the CDF analysis of single photon and diphoton events, it might be that the SSC events have more underlying minimum bias structure. In this case, the efficiency for such a cut on the Higgs signal might not be very large. This is why the SSC/LHC analyses have adopted the more conservative cut on extra hadronic energy in the isolation cone. Thus, the CDF results for isolation rejection are not immediately applicable to the SSC/LHC analyses. However, studies by the GEM collaboration have suggested that realistic (in the SSC context) cuts yield isolation and detector-dependent rejection factors which combine to give a net jet rejection factor of order $10^{-4}$. Similar claims have been made by the LHC collaborations. Ultimately, this issue must be settled after the SSC/LHC begins operation.

Let us now turn to the impact of NLO corrections. In one contribution, B. Bailey and D. Graudenz have evaluated the NLO corrections to both the inclusive $\gamma\gamma$ Higgs signal and the $\gamma\gamma$ continuum background (including $gg \rightarrow \gamma\gamma$ via the box diagram). The computation was done in the Monte Carlo context as delineated in Ref. 15, for both signal and background. This allows an accurate implementation of the standard cuts (e.g. $p_T \gtrsim 20$ GeV, $|\eta| \lesssim 2.5$ and hadronic energy inside isolation cone $\lesssim 4 - 5$ GeV, for the photons) required to establish a signal. Their results show that both signal and background are enhanced by K-factors of order 1.5 to 1.7, with the K-factor for the $\phi^0$ signal being larger than that for the background for $m_{\phi^0} \lesssim 120$ GeV. Thus, not only is the absolute event rate enhanced, but also the signal to background ratio, $S/B$. The result is that the number of standard deviations of significance of the signal, $N_{SD} \equiv S/\sqrt{B}$, is larger according to NLO computations than when only leading order (LO) computations are employed, typically by a factor of order 1.3, rising to 1.4 for lower $m_{\phi^0}$ values near 80 GeV. This represents an important addition to our confidence level for the viability of the inclusive $\gamma\gamma$ mode.

A second contribution, by D. Summers, focuses on the associated $W\phi^0$ and
production modes, leading to the $\ell\gamma\gamma$ signature for the Higgs boson. His reanalysis confirms earlier results for signal and background rates for the two modes. In particular, for an integrated luminosity of $L = 10 \text{ fb}^{-1}$ at the SSC, we recall that the $W\phi^0 \rightarrow \ell\gamma\gamma X$ mode has a rather low event rate, and only by virtue of the $t\bar{t}\phi^0$ production mode is a viable signal achieved in the $\ell\gamma\gamma$ channel. In comparison, at the LHC $L = 100 \text{ fb}^{-1}$ is required in order that the combined signal be viable. At the SSC, for $L = 100 \text{ fb}^{-1}$ event rates in the $\ell\gamma\gamma$ channel coming from $W\phi^0$ and $t\bar{t}\phi^0$ are both quite large (e.g. 85 and 170 events, respectively, for $m_t = 140 \text{ GeV}$ and $m_{\phi^0} = 110 \text{ GeV}$). Backgrounds at $L = 100 \text{ fb}^{-1}$ from $W\gamma\gamma$ and $t\bar{t}\gamma\gamma$ (where $q\bar{q}$ collisions can be as important as $gg$ collisions) are only 24 and 23 events, respectively (for $\Delta m \sim 0.03m_{\phi^0}$). Thus, there is an excellent chance that with high luminosity at the SSC the $W\phi^0$ and $t\bar{t}\phi^0$ production mechanisms can be separated (on the basis of much larger jet activity for the $t\bar{t}\phi^0$ mechanism), allowing a determination of the ratio of the $W^+W^-$ and $t\bar{t}$ couplings of the $\phi^0$.

Exactly how low in $m_{\phi^0}$ one can go in these modes is an important issue. If it is only the irreducible backgrounds that are important, then, with $L = 10 \text{ fb}^{-1}$ at the SSC, the $t\bar{t}\phi^0$ mode remains viable all the way down to the LEP upper limit of $m_{\phi^0} \sim 60 \text{ GeV}$. For $L = 100 \text{ fb}^{-1}$ at the SSC, the $W\phi^0$ mode also yields a viable signal for such low masses. This is nicely illustrated in the tables of Ref. 19. However, for $m_{\phi^0}$ below about 80 GeV, the reducible backgrounds from $W\gamma j + Wjj$ and $t\bar{t}\gamma j + t\bar{t}jj$ production, where the $j$’s appear to be photons, become much more problematical. We will only know how low in mass these searches can be carried out after the detector is up and running and the jet/photon discrimination factor can be experimentally measured.

As Summers emphasizes, the above results are based on LO computations. NLO corrections need to be examined. It is well known that NLO corrections to the $W\phi^0$ process are closely related to those for inclusive on-shell $W$ production; one need only extend the calculation to an off-shell $W^*$ which decays to $W\phi^0$. Perturbative calculations to $\mathcal{O}(\alpha_s)$\cite{20} yield a $\sim 10\%$ increase in the $W\phi^0$ cross section at NLO compared to LO. At the workshops, P. Agrawal and C.P. Yuan improved these calculations by taking into account the effect of multiple soft gluon emissions on the kinematics of the decay products, using the Collins-Soper-Sterman techniques for gluon resummation.\cite{21} They find that the net K-factor lies in the 5-15% range, depending upon the choice of the parton distribution functions employed in the tree-level calculation. (The tree-level result is larger for NLO distributions than for LO distributions, leading to a smaller effective K-factor if NLO distributions are used for both the tree-level and the full $\mathcal{O}(\alpha_s)$ computations.) Also important is the extent to which the NLO QCD corrections alter the distributions for the transverse momentum and rapidity of the Higgs boson or of its photon decay products. Agrawal and Yuan show that resummation effects do not affect the earlier conclusions found in Ref. 20, namely that NLO corrections do not alter these distributions relative to the naive LO computations. More details can be found in Ref. 22.
NLO corrections to $t\bar{t}\phi^0$ production are likely to be much larger than those for $W\phi^0$, with a K-factor of order 1.5 to 2 not being improbable. However, an explicit computation has not yet appeared. What about the background processes? There seems to be no reason to expect the K-factor for the $t\bar{t}\gamma\gamma$ to be any larger than that for the corresponding signal. However, the situation for the $W\gamma\gamma$ background to $W\phi^0$ is likely to be different. In Ref. 23 it is shown that the NLO corrections to $q'\bar{q} \to W\gamma$ enhance the cross section by a factor of about 3. This can be understood by noting that the LO $W\gamma$ cross section is suppressed by the presence of radiation zeroes for certain angles of the final state $W$ and $\gamma$ with respect to the incoming quarks. At next-to-leading order these radiation zeroes are absent. Thus the NLO cross section is of standard size, while the LO cross section is artificially suppressed. This same pattern is expected to continue in the case of the $W\gamma\gamma$ background to the $W\phi^0 \to W\gamma\gamma$ signal. Calculations of the real radiation processes (e.g. $gg \to W\gamma q'$) yield cross sections of order 3.5 times the LO $q'\bar{q} \to W\gamma\gamma$ cross section even after a $p_T$ cut of $\geq 50$ GeV is imposed on the extra jet.\textsuperscript{[24]} Although the virtual diagram components of the NLO computation have not yet been completed, it is clear that a net K-factor of order 2 to 3 is not out of the question.

While such an increase in the $W\gamma\gamma$ background would not be a particular problem for the overall $\ell\gamma\gamma$ signal (especially at the SSC where the $t\bar{t}\phi^0 \to \ell\gamma\gamma X$ process is much larger), one might worry that it would make the $W\phi^0$ signal much more difficult to isolate. However, this need not be the case either. To isolate the $W\phi^0$ signal in the $\ell\gamma\gamma$ channel we have already noted that one must veto against extra jet activity in order to eliminate most of the $t\bar{t}\phi^0$ component. In other words, what is crucial is the $W\gamma\gamma + 0$ jets cross section. As Summers notes in his contribution, this will have far smaller NLO corrections than fully inclusive $W\gamma\gamma$ production, since the largest NLO corrections are due to real radiation processes. A full assessment of the ability to isolate $W\phi^0$ process requires the computation of the virtual NLO contributions to $W\gamma\gamma$ as well as a simulation of the $W\phi^0$ signal events which includes minimum bias low-$p_T$ jets and the gluon resummation effects discussed earlier. Hopefully, a definition of ‘0 jets’ can be found which maintains high efficiency for the $W\phi^0 \to W\gamma\gamma$ signal events of interest, while reducing the $W\gamma\gamma$ background to more or less the LO level.

In summary, the $\gamma\gamma$ and $\ell\gamma\gamma$ detection modes both appear to be highly viable for the $\phi^0$ in the MSM context, provided appropriate detectors are available. However, as noted earlier, if an extra heavy generation is present, the $\ell\gamma\gamma$ rates would be very suppressed and only the inclusive $\gamma\gamma$ channel is certain to be viable, and then only if 1% mass resolution and a jet-rejection factor of $\sim 10^{-4}$ are achieved. As discussed, our ability to achieve the latter cannot yet be regarded as proven. Other scenarios could also impact these modes. For instance, if the Higgs has only bosonic couplings (in which case a second Higgs or some other mechanism must be responsible for giving mass to fermions) then only the $W\phi^0$ production mechanism would survive intact. Since the $\phi^0 \to \gamma\gamma$ branching ratio is likely to be quite large (if not nearly 100%) in such a scenario, the resulting $\ell\gamma\gamma$ event rate would be truly enormous and
2.2. Detection of $t\bar{t}\phi^0 \to t\bar{t}b\bar{b}$

Although the importance of the $\gamma\gamma$ and $\ell\gamma\gamma$ modes is apparent, it would be highly desirable to be able to detect a light $\phi^0$ in its primary $b\bar{b}$ decay mode. Not only would this provide an alternative detection mode, it would allow us to normalize any Higgs signal that is seen in the $\gamma\gamma$ channels. Normally, the primary decay mode, $\phi^0 \to b\bar{b}$, is rejected as having too large a QCD background. In the MSM, this is certainly the case for inclusively produced $\phi^0$'s. (If an extra generation is present, the $gg \to \phi^0 \to b\bar{b}$ rate might be sufficiently enhanced that the mass of spectrum of two tagged $b$ jets would reveal a signal in the inclusive case.) However, at least two groups decided that the associated production/decay mode $t\bar{t}\phi^0 \to t\bar{t}b\bar{b}$ might hold promise if $m_t$ is large.

At this meeting we heard reports from D. Wu[25] and J. Gunion[26] which summarized the situation. Related work on the $W\phi^0$ associated production mode at the Tevatron was reported by S. Willenbrock[27].

Let us focus first on the $t\bar{t}\phi^0 \to t\bar{t}b\bar{b}$ process, which is relevant for the SSC and LHC. Two possible approaches to isolating the Higgs signal in the $t\bar{t}b\bar{b}$ final state can be considered. In the first, both top quarks are required to decay leptonically and one looks at the two-jet mass spectra for a Higgs bump. This was considered in both Ref. 25 and Ref. 26. In Ref. 25 the backgrounds considered were $t\bar{t}Z$, with $Z \to 2$ jets, and $t\bar{t}gg$, which is certainly the largest background in the $t\bar{t}2j$ channel. Approximate calculations of signal and backgrounds were performed. Not considered were the $t\bar{t}q\bar{q}$ backgrounds. The conclusion of the analysis of Ref. 25 is that after appropriate cuts (especially a cut requiring large $p_T$ for any jet pair) a signal for $m_{\phi^0} \sim 100$ GeV could be detected. In Ref. 26 the $t\bar{t}\phi^0$ signal and $t\bar{t}Z(\to q\bar{q}) + t\bar{t}q\bar{q}$ backgrounds were computed using exact matrix elements. The latter give a lower bound on the background level — the $t\bar{t}gg$ background can only worsen the situation. In order to compare to Ref. 25, exactly the same cuts were employed. The result (for one SSC year of $L = 10$ fb$^{-1}$) is that in the central 10 GeV bin centered on the $m_{\phi^0} = 100$ GeV Higgs mass peak a signal of $S = 221$ events is present over a background of $B = 2400$. (The latter figure includes combinatoric backgrounds of all types, including those from the signal events themselves.) This yields a nominal statistical significance of $S/\sqrt{B} = 4.2$. Even if an exact computation of the $t\bar{t}gg$ rate were to only double $B = 2400$ to $B = 4800$, the signal becomes extremely marginal. In comparison to Ref. 25, the signal rate of Ref. 26 is substantially smaller despite the inclusion of a ‘K’ factor of 1.6 in the latter computations. The estimate of Ref. 25 for the signal rate in a 10 GeV bin centered at 100 GeV is $S \sim 650$. The $t\bar{t}gg$ background over this same mass interval is estimated to be of order 9900 events. This nominally yields $S/\sqrt{B} \sim 6.5$, before inclusion of $t\bar{t}q\bar{q}$ backgrounds.
The second approach to isolating the $t\bar{t}b\bar{b}$ final state is to employ $b$-tagging. At least three $b$’s must be tagged in order to make headway against the general $t\bar{t}qq+ttgg$ backgrounds. This cuts down on the event rate to such an extent that demanding double lepton tagging does not yield a viable signal. In Ref. 26, a single lepton trigger is employed. The other top quark is allowed to decay either leptonically or hadronically. The signal and backgrounds were examined in the two cases where a) 3 or more $b$’s are tagged and b) 4 or more $b$’s are tagged. Crucial ingredients are the efficiency for tagging a $b$ ($e_{b-tag}$) and the probability that a light quark or gluon jet is mis-identified (i.e. mis-tagged) as a $b$ ($e_{mis-id}$). After appropriate kinematic cuts two cases were examined: i) $e_{b-tag} = 0.3$ and $e_{mis-id} = 0.01$, and ii) $e_{b-tag} = 0.4$ and $e_{mis-id} = 0.005$. After a few other cuts, none of which required reconstruction of either the hadronically decaying $W$ or $t$ quark, the invariant mass spectrum for two tagged $b$ jets was examined for signal vs. background. The results are encouraging for $m_t \sim 130 - 140$ GeV and $m_{\phi^0} \lesssim 110$ GeV. Table 1 gives the number of years required to see a 5 sigma signal (along with the associated number of signal events in parentheses) for different Higgs masses and top quark masses in the four scenarios: I=a)+(i); II=a)+(ii); III=b)+(i); and IV=b)+(ii). Typically, the $\phi^0$ is observable in $\lesssim 3$ SSC years if $m_{\phi^0} \lesssim 100$ GeV for $m_t = 140$, and if $m_{\phi^0} \lesssim 120$ GeV for $m_t = 180$ GeV. (The actual peaks for $m_t = 180$ GeV are quite dramatically visible relative to the
background — the only issue is event rate.)

One virtue of this mode that should be apparent from Table 1 is that an $m_{\phi^0}$ value below 80 GeV would actually yield an even more significant signal. This is in contrast to the $\gamma \gamma$ modes which start to die out below $m_{\phi^0} \sim 80$ GeV due to decreasing $\gamma \gamma$ branching ratio and increasing $\gamma \gamma$ continuum backgrounds. As discussed earlier, our ability to extend the $\gamma \gamma$ and $\ell \gamma \gamma$ mode searches to masses as low as the $\sim 60$ GeV upper limit of LEP is crucially dependent upon the level of reducible $\gamma j$ and $jj$ type backgrounds.

What are the expectations for $b$-tagging? The scenario of 30% tagging efficiency and 1% mis-identification probability, referred to as a) above, is quite close to that obtained in the SDC TDR study of $b$ vertex tagging. In fact, if the actual efficiencies computed in the SDC TDR are employed, the results of Table 1 improve somewhat. Our only current experience with $b$ tagging at a hadron collider is that for the SVX vertex detector at CDF. These $b$-tagging results were summarized by N.M. Shaw in his contribution to these proceedings. Three different algorithms have been explored so far, labelled 'jet vertexing', 'jet probability', and '$d\phi$ clustering'. All give similar results. The net efficiency for tagging at least one of the two $b$ quarks in a $t\bar{t}$ event is about 22% at $m_t = 120$ GeV, corresponding to roughly 12% per $b$-jet. This efficiency includes SVX acceptance (which extends to roughly $|\eta| \lesssim 1$), but does not include any kinematic or high $p_T$ lepton cuts. Given the limited $|\eta|$ range and absence of momentum cuts, this number does not appear to be significantly worse than what is obtained in the SDC TDR Monte Carlo study. In any case, the SDC tracker is significantly better than that of CDF, and so a better efficiency for the SDC detector should be anticipated.

Tagging of $b$-jets using the $\ell$ from the semi-leptonic $b \rightarrow c\ell\nu$ decay mode has also been employed by CDF. For $m_t = 140$ GeV a cut on the secondary lepton’s $p_T$ of 2 GeV (relative to the jet axis) yields a net efficiency (not including kinematic and primary high-p$_T$ trigger lepton cuts) for tagging one or more $b$-jets of 23%, i.e. again of order 12% per $b$. (The fake tag rate is of order $10^{-2}$ per track.) If this same efficiency were to apply in the SDC context, by combining with a vertex tagging efficiency of order 30% quite large net $b$ tagging efficiencies at the SSC could be achieved, perhaps as large as the 40% assumed in scenario b) in the $t\bar{t}b\bar{b}$ channel study.

In any case, these respectable CDF $b$-tagging results and the promising results outlined above for Higgs detection in the $t\bar{t}b\bar{b}$ channel using $b$-tagging should encourage the SSC/LHC detector groups to pay close attention to their $b$-vertex detectors. It is clear that maximizing the efficiency and purity of $b$-tagging could have crucial benefits in the arena of Higgs detection. Indeed, our group concluded that this was probably the single most important focus for improvement of detectors that might still be possible relatively late in the detector development and construction process.

The related work of Ref. 27, as reported by S. Willenbrock, focuses on the $W\phi^0 + Z\phi^0$ production modes, with $\phi^0 \rightarrow b\bar{b}$, at the Tevatron. (The current lower bound on $m_t$ rules out any significant cross section for $t\bar{t}\phi^0$ production at Tevatron energies.)
Leptonic decays of the $W$ or $Z$ (including $Z \rightarrow \nu \bar{\nu}$) are used to trigger the event (with appropriate cuts and isolation criteria for charged leptons) and at least one of the $b$ quarks is required to have $|\eta| < 2$ and large enough $p_T$ that it could be tagged. Before including the actual efficiency for $b$-tagging, the combined $W \phi^0 + Z \phi^0$ cross section after all cuts and branching ratios are included is about 0.2 pb at $m_{\phi^0} = 60$ GeV, falling to about 0.07 pb by $m_{\phi^0} = 100$ GeV. Two important backgrounds are $Wb\bar{b}$ and $Zb\bar{b}$. When integrated over an invariant mass bin of $\pm 2 \Delta M_{b\bar{b}}$ (where $\Delta M_{b\bar{b}} \sim 0.8 \sqrt{M_{b\bar{b}}}$) centered on $m_{\phi^0}$, these give a cross section of very nearly the same size as the signal (after the same cuts etc.). Backgrounds of similar size come from $ZZ$ and $WZ$ production. Less troublesome are the backgrounds arising from $t\bar{t}$ production by virtue of missing one of the $W$'s in the $t$ decays. $Wc\bar{c}$ and $Zc\bar{c}$ backgrounds could be comparable to their $b\bar{b}$ analogues, depending on the probability for the vertex tagger to mis-identify a $c$ quark as a $b$ quark. After all cuts etc., the $Wjj$ and $Zjj$ ($j =$ light quark or gluon) are about 100 times as large as the $Wb\bar{b}$ and $Zb\bar{b}$ backgrounds. Thus, light-quark/gluon vs. $b$ discrimination must be better than about $10^{-2}$. Assuming 50% efficiency for detecting a displaced vertex from one or both of the $b$ quarks in the signal (equivalent to 30% efficiency per $b$ quark, which is higher than the current efficiency for the CDF SVX, but may be achieved by the time the main injector is in operation), and a 1% mis-tag probability for light quark and gluon jets, one finds at $m_{\phi^0} = 60$ GeV and $L = 1000$ pb$^{-1}$ 50/58 $W\phi^0/Z\phi^0$ events compared to 60/52 $Wb\bar{b}/Zb\bar{b}$ irreducible background events and 200/200 $Wjj/Zjj$ mis-tagged events. The resulting statistical significances in the two channels are of order 3 standard deviations. These signals will be further diluted by inclusion of the $Wc\bar{c}/Zc\bar{c}$ backgrounds. Significant signal enhancement would occur if the mis-tagging probability for light quark and gluon jets could be reduced below 1%.

2.3 A New Approach to the $\phi^0 \rightarrow \ell^+ \ell^- \nu \bar{\nu}$ Signal

One of the modes suggested for detecting the $\phi^0$ is $pp \rightarrow ZZ \rightarrow \ell^+ \ell^- \nu \bar{\nu}$, where the Higgs appears as a resonance on a Jacobian background. Unfortunately, there are QCD background processes which mimic the final state that can only be removed via stringent kinematic cuts. In the process, a significant fraction of the signal is lost. One approach to relaxing these stringent cuts and extracting a better signal is that of tagging one of the spectator jets that arise when the $\phi^0$ is produced via the $WW$ fusion mechanism.$^{[26]}$ In a contribution to the workshops,$^{[30]}$ Duncan and Reno have examined a method, that employs the distribution of the final state charged leptons, which could enhance the gluon fusion component of the signal in this same channel. It relies on the fact that the $Z$'s coming from the $\phi^0$ are mainly longitudinally polarized, whereas all backgrounds (e.g. $qq \rightarrow qZ$, where the $q$ disappears down the beam pipe, as well as the irreducible $q\bar{q} \rightarrow ZZ$ process) yield transversely polarized $Z$'s. Also important is the fact that the $\phi^0$ tends to be produced nearly at rest in $gg$ fusion.
If the center of mass of the $\ZZ$ system could be determined, then the distribution in $z = \cos \theta$ for the $\ell^+\ell^-$ pair (where $\theta$ is the angle of one of the leptons with respect to the boost direction of the parent $Z$) could be determined. The shape of this distribution would then yield a good estimate for the fraction $f_L$ of events in which the visible $Z$ is longitudinally polarized. However, since one of the $Z$’s decays invisibly, the $\ZZ$ center of mass cannot be determined. Duncan and Reno define an alternative variable $z^* \equiv 2|p_\ell \cdot \epsilon_L|/m_Z$ (where $p_\ell$ is the momentum of one of the charged leptons and $\epsilon_L = (|\vec{p}_Z|, \vec{p}_Z E_Z / |\vec{p}_Z|)/m_Z$), which coincides with $z$ when the $\ZZ$ center of mass and laboratory frames are the same. Since, in $gg$ fusion, a heavy Higgs boson is largely produced nearly at rest, $z^*$ should be a good approximation to $z$ in the case of the signal contribution to this channel. If one computes the average value of $z^*$, then theoretically $0.375 < \langle z^* \rangle < 0.5625$, the lower (upper) bound being reached if only purely longitudinal (transverse) $Z$’s are produced. Combining the $q\bar{q} \rightarrow ZZ$ background and the $gg \rightarrow \phi^0 \rightarrow ZZ$ signal, a plot of $\langle z^* \rangle$ shows a pronounced dip as a function of the transverse mass, $m_T$, in the vicinity of the Higgs resonance. Of course, the vector boson fusion contributions to the $ZZ$ channel (both resonance and continuum background) must be included, as well as backgrounds such as that from $qg \rightarrow qZ$; and errors must be analyzed. This work is in progress.

2.4 Heavy Higgs and Strongly Interacting $W$ Scenarios

As noted earlier, our ability to probe a strongly interacting EWSB sector is reviewed in contributions by Bagger,\cite{31} Cheung,\cite{32} and Morris.\cite{33} Duncan Morris reported on work done in collaboration with R.D. Peccei and R. Rosenfeld in which they considered the phenomenology of a strongly interacting Higgs sector patterned after hadronic pion physics.\cite{34} The central idea is to go beyond the low energy theorems relating the behaviour of pions and longitudinal weak gauge bosons and to assume, essentially by fiat, that nonperturbative pion physics can serve as a model for a strongly interacting $W_L$ ($W \equiv W^\pm, Z$) sector. By scaling up multipion energy scales by a factor of $\sqrt{v/f_\pi} \simeq 2600$ one can relate multipion production on the GeV scale to multi-$W_L$ production on the TeV scale. Though there is little reason to expect that a strongly interacting Higgs sector should mimic pion physics in every detail, the assumed equivalence permits a useful exploration of the phenomenological consequences. A literal interpretation results in $O(1-30)$ fb cross sections at the SSC and LHC for the production of multiple $W_L$’s. When cuts are imposed to appropriately reduce background processes such as multiple top quark production, they find multi-$W_L$ signatures of multiple high-$p_T$ leptons and high-$p_T$ jets with cross sections of $O(1$ fb) at the SSC. Multiple gauge boson production through generic perturbative processes, strongly interacting Higgs sectors and a possible breakdown in electroweak perturbation theory are also discussed in the contribution by Morris.

Jon Bagger reviews the work of a large collaboration which explored strong $W_LW_L$ interactions in a very wide class of models.\cite{35} More specifically, the processes $pp \rightarrow$
We studied, and the rates for the “gold-plated” channels, where \( W^\pm \rightarrow \ell^\pm \nu \) and \( Z \rightarrow \ell^+\ell^- \) (\( \ell = e, \mu \)), were computed for each model. Using a forward jet-tag, a central jet-veto and a back-to-back lepton cut to suppress the Standard Model backgrounds, it was demonstrated that the SSC and LHC have substantial sensitivity to strong interactions in the electroweak symmetry breaking sector. Of course, the channels examined, \( W^+W^- \rightarrow \ell^+\ell^-\nu\bar{\nu}, W^+Z \rightarrow \ell^+\ell^-\ell\nu, ZZ \rightarrow \ell^+\ell^-\ell^+\ell^- \), and \( W^+W^+ \rightarrow \ell^+\ell^+\nu\nu \), do not all yield adequate signals in 1-2 years of canonical SSC or LHC luminosity for all models. Instead, a significant signal can always be found in the channels that most naturally complement the particular type of model considered. In particular, models with a resonance of definite isospin are most easily probed using the \( WW \) channels that have resonant contributions from that same isospin. Non-resonant models are often best probed in the \( W^+W^+/W^-W^- \) channels. Indeed, one important conclusion is that different types of models can be distinguished experimentally by determining the relative magnitude of the \( LL \) signals in the four channels listed above. A large part of the work focused on the techniques required to suppress reducible and, especially, irreducible backgrounds to a level such that the low \( LL \) signal event rates in the purely leptonic channels can be isolated. In particular, the irreducible backgrounds from production of \( WW \) pairs with \( TT \) and \( LT \) polarizations end up being most important, and the techniques developed in Ref. 35 are particularly focused on suppressing them. Although the calculations do not include detector effects, they should survive more sophisticated Monte Carlo analyses. In particular, the types of cuts employed should be directly applicable in the experimental analyses that will be performed when actual data becomes available.

In the contribution by K. Cheung, the types of cuts employed in Ref. 35 are discussed in greater depth, emphasizing the fact that in the \( WW \) channel the appropriate cuts for minimizing the irreducible background from \( TT + LT \) modes relative to the \( LL \) signal are model independent. Detailed graphs explaining how the cuts were approximately optimized are also given. As an aside, we also note that there is now good agreement between the leading-log Monte Carlo treatments and parton-level treatments of \( t\bar{t} \)-related and other backgrounds. Thus, we have good reason for confidence that the purely leptonic signals for a strongly interacting \( WW \) sector can be extracted.

Overall, it certainly appears to be possible to probe a strongly interacting electroweak symmetry breaking sector at the SSC or LHC using the “gold-plated” purely-leptonic modes. Even if a light Higgs boson is found, it will be important to measure the event rates at high \( WW \) mass in all the various channels in order to make certain that the Higgs boson completely cures the bad high-energy behavior in all \( WW \) scattering subprocesses. The low event rates for the purely-leptonic final states imply that of order 2-3 years of \( 10^4 \text{ pb}^{-1} \) yearly luminosity will be required to conclude that there is no obvious \( W_L W_L \) enhancement in any of the four channels. Because of the relative cleanliness of these final states, the option of achieving this required integrated luminosity via enhanced instantaneous luminosity should be strongly considered.
2.5 Conclusions for the Standard Model

There seems to be little question that the SSC (and LHC) can probe electroweak symmetry breaking for all the most attractive scenarios that can be envisioned in the context of the Standard Model, including the renormalization-group-motivated light intermediate-mass range, as well as the case of a strongly interacting $WW$ sector.

3. The Minimal Supersymmetric Model Higgs Sector

As outlined in the Introduction, our goal is to establish a no-lose theorem according to which we are guaranteed to see at least one of the MSSM Higgs bosons regardless of the model parameter choices. For much of parameter space, the $h^0$ is relatively light and rather SM-like, and is the one that we could be certain of detecting. Going beyond the no-lose theorem, it is desirable to be able to see one or more of the heavier $A^0$, $H^0$ and $H^\pm$ Higgs states. However, let us first focus on issues relevant to the no-lose theorem. Of special concern will be the $\gamma\gamma$, $\ell\gamma\gamma$ channels and the $t\bar{t}b\bar{b}$ channel. Our discussion of these modes will temporarily ignore the possibility of superparticle decays of a light neutral $h^0$ or $H^0$.

3.1. The $\gamma\gamma$ and $\ell\gamma\gamma$ Detection Modes

A particularly interesting question is the extent to which the $\gamma\gamma$ widths of the MSSM Higgs bosons, especially that of the $h^0$, depend upon the SUSY context and/or superpartner masses. Some exploration of this issue has appeared in Refs. 11 and 36. 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).) In Fig. 1 several cases are illustrated. In Fig. 1 we show that if the MSSM parameters are chosen such that all new particles beyond the SM are heavier than about 250 GeV (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, the $WW$ coupling of the $h^0$ rapidly approaches its SM value and the squark and chargino loops become negligible. However, Fig. 1 also shows that if the lightest chargino is allowed to be as light as $\sim m_Z/2$ (by taking $M = 100$ GeV and $\mu = -65$ GeV), so that it only just evades detection at LEP, then the $h^0 \to \gamma\gamma$ width ratio can be suppressed to as little as half the corresponding SM value. In the case of the $H^0$, such a light chargino can greatly enhance the width. This is also illustrated in Fig. 1, where $\Gamma(H^0 \to \gamma\gamma)/\Gamma(\phi^0 \to \gamma\gamma)$ is plotted in the same cases as considered for the $h^0$.

The effect upon the $\gamma\gamma$ widths of the $h^0$ and $H^0$ of reducing further the common squark-slepton mass scale, while maintaining a light chargino, is also illustrated in
Fig. 1. While modifications in the $h^0$ width are not dramatic in the $\gtrsim 50$ GeV mass range of interest, the $H^0\gamma\gamma$ coupling can be significantly suppressed when squark and slepton masses are made light.

To determine the implications of these results for the discovery potential for the MSSM Higgs bosons in the $\gamma\gamma$ and $\ell\gamma\gamma$ channels requires additional ingredients. Consider first the inclusive $\gamma\gamma$ mode. In the case of inclusive $gg$ fusion production of the Higgs boson, followed by $\gamma\gamma$ decay, the event rate is determined by $\Gamma(\text{Higgs} \to gg) \times BR(\text{Higgs} \to \gamma\gamma)$. In the MSSM the effects of including squark loops in the $gg \to h$ ($h = h^0$ or $H^0$) coupling computations are generally numerically small. This was already evident from the early work of Ref. 37. Potentially more important are modified couplings of the $h$ to the fermions in the loops which dominate the $gg \to h$ coupling. However, the $h^0t\bar{t}$ coupling tends to be fairly SM-like for large $m_{A^0}$, and, consequently, the $gg \to h^0$ coupling tends to be rather SM-like. Similarly, for $m_{H^0}$ near its lower limit (the only region where the $H^0 \to \gamma\gamma$ mode is viable) the $H^0t\bar{t}$ and $gg \to H^0$ couplings tend to be SM-like. It turns out that a more significant impact on the inclusive $\gamma\gamma$ event rate occurs as a result of modifications to the $h \to \gamma\gamma$ branching ratios due to other channels. For instance, the $h^0b\bar{b}$ coupling at $m_{A^0} = 400$ GeV is still significantly above the SM value when $\tan \beta$ is large. Thus, $\Gamma(h^0 \to b\bar{b})$ will be larger than in the SM and $BR(h^0 \to \gamma\gamma)$ will be correspondingly suppressed. In addition, supersymmetric particle pair channels can enter into the $h$ decays. For the light chargino/neutralino sector scenario considered in Fig. 1, this is possible even for the $h^0$ when $m_{h^0}$ approaches its upper limit. Finally, in the case of the $H^0$, if $\tan \beta$ is large then $b\bar{b} \to H^0$ fusion production can be comparable to $gg \to H^0$ production.

To illustrate the importance of these additional effects, Fig. 2 shows the ratio of $\Gamma(h \to gg) \times BR(h \to \gamma\gamma)$ (for $h = h^0$ or $H^0$) to the same quantity for $h = \phi^0$. The general viability of the $\gamma\gamma$ mode for the $h^0$ (indicated by a ratio not too far below one) is apparent, the exception being for $m_{h^0}$ near its upper limit in the light chargino/neutralino mode cases (when $\tilde{\chi}\tilde{\chi}$ modes become kinematically allowed). The very narrow region for $H^0$ discovery in this mode is also apparent. Clearly, the enhanced $\gamma\gamma$ width found for the $H^0$ when charginos are light (see Fig. 1) does not lead to any significant expansion of the very narrow region for which $H^0$ discovery in the inclusive $\gamma\gamma$ mode is possible.

Turning to the $\ell\gamma\gamma$ channel, we first note that the $Wh$ and $t\bar{t}h$ production rates are determined by the $hW^+W^-$ and $ht\bar{t}$ couplings, which in the relevant $m_{h^0}$ and $m_{H^0}$ mass ranges tend to be fairly SM-like, as noted above. The $h \to \gamma\gamma$ branching ratio is subject to the modifications outlined above. Differences from the $\ell\gamma\gamma$ rate expected for the $\phi^0$ are to be expected, although over much of the relevant mass range such differences will be in the range of a factor of two to three.

Certainly, these results show that the potential for MSSM Higgs discovery in the $\gamma\gamma$ and $\ell\gamma\gamma$ channels must be reassessed as LEP-II and, eventually, the SSC provide information on the masses of the supersymmetric particles that enter into the Higgs
decays and the loop diagrams responsible for the $h\gamma\gamma$ and $hgg$ couplings. Even if supersymmetric decays of the $h^0$ and $H^0$ are not allowed, if either is detected in these channels, the exact rate(s) could yield valuable self-consistency tests for the masses and Higgs boson couplings of the supersymmetric particles.

The background to detection of a neutral MSSM Higgs boson in the inclusive $\gamma\gamma$ mode are also affected by superpartner particles. In particular, there are squark loop contributions to the $gg \to \gamma\gamma$ background process (which provides roughly 30-50% of the continuum background in the $\gamma\gamma$ inclusive channel). These were examined by A. Djouadi and G. Belanger. Their results are summarized below.

In the limit where the mass of the internal particle is much larger than the invariant mass of the two photons $\hat{s}$, $\kappa \equiv \hat{s}/(4m^2) \ll 1$, the helicity amplitudes for loops of scalar and fermion particles belonging to the fundamental triplet representation are easily extracted. In terms of the angle $\theta$ of one of the outgoing photons relative to an incoming gluon one finds:

\begin{align}
\text{scalars} & : \frac{1}{2} \sum_{\lambda_i} |M_{\lambda_1,\lambda_2,\lambda_3,\lambda_4}|^2 = \frac{272}{2025} \kappa^4 (3 + \cos^2 \theta)^2, \\
\text{fermions} & : \frac{1}{2} \sum_{\lambda_i} |M_{\lambda_1,\lambda_2,\lambda_3,\lambda_4}|^2 = \frac{4448}{2025} \kappa^4 (3 + \cos^2 \theta)^2. \tag{1}
\end{align}

For the cross section one has to include the factors: 1/2 for two identical photons in the final state; 2 for the color factor $(TrT^aT^b)^2$; color and polarization averaging factors 1/8 and 1/2 for each gluon in the initial state. The result is:

\begin{equation}
\frac{1}{256} \alpha^2 \alpha_S^2 \sum_{\lambda_i} \sum_q (e^2 \lambda_{\lambda_1,\lambda_2,\lambda_3,\lambda_4})^2 dP S. \tag{2}
\end{equation}

Note that the final result for both the fermion and scalar amplitudes goes like $1/m^4$. ($1/m^2$ should have been enough for the decoupling, but the amplitude must be proportional to $F_{\mu\nu}$, i.e. to $s^2$, and to get the correct dimension we need $1/m^4$ from the loop.) Since the decoupling occurs so quickly, the approximation above is likely to work even for invariant masses $\hat{s}$ that are just slightly smaller than $4m^2$. Even more noteworthy is the small size of a scalar loop as compared to a fermion loop, which is roughly 16 times larger for the same mass. Even if we set the squark masses all equal to the top quark mass, the top contribution is still larger. Light quark loops make contributions that behave as $\log(\hat{s}/m^2)$ (multiplied by some $\pi^2$ factors), and will dominate by far both the top loop and the sum of all the squark loops. Thus, it seems that squark loop contributions to $gg \to \gamma\gamma$ are not numerically significant.

\* 6 boxes, 12 three-point and 3 two-point functions have to be taken into account.
Of course, the $Wh$ and $tth$ production processes responsible for the $\ell\gamma\gamma$ detection modes are tree-level processes, and will not be significantly influenced by superparticle loop corrections.

We turn now to QCD corrections for the $gg \to h$ production processes. As already emphasized in the SM $\phi^0$ section, when discussing the $\gamma\gamma$ inclusive mode, the NLO corrections provide a substantial enhancement to the $gg$ fusion cross section, and must be included. For the $h^0$ and $H^0$ CP-even Higgs bosons, these are the same as found for a $\phi^0$ of the same mass (neglecting squark loops). In a contribution to these proceedings [38] and another recent paper, [39] two groups have computed the NLO corrections (in the large $m_t$ limit) for the CP-odd $A^0$. They find that the real diagram corrections (i.e. those arising from diagrams in which a real gluon is radiated into the final state) are the same for the CP-even and CP-odd Higgs bosons (when written in terms of the respective LO cross sections). The only difference is in the virtual component of the NLO corrections. For instance, in the contribution by Kauffman and Schaffer, the virtual correction is characterized by the constant $N_c(\pi^2/3 + 2)$ for a CP-odd Higgs boson, as opposed to $(N_c\pi^2/3 + 5N_c/2 - 3C_F/2)$ for a CP-even Higgs boson. These too are quite similar. Thus, the K-factor for CP-even and CP-odd Higgs bosons will be almost the same.

3.2. The $t\bar{t}b\bar{b}$ Detection Mode for $h^0$ and $H^0$

As we have seen, the $t\bar{t}b\bar{b}$ channel, arising via associated Higgs+$t\bar{t}$ production with Higgs$\to b \bar{b}$ will be viable for a light SM $\phi^0$. Thus, it should be no surprise that it is quite likely that the $h^0$ or $H^0$ of the MSSM can also be detected in this way. [40] Indeed, over essentially all of parameter space either the $h^0$ or the $H^0$ has SM-like couplings.

In addition, if $m_t$ is not too large, the Higgs with SM-like couplings will be quite light, $m_h \lesssim 120$ GeV. Thus, based on the analysis of the SM Higgs boson, [26] but correcting for branching ratio and production rate differences, Ref. 40 obtains the discovery contours in $\tan \beta - m_{A^0}$ parameter space shown in Fig. 4, in the case $m_t = 150$ GeV and $m_\gamma = 1$ TeV. There, we have adopted the same contour labels as used in J. Gunion and L. Orr, Ref. 7, and Refs. 8 and 9, in order to facilitate comparison with these previous treatments of the no-lose theorem. Scenario (A) refers to the notation established in Ref. 9; it corresponds to the case in which decays of the Higgs bosons to chargino and neutralino pair channels are not kinematically allowed. In Fig. 4, we have chosen to display only those earlier-obtained contours that define the region (labelled by the large $\times$) where the SSC would not be able to detect any of the MSSM Higgs bosons using the $4\ell$, $\gamma\gamma$ (and/or $\ell\gamma\gamma$), and $t \to H^+b$ modes. (Discovery survey plots in the above-noted references include a fuller selection of channels.)

From these contours the following conclusions can be drawn. If $m_t \sim 150$ GeV, then detection of the $h^0$ in the $t\bar{t}b\bar{b}$ mode will be possible for any $m_{A^0} \gtrsim 110$ GeV, the precise lower limit being slightly $\tan \beta$ dependent. (Note that for such moderate to large values of $m_{A^0}$ it is the $h^0$ which is SM-like). For $m_{A^0} \gtrsim 50$ GeV essentially
Figure 4: Discovery contours (at the 4σ level) in $m_{A^0}$–tan$\beta$ parameter space for the SSC with $L = 30$ fb$^{-1}$ and LEP-200 with $L = 500$ pb$^{-1}$ for the reactions: a) $e^+e^- \rightarrow h^0Z$ at LEP-200; e) $Wh^0X \rightarrow l\gamma\gamma X$; g) $t \rightarrow H^+b$; h) $t\bar{t}H^0$, with $H^0 \rightarrow b\bar{b}$; and i) $t\bar{t}h^0$, with $h^0 \rightarrow b\bar{b}$. Each contour is labelled by the letter assigned to the reaction above, on the side of the contour for which detection of the particular reaction is possible. The large $\times$ indicates the location of the window where no MSSM Higgs could be discovered at LEP-II or the SSC without processes h) and i). We have taken $m_t = 150$ GeV, $m_{\tilde{t}} = 1$ TeV and neglected squark mixing. Charginos and neutralinos are taken to be heavy.
all the way up to the lower limit in \( m_{A^0} \) at which \( h^0 \) detection becomes possible, the \( H^0 \) is relatively light and is SM-like and can be detected in this mode. That is, either the \( H^0 \) or the \( h^0 \) can be detected in this way for \( m_{A^0} \gtrsim 50 \text{ GeV} \). This completely closes the no-lose theorem parameter space hole (indicated by the big \( \times \) in Fig. 4) that arises if only \( \gamma\gamma, \ell\gamma\gamma \) and \( 4\ell \) final states are employed. In fact, by combining just the \( t \to H^+b \) 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 \( m_t \sim 200 \text{ GeV} \), the parameter space region for which the \( t\bar{t} + h^0, H^0 \to t\bar{t} + b\bar{b} \) mode is viable is smaller; but this is simply correlated with the fact that other decay modes, most notably the \( ZZ^* \to 4\ell \) mode, of the \( H^0 \) and (rather heavy) \( h^0 \) acquire larger branching ratios, and become viable over a large range of parameter space. Thus, the no-lose theorem continues to hold, but a larger set of final state modes must be employed. See Ref. 40 for details.

Of course, it should be noted that this \( t\bar{t}b\bar{b} \) mode is not viable for observing a heavy \( H^0 \) or heavy \( A^0 \); at large \( \tan \beta \), although the \( b\bar{b} \) channel would dominate the decays of the \( H^0 \) and \( A^0 \), the \( t\bar{t}A^0 \) and \( t\bar{t}H^0 \) production processes are suppressed. At small \( \tan \beta \), the \( A^0 \) and \( H^0 \) decays would be dominated by other modes such as \( H^0 \to h^0h^0 \) and \( A^0 \to Zh^0 \) for masses below \( 2m_t \) and by the \( t\bar{t} \) channel when kinematically allowed.

To reiterate, combining all modes, the SSC/LHC alone will detect at least one and most probably several of the MSSM Higgs bosons. The detected Higgs boson is most likely to be the \( h^0 \). Indeed, it is not altogether 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. The other Higgs bosons are most likely to be moderately heavy, in which case they could not be detected using the modes discussed so far. We will return to these heavier Higgs bosons shortly.

3.3. Detection of an Invisibly Decaying Higgs Boson

The above discussion does not take into account alternative decay channels for the MSSM Higgs bosons. Generally speaking, dominance of the decays of a light Higgs boson by invisible channels is quite possible in supersymmetric models. In the case of the MSSM, decays to \( \tilde{\chi}_1^0\tilde{\chi}_1^0 \), where \( \tilde{\chi}_1^0 \) is the lightest supersymmetric particle, can be dominant and would be invisible if R-parity is conserved. In the MSSM, \( \tilde{\chi}_1^0\tilde{\chi}_1^0 \) dominance is possible for both the \( h^0 \) and the \( A^0 \). In supersymmetric models with spontaneously broken R-parity, the dominant decay mode of the lightest scalar Higgs boson is predicted to be \( h \to JJ \), where \( J \) is the (massless) Majoron. \( J \) interacts too weakly to be observed in the detector. In such models, the decays of the

\[ \text{† There is a crossover in the vicinity of } m_{A^0} \sim 110 \text{ GeV where the } h^0 \text{ and } H^0 \text{ interchange roles as being most SM-like.} \]
second lightest scalar Higgs boson can also be predominantly invisible, the two most important modes being \( JJ \) and \( hh (\rightarrow JJJJ) \).\(^{[43]}\)

Two possible detection modes for an invisibly decaying Higgs boson can be envisioned at a hadron collider. Associated \( Zh \) production was considered in Ref. 44, with the rough conclusion that a viable signal for \( h \rightarrow I \) (\( I \) being any invisible channel) can be extracted for \( m_h \lesssim 150 \text{ GeV} \) provided the \( hZZ \) coupling is Standard Model (SM) strength and \( BR(h \rightarrow I) \sim 1 \). As part of the workshops, an alternative based on associated production of top plus anti-top plus Higgs was considered.\(^{[45]}\) There, it is found that if the top quark is not too light (\( m_t \gtrsim 130 \text{ GeV} \)) and if \( BR(h \rightarrow I) \sim 1 \), then a viable signal for \( h \rightarrow I \) can be extracted for \( m_h \lesssim 250 \text{ GeV} \) when the \( ht \) coupling is of SM strength. Clearly, the \( Zh \) and \( t\bar{t}h \) modes are complementary in the sense that they rely on the vector boson vs. fermion couplings, respectively, of the Higgs boson. For a CP-even Higgs boson, which has both types of coupling, both modes tend to be viable, but for a CP-odd Higgs boson the \( ZZ \) coupling is absent at tree-level and only the \( t\bar{t}h \) mode outlined below could lead to a visible signal.

The procedure for the \( t\bar{t}h \) case is quite simple. One triggers on \( t\bar{t}h \) events via an isolated lepton (\( e \) or \( \mu \)) from one \( t \) decay. In order to further single out events containing a \( t\bar{t} \) pair, at least one of the \( b \)-quarks must be vertex tagged. The procedure is that outlined earlier in the \( t\bar{t}bb \) channel case. Results given below are for \( e_{b-tag} = 0.3 \). Mis-identification backgrounds are not significant so long as \( e_{mis-tag} \sim 0.01 \), and the corresponding number for \( c \)-quark jets is of order 0.05.

The invariant mass of each pair of jets, \( M_{jj} \), is computed and at least one pair not containing the tagged \( b \)-quark is required to have \( m_W - \Delta m_W/2 \leq M_{jj} \leq m_W + \Delta m_W/2 \). In addition, each pair of jets satisfying this criteria is combined with the tagged \( b \) jet(s) to compute the three-jet invariant mass, \( M_{bjj} \). One demands that \( m_t - \Delta m_t/2 \leq M_{bjj} \leq m_t + \Delta m_t/2 \) for at least one \( b\bar{b}j \) combination. Together, these two cuts greatly reduce the likelihood that the second top in a \( t\bar{t} \) event can decay leptonically and satisfy all our criteria. In fact, if both \( t \)'s decay leptonically, to leading order only \( t\bar{t}g \) events in which the non-tagged \( b \)-quark and the \( g \) combine to yield an invariant mass near \( m_W \) can pass the \( M_{jj} \) cut, and this non-tagged \( b \) plus the \( g \) jet must combine with the tagged \( b \) to give a mass near \( m_t \). If mass cuts of \( \Delta m_W = 15 \text{ GeV} \) and \( \Delta m_t = 25 \text{ GeV} \) are used, only a small fraction of signal events are eliminated when typical SDC jet and lepton energy resolutions are employed, whereas the reducible backgrounds are significantly decreased.

Finally, to reveal the invisibly decaying Higgs, one employs the missing transverse momentum, \( \vec{p}_T^{miss} \), for the event and computes \( M_{miss-\ell} \), the transverse mass obtained by combining the transverse components of the missing momentum and the lepton momentum, \( M_{miss-\ell}^2 \equiv (E_T^{miss} + E_T^\ell)^2 - (\vec{p}_T^{miss} + \vec{p}_T^\ell)^2 \). The \( t\bar{t}h \) events of interest are characterized by very broad distributions in \( M_{miss-\ell} \) and \( E_T^{miss} \). Cuts on both variables are made.

There are several sources of background. The most obvious is the irreducible
Table 1: Number of 10 fb$^{-1}$ years (signal event rate) at the SSC required for a 5$\sigma$ confidence level signal, for $m_t = 140$ GeV, $E_T^{miss} > 200$ GeV and $M_{miss-\ell} > 150$ GeV if $\Delta m_W = 15$ GeV and $\Delta m_t = 25$ GeV. $BR(h \to I) = 1$ is assumed.

| $m_t \setminus m_h$ (GeV) | 60  | 100  | 140  | 200  | 300  |
|---------------------------|-----|------|------|------|------|
| 110                       | 1.8(26) | 2.6(30) | 3.3(34) | 8.4(55) | 15.5(74) |
| 140                       | 0.3(19) | 0.4(22) | 0.7(29) | 1.4(41) | 2.9(60) |
| 180                       | 0.2(27) | 0.3(34) | 0.6(45) | 1.3(68) | 3.9(118) |

The number of SSC years required for a $N_{SD} = 5$ sigma significance of the signal compared to background for $BR(h \to I) = 1$ is given in Table 2, for a variety of $m_h$ and $m_t$ values. The statistical significance $N_{SD}$ is computed as $S/\sqrt{B}$. For any given integrated luminosity, $S$ is the total $t\bar{t}h$ event rate and $B$ the total $t\bar{t}Z + t\bar{t}g - \ell\nu + t\bar{t}g - b\nu\nu$ event rate, with $E_T^{miss} > 200$ GeV and $M_{miss-\ell} > 150$ GeV (and all other cuts) imposed. Also given (in parentheses) is the associated number of signal events ($S$). The associated number of background events ($B$) can be obtained from the relation $B = S^2/25$.

From this table, it is immediately apparent that detection of an invisibly decaying Higgs boson should be possible within 1 to 2 SSC years for $m_h \lesssim 200 - 250$ GeV if (as assumed in these calculations) its coupling to $t\bar{t}$ is of Standard Model strength and $m_t \gtrsim 130$ GeV. (The required number of years for non-SM coupling is obtained simply by dividing the results of Table 1 by the ratio of the $t\bar{t}$ coupling strength
squared to the SM strength squared.) For \( m_h \gtrsim 300 \text{ GeV} \), the \( t\bar{t}h \) event rate drops to a lower level such that more than 2 SSC years are required. However, it is rather unlikely that invisible decays would be dominant for a Higgs boson with mass above 200 GeV or so. The table shows that detection of the \( h \) generally becomes easier for heavier top quark masses. This is because the \( t\bar{t}Z \) and \( t\bar{t}(g) \)-related backgrounds are smaller and the signal rates somewhat larger than for smaller \( m_t \). However, even for \( m_t \) as low as 110 GeV, Higgs bosons with mass below about 140 GeV that decay invisibly should be detectable in less than 3 SSC years.

Of course, the \( N_{SD} \) values quoted assume that the normalization of the expected background will be well-determined by the time that the experiments are performed. This will require a good understanding of the missing energy tails as they actually appear in the detectors, calculation of the higher-order QCD corrections that were only estimated, and accurate knowledge of the parton (especially gluon) distribution functions. With the availability of HERA data, and through the analysis and study of \( t\bar{t} \) events in the actual detectors, it is likely that uncertainties in the relevant backgrounds can be brought down to the 20% level by the time that adequate luminosity has been accumulated that an invisible Higgs signal would become apparent at the SSC.

The above study was performed for a \( h \) that is a CP-even Higgs mass eigenstate. The \( t\bar{t}h \) rates would be somewhat different as a function of \( m_h \) for a mixed CP or CP-odd eigenstate. However, we do not anticipate that the results for such cases would differ by very much from those obtained here.

The immediate relevance of these results to the MSSM is somewhat model dependent. In ongoing work, H. Pois and J. Gunion have begun exploration of a variety of grand unification scenarios, in particular the so-called no-scale scenario. Especially in this latter, it is quite possible for neutralinos to be light enough that SUSY decay modes of the \( h^0 \) are allowed. Typically, the \( b\bar{b} \) and \( \tilde{\chi}_1^0\tilde{\chi}^0_1 \) modes then compete with one another. After also accounting for the \( t\bar{t}h^0 \) coupling, it is found that one can guarantee detection of the \( h^0 \) at the 5 sigma level after 3 SSC years in either the \( t\bar{t}b\bar{b} \) channel discussed earlier or the invisible decay channel, if not both.

### 3.4. Other Superparticle Decay Modes

More generally, 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. 46, 8, and 9. In particular, the above references show that a significant branching ratio for \( \tilde{\chi}\tilde{\chi} \) decays would make detection of the \( H^0 \) in the \( 4\ell \) mode impossible. At the same time, however, these same decays can lead to final states containing multiple leptons, which would be relatively free of background. The preliminary exploration of Ref. 46 indicates that detection of the excess events in such channels due to Higgs→\( \tilde{\chi}\tilde{\chi} \) decays may be possible.
3.5. Detection of the $H^+$ in the $t\bar{t}b\bar{b}$ Final State

Of course, if $\tilde{\chi}\tilde{\chi}$ channels do not dominate the decays of the $H^0, A^0$ and $H^+$, detection of these latter Higgs bosons must be via SM particle channels. As emphasized earlier, this appears to be a non-trivial task. One mode that has been investigated is $H^0, A^0 \rightarrow \tau^+\tau^-$, originally suggested by Kunszt and Zwirner.\cite{47} It potentially 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 absence of $\tilde{\chi}\tilde{\chi}$ decays). In the L3P simulation study,\cite{48} $A^0$ and $H^0$ detection in this mode is claimed to be viable for all $m_{A^0} \gtrsim 100$ GeV and $\tan\beta \gtrsim 7$.

What about the $H^+$? In many ways a charged Higgs boson is the hallmark of a truly non-minimal Higgs sector, and in particular of two-doublet models such as the MSSM. In contrast, the presence of more than one neutral Higgs boson could be due to additional singlet Higgs representations beyond the single MSM doublet. Work was begun during the course of these workshops,\cite{49} to assess the possibility of detecting a charged Higgs boson in the $gg \rightarrow b\bar{t}H^+ \rightarrow b\bar{t}t\bar{b}$ production/decay mode. A preliminary report on this work is given here. Two-doublet model-II type fermion couplings\cite{1} for the charged Higgs are assumed.

Events are tagged by requiring that one of the $t$'s decay to a leptonically decaying $W$. Three $b$-jets are required to be tagged. Cuts and efficiencies for these tags are the same as those used in previous studies of this final state. The second $W$ from $t$ decay is required to decay hadronically. $M_{jj}$ and $M_{bjj}$ cuts are imposed as described in Section 3.3. Finally, a plot of the $M_{bbjj}$ mass distribution is made, where both $b$'s are required to be tagged and the two $j$'s must not have been tagged. The only important backgrounds, after such cuts, turn out to be the $t\bar{t}b\bar{b}$ continuum QCD background, and $t\bar{t}g$ where the $g$ is mis-tagged (with $1\%$ probability) as a $b$-jet. The $t\bar{t}Z$, with $Z \rightarrow b\bar{b}$, background is much smaller than either. Typical results for the $M_{bbjj}$ distribution are shown in Fig. 5. In this figure, the $gg \rightarrow b\bar{t}H^+ + \bar{b}tH^-$ signal and $gg \rightarrow t\bar{t}b\bar{b}$ backgrounds are computed at LO. For three $b$ tagging, there will be a large K-factor by which these should be multiplied. By comparing the exact computation of the $2 \rightarrow 2$ processes $gb \rightarrow tH^+ + g\bar{b} \rightarrow \bar{t}H^+$ as performed in Ref. 50 to the results obtained from the LO $gg \rightarrow b\bar{t}H^+ + \bar{b}tH^- \rightarrow 2 \rightarrow 3$ calculation, before any cuts, this K-factor is estimated to be in the range 2-2.5. We have employed the $2 \rightarrow 3$ computation in order to more correctly account for $b$-tagging, including multiple tag possibilities and kinematic cuts.

Fig. 5 shows that a significant $H^+ + H^-$ signal will be present (especially if it is appropriate to multiply the $tbH^\pm$ signal (and $t\bar{t}b\bar{b}$ background) by a K-factor of order 2.5). At $m_t = 180$ GeV, the same type of plot would show $H^\pm$ signals that are dramatically above background at $\tan\beta = 1$. For $m_t = 110$ GeV, $H^\pm$ detection in this mode would be very difficult. The magnitude of the signal for an $H^\pm$ of any given mass depends significantly on $\tan\beta$. Assuming model-II couplings, as $\tan\beta$ increases above 1 the signals decrease to a minimum level at $\tan\beta \sim 5-6$ rising back
Figure 5: $dN/dM_{b\bar{b}j}$ is plotted as a function of $M_{b\bar{b}j}$ for: the $gg \rightarrow b\bar{t}H^+ + \bar{t}tH^-$ signal (solid); the $gg \rightarrow t\bar{t}b\bar{b}$ background (dots); and the $t\bar{t}g$ mis-tagged background (dashes). For this plot we have taken $\tan\beta = 1$ and $m_t = 140$ GeV and integrated luminosity of $L = 10$ fb$^{-1}$ at the SSC. Signal curves are given for $m_{H^+} = 180, 200, 300, 400$ and 500 GeV. Results do not include any QCD K-factors for the $tbH^\pm$ signal or $t\bar{t}b\bar{b}$ background. No additional K-factor for the $t\bar{t}g$ background is appropriate.

up to the $\tan\beta = 1$ level for $\tan\beta \gtrsim 20$. For $\tan\beta$ values above this level, it should be possible to detect the $H^+$ in this manner if $m_t \gtrsim 140$ GeV. Since in many GUT scenarios the large value of $m_t$ is correlated with a large $\tan\beta$ value, this mode holds considerable promise. More details on precise statistical significances etc. will appear in a forthcoming paper. Of course, it should be noted that the above discussion is largely independent of whether or not the $H^+$ is part of the MSSM Higgs sector, or simply the charged member of a more general Higgs doublet with couplings to fermions of model-II type.

3.6. Conclusions for the MSSM

We have made some remarkable strides in MSSM Higgs detection in the last
With the addition of the \( t\bar{t}b\bar{b} \) and \( t\bar{t} + \text{invisible} \) detection techniques, the no-lose theorem has no remaining loop-holes. Detection of one of the MSSM Higgs bosons (most likely the \( h^0 \)) will be possible at the SSC for all possible MSSM parameter choices. Further, progress on detection of the heavier Higgs bosons of the MSSM continues.

4. A Higgs Boson with Negligible Fermionic Couplings

It is not impossible that the symmetry breaking mechanism responsible for giving masses to the gauge bosons is separate from that which generates the fermion masses. In this case, there could exist a neutral Higgs boson, \( h \), which couples at tree-level only to gauge bosons and not to fermions. Such a neutral Higgs boson can also arise in the context of triplet Higgs representations.\(^{[1]}\) There, the triplet Higgs bosons do not have the appropriate quantum numbers to allow tree-level fermion couplings, but they can couple to the gauge bosons. In both cases, fermion couplings will be generated by one-loop diagrams. Fine-tuning is required to enforce small fermion couplings to all orders in perturbation theory. Nonetheless, it is amusing to consider the phenomenology of a neutral \( h \) with only tiny \( f\bar{f} \) couplings. Also important is the phenomenology forced upon other members of the Higgs sector by choosing parameters such that the \( h \) be of this special character.

Turning first to the phenomenology of the \( h \), we note that, if such an \( h \) has large enough mass, its decays to \( WW^* \) channels will certainly be dominant, and the phenomenology for the \( h \) would not be very different from that of the \( \phi^0 \) (at the same mass). But, far enough below \( WW^* \) threshold the \( \gamma\gamma \) decay of the \( h \) (as induced by the \( W \) loop diagram) will be dominant, rather than \( f\bar{f} \) channels. The exact value of \( m_h \) below which \( \gamma\gamma \) decays of the \( h \) begin to dominate is model dependent, but is generally below 100 GeV. Since the \( ZZ \) coupling of the \( h \) is by assumption not suppressed, if the \( h \) were light enough it could be produced with significant rate in the \( Z \to Z^+ h \) channel at LEP and there would be an excess of \( \nu\bar{\nu}\gamma\gamma \), \( \ell^+\ell^-\gamma\gamma \) and \( q\bar{q}\gamma\gamma \) events there. If we assume that there is no such excess, then most likely \( m_h \gtrsim 60 \) GeV.

We have already noted that the signature for an \( h \) with mass in the 60 – 130 GeV range that decays primarily to \( \gamma\gamma \) would be a truly spectacular enhancement in the \( Wh \to \ell\gamma\gamma X \) channel. There would be no difficulty in seeing such an \( h \) at the SSC. In addition, there would be a not-insignificant inclusive \( \gamma\gamma \) signal coming from \( WW \) fusion production of the \( h \) followed by \( h \to \gamma\gamma \) decay. Even though the \( WW \) fusion cross section is normally neglected in comparison to \( gg \) fusion when discussing a light \( \phi^0 \), the \( WW \) fusion production rate is only about a factor of 10 lower. If the \( \gamma\gamma \) branching ratio of the \( h \) is enhanced by a factor of 10 to 1000 compared to the SM \( \phi^0 \), the signal in the \( qg \to qgW^+W^- \to qq\gamma\gamma \) channel would be easily visible. Ref. 27 has examined the feasibility of detecting the \( h \) in the \( Wh \) mode even at the Tevatron. Their conclusion is that with \( L = 100 \) pb\(^{-1} \), the \( W\gamma\gamma \) channel can be used to detect such an \( h \) for \( m_h \lesssim 100 \) GeV.

What about other Higgs bosons associated with the \( h \). This is a model-dependent
issue. An investigation is reported in one contribution,\(^{[54]}\) in the case where the \(h\) is required to be part of a two-doublet model of type-I. If the Higgs potential is to realize the above described scenario for the lighter of the two neutral Higgs bosons (i.e. \(h = h^0\)) naturally, then \(\beta\) must be small, and the neutral sector mixing angle \(\alpha\) must be near \(-\pi/2\). Thus, the couplings of the \(H^0\) and \(A^0\) to fermions (\(\propto 1/\sin \beta\)) will be enhanced. These fermion couplings enter into the loop computations of both the \(gg\) and the \(\gamma\gamma\) couplings of the \(H^0\) and \(A^0\). The fermionic coupling enhancement is such, for instance, that below \(t\bar{t}\) threshold the \(H^0\) has a \(\gamma\gamma\) branching ratio that generally exceeds \(10^{-3}\), while the \(A^0 \rightarrow \gamma\gamma\) branching ratio rises from \(\sim 10^{-4}\) at \(m_{A^0} = 80\ \text{GeV}\) to above \(10^{-3}\) for \(m_{A^0}\) just below \(2m_t\).

The result is that the ratio (as computed at the SSC)

\[
R_{H^0,A^0} \equiv \frac{\sigma(gg \rightarrow H^0, A^0 \rightarrow \gamma\gamma)}{\sigma(gg \rightarrow \phi^0 \rightarrow \gamma\gamma)} \quad (3)
\]

is always larger than 1 for the \(H^0\) (substantially so for \(m_{H^0}\) below 100 GeV), while \(R_{A^0}\) can be almost 1 in the \(m_{A^0} < 120\ \text{GeV}\) range. (Both \(R_{A^0}\) and \(R_{H^0}\) rise to values substantially above 1 for masses above 150 GeV.) Meanwhile, the \(A^0\) and \(H^0\) widths remain well below 1 GeV for masses below 150 GeV. Since the \(\gamma\gamma\) inclusive mode is viable (in those detectors with \(\sim 1\%\ \gamma\gamma\) mass resolution) for a \(\phi^0\) with \(80 < m_{\phi^0} < 130\ \text{GeV}\), it is clear that both the \(H^0\) and the \(A^0\) of this model would be detectable over this same mass range. Indeed, because \(R_{H^0}\) becomes large for low \(m_{H^0}\), \(H^0\) detection would be possible well below the \(\sim 60 - 80\ \text{GeV}\) lower limit of the inclusive \(\gamma\gamma\) mode that applies in the case of the \(\phi^0\).

The enhanced fermionic couplings of the \(H^0\) and \(A^0\) also imply that the production processes \(gg \rightarrow t\bar{t}H^0\) and \(gg \rightarrow t\bar{t}A^0\) will be at least an order of magnitude larger than for the SM \(\phi^0\). For \(m_{A^0}\) and \(m_{H^0}\) below \(2m_t\), the above quoted \(\gamma\gamma\) branching ratios of the \(H^0\) and \(A^0\) are then such that the \(\ell\gamma\gamma\) channel will provide signals for the \(H^0\) and \(A^0\) that are at least as good as those obtained for the SM \(\phi^0\), and often much better. Thus, the \(\ell\gamma\gamma\) and \(\gamma\gamma\) channels each have the potential of allowing discovery of all of the neutral Higgs bosons of such a model.

For \(m_{H^0}\) and \(m_{A^0}\) above \(2m_t\), the \(H^0\) and \(A^0\) decay primarily to \(t\bar{t}\). The enhanced \(t\bar{t}H^0\) and \(t\bar{t}A^0\) cross sections imply large, and probably observable, signals for both the \(H^0\) and \(A^0\) in the \(t\bar{t}t\bar{t}\) final state at the SSC.

In general, this model illustrates the fact that a Higgs sector which differs significantly from those of the SM or MSSM may well be much more easily explored. The SM and MSSM Higgs sectors could not have been more cunningly constructed if the goal were to make Higgs detection at hadron colliders as challenging as possible.
5. Conclusions

The techniques for probing the Higgs or EWSB sector at the SSC continue to evolve and expand. New modes and refinements give us great confidence that the SSC will be a powerful tool for revealing the Higgs bosons associated with the mechanism for mass generation. In addition to having demonstrated that we can either discover the Higgs boson of the Standard Model or detect strong $WW$ interactions if the SM Higgs is heavy, we have now clearly established a no-lose theorem for the Higgs sector of the Minimal Supersymmetric Model. At least one of the MSSM Higgs bosons can be found at the SSC, regardless of model parameter choices. These are especially important milestones given that a Higgs sector that differs substantially from those of the SM and MSSM is often far more easily probed, using a selection of the same production and detection procedures as developed for the SM and MSSM.

Experimentally, two crucial items arise over and over again. First, it is vital that the detectors continue to stress excellent mass resolution in the $\gamma\gamma$ channel and excellent jet/photon discrimination. Secondly, if there is a single experimental lesson to be drawn from the most recent theoretical efforts, it is the great importance that should be attached to the construction of a $b$-jet vertex detector and the development of associated $b$-identification algorithms that together yield the highest possible efficiency and purity for $b$-tagging.

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