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Invisible Higgs bosons at present and future colliders*

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In many extensions of the standard model the Higgs boson can have substantial *invisible* decay modes, for example, into light or massless weakly interacting Goldstone bosons associated to the spontaneous violation of lepton number below the weak scale. In this work, we first report on the model independent limits on the Higgs boson from the analysis of the present LEP samples after including the possibility of invisible decays and study the prospects for LEP II. Next, we review the detectability prospects for such invisible Higgs boson at the Next Linear Collider.

1. Introduction

Recently the LEP experiments on e^+e^- collisions around the Z peak have placed important restrictions on the mass of the standard model Higgs boson [1]

$$m_{H_{SM}} \gtrsim 60 \text{ GeV}. \tag{1}$$

There are, however, many reasons to think that there may exist additional Higgs bosons in nature. One such extension of the minimal standard model is provided by supersymmetry and the desire to tackle the hierarchy problem [2]. Another interesting motivation for an enlargement of the Higgs sector is to generate the observed baryon excess by electroweak physics [3]. This, in principle, requires $m_{H_{SM}} \lesssim 40$ GeV [4] in conflict with eq. (1). This limit can be avoided in models with new Higgs bosons [5]. These could be intimately related to the question of neutrino masses [6]. Indeed, most extensions of the minimal standard model require the presence of new Higgses to induce neutrino masses [7].

Amongst the extensions of the standard model which have been suggested to generate neutrino masses, the so-called majoron models are particularly interesting and have been widely discussed [7]. The majoron is a Goldstone boson associated with the spontaneous breaking of the lepton number. Astrophysical arguments based on stellar cooling rates constrain its couplings to the charged fermions [8], while the LEP measurements of the invisible Z width substantially restrict the majoron couplings to the gauge bosons. In particular, models where the majoron is not a singlet under the $SU(2) \otimes U(1)$ symmetry [9] are now excluded [1].

There is, however, a wide class of models motivated by neutrino physics [10] which are characterized by the spontaneous violation of a global $U(1)$ lepton-number symmetry by a singlet vacuum expectation value. Unlike the original model of this type [11], this new class of models may naturally explain the neutrino masses required by astrophysical and cosmological observations without introducing any very high mass scale. Another example of this type is provided by supersymmetric extensions of the standard model where R parity is spontaneously violated [12–15].

In any of these models with the spontaneous violation of a global $U(1)$ symmetry around (or below) the weak scale the corresponding Goldstone boson has significant couplings to the Higgs bosons, even if its other couplings are suppressed. This implies that the Higgs boson can decay, with a substantial branching ratio, into the invisible mode [10, 15, 16]

$$h \to J + J \,, \tag{2}$$

where J denotes the majoron.

Such an invisible Higgs decay would lead to events with large missing energy that could be observable at LEP and affect the corresponding Higgs mass bounds. Here, we first review how one can derive, *in a model independent way*, limits on the Higgs boson from the analysis of the present LEP samples and study the prospects for LEP II. Next, we present the detectability prospects for the Higgs boson at the Next Linear Collider[17].

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2. The simplest model

2.1. The Higgs spectra

In order to illustrate the main points, we consider the simplest model which contains, in addition to the standard model scalar Higgs doublet, a complex singlet σ carrying a nonzero vacuum expectation value $\langle\sigma\rangle$, which breaks a global symmetry. The scalar potential is given by [10, 16]

$$V = \mu_\phi^2 \phi^\dagger\phi + \mu_\sigma^2 \sigma^\dagger\sigma + \lambda_1(\phi^\dagger\phi)^2 + \lambda_2(\sigma^\dagger\sigma)^2 + \delta(\phi^\dagger\phi)(\sigma^\dagger\sigma) . \tag{3}$$

Terms like σ^2 are omitted above in view of the imposed $U(1)$ invariance under which we require σ to transform nontrivially and ϕ to be trivial. Let $\sigma \equiv \frac{w}{\sqrt{2}} + \frac{R_2+iI_2}{\sqrt{2}}$, $\phi^0 \equiv \frac{v}{\sqrt{2}} + \frac{R_1+iI_1}{\sqrt{2}}$, where we have set $\langle\sigma\rangle = \frac{w}{\sqrt{2}}$ and $\langle\phi^0\rangle = \frac{v}{\sqrt{2}}$. The above potential leads to a physical massless Goldstone boson, namely the majoron $J \equiv \mathrm{Im}\ \sigma$, and two massive neutral scalars H_i (i= 1,2)

$$H_i = \hat{O}_{ij}\ R_j , \tag{4}$$

where $\hat{O}_{ij}$ is an orthogonal mixing matrix.

2.2. Higgs production

In order to be able to predict the production rates of these particles in e^+e^- collisions one needs to know their couplings to the Z boson. In the simplest model only the doublet Higgs boson ϕ has a coupling to the Z in the weak basis, not the $SU(2)\otimes U(1)$ singlet field σ. After diagonalizing the scalar boson mass matrix, one finds that the two CP even mass eigenstates H_i (i=1, 2) have couplings to the Z involving the mixing matrix,

$$\mathcal{L}_{HZZ} = (\sqrt{2}G_F)^{1/2} M_Z^2 Z_\mu Z^\mu \hat{O}_{i1} H_i . \tag{5}$$

Through these couplings both CP even Higgs bosons may be produced via the Bjorken process. As long as the mixing appearing in eq. (5) is $\mathcal{O}$ (1), all Higgs bosons can have significant production rates that are smaller than in the standard model by a factor $\epsilon_i^2 = \hat{O}_{i1}^2$. This is a general result which is valid for a large variety of models we are interested in.

2.3. Invisible Higgs boson decay

We now turn to the Higgs boson decay rates, which are sensitive to the details of the mass spectrum and to the Higgs potential. For definiteness we focus on the simplest potential, given in eq. (3). In this case, the width for the invisible H_i decay can be parametrized by

$$\Gamma(H \to JJ) = \frac{\sqrt{2}G_F}{32\pi} M_{H_i}^3 g_{H_iJJ}^2 , \tag{6}$$

where the corresponding couplings are given by

$$g_{H_iJJ} = \tan\beta\ \hat{O}_{i2} . \tag{7}$$

The rate for $H \to b\bar{b}$ also gets diluted compared to the standard model prediction, because of the mixing effects. Explicitly one has,

$$\Gamma(H \to b\bar{b}) = \frac{3\sqrt{2}G_F}{8\pi} M_H m_b^2\, g_{Hb\bar{b}}^2 (1-4m_b^2/M_H^2)^{3/2} , \tag{8}$$

(9)

which is smaller than the standard model prediction by the factor $g_{H_ib\bar{b}}^2$, where

$$g_{H_ib\bar{b}} = \hat{O}_{i1} . \tag{10}$$

The width of the Higgs decay to the JJ relative to the conventional $b\bar{b}$ mode depends upon the mixing angles. For this simple model it was shown [10] that in large regions of parameter space the Higgs field decays mainly invisibly to majorons and is produced without any substantial suppression relative to the standard model predictions. The same conclusion holds for other models, like the model described in the next section [12–15].

3. Another example: spontaneously broken R-Parity

3.1. The model

Here we are going to give another example of the situation described above. The global symmetry is now R-parity. In order to set up our notation we recall the basic ingredients of the model for spontaneous violation of R parity and lepton number proposed in Ref. [12]. The superpotential is given by

$$\begin{aligned} W &= h_u Q H_u u^c + h_d H_d Q d^c + h_e L H_d e^c \\ &+(h_0 H_u H_d - \varepsilon^2)\Phi + \hat{\mu} H_u H_d + h_\nu L H_u \nu^c \\ &+ h\Phi S\nu^c + M\nu^c S + m_\Phi \Phi\Phi \end{aligned} \tag{11}$$

This superpotential conserves total lepton number and R_p. The superfields $(\Phi,{\nu^c}_i,S_i)$ are singlets under $SU_2 \otimes U(1)$ and carry a conserved lepton number assigned as (0 ,−1 ,1), respectively. All couplings $h_u, h_d, h_e, h_\nu, h, M$ are described by arbitrary matrices in generation space which explicitly break flavour conservation.

As we will show in the next section these singlets may drive the spontaneous violation of R-parity[12, 13] leading to the existence of a Majoron given by the imaginary part of

$$\frac{v_L^2}{Vv^2}(v_u H_u - v_d H_d) + \frac{v_L}{V}\tilde{\nu}_\tau - \frac{v_R}{V}\tilde{\nu^c}_\tau + \frac{v_S}{V}\tilde{S}_\tau \quad (12)$$

where the isosinglet VEVs

$$v_R = \langle \tilde{\nu}_\tau^c \rangle \quad , \quad v_S = \left\langle \tilde{S}_\tau \right\rangle \quad (13)$$

with $V = \sqrt{v_R^2 + v_S^2}$, characterise R_p or lepton number breaking and the isodoublet VEVs

$$v_u = \langle H_u \rangle \quad , \quad v_d = \langle H_d \rangle \quad , \quad v_L = \langle \tilde{\nu}_{L\tau} \rangle \quad (14)$$

drive electroweak breaking and give masses to the fermions.

3.2. The existence of vacuum solutions

In this section we are going to show that the scalar potential has vacuum solutions that break R-parity. The model described by eq. (11) is a 3-generation model and as we will see some mixing among generations is needed for consistency of the model But for the analysis of the scalar potential we are going to consider, for simplicity, a 1-generation model.

Before we write down the scalar potential we need to specify the soft breaking terms. We write them in the form given in the spontaneously broken $N=1$ supergravity models, that is

$$\begin{aligned} V_{soft} = \ & \tilde{m}_0 \left[-A h_0 \Phi H_u H_d - B\varepsilon^2 \Phi \right. \\ & \left. + C h_\nu \tilde{\nu}^c \tilde{\nu} H_u + D h \Phi \tilde{\nu}^c S + h.c. \right] \\ & + \tilde{m}_u^2 |H_u|^2 + \tilde{m}_d^2 |H_d|^2 + \tilde{m}_L^2 |\tilde{\nu}|^2 \\ & + \tilde{m}_R^2 |\tilde{\nu}^c|^2 + \tilde{m}_S^2 |S|^2 + \tilde{m}_F^2 |\Phi|^2 \end{aligned} \quad (15)$$

At unification scale we have

$$\begin{aligned} C = D = A \quad &; \quad B = A - 2 \\ \tilde{m}_u^2 = \tilde{m}_d^2 &= \cdots = \tilde{m}_0^2 \end{aligned} \quad (16)$$

At low energy these relations will be modified by the renormalization group evolution. For simplicity we take $C = D = A$ and $B = A - 2$ but let[2] $\tilde{m}_u^2 \neq \tilde{m}_d^2 \neq \cdots \neq \tilde{m}_0^2$. Then the neutral scalar potential is given by

$$\begin{aligned} V_{total} \ &= \frac{1}{8}\left(g^2 + g'^2\right)\left[|H_u|^2 - |H_d|^2 - |\tilde{\nu}|^2\right]^2 \\ & + |h\Phi S + h_\nu \tilde{\nu} H_u|^2 + |h_0 \Phi H_u|^2 \\ & + |h\Phi \tilde{\nu}^c|^2 + |-h_0 \Phi H_d + h_\nu \tilde{\nu}\tilde{\nu}^c|^2 \\ & + |h\Phi\tilde{\nu}^c|^2 + |-h_0 H_u H_d + h\tilde{\nu}^c S - \varepsilon^2|^2 \\ & + \tilde{m}_0 \left[-A\left(-h\Phi\tilde{\nu}^c S + h_0 \Phi H_u H_d \right.\right. \\ & \left.\left. - h_\nu \tilde{\nu} H_u \tilde{\nu}^c\right) + (2-A)\varepsilon^2\Phi + h.c.\right] \\ & + \sum_i \tilde{m}_i^2 |z_i|^2 \end{aligned} \quad (17)$$

where z_i stand for all the neutral scalar fields. The stationary equations are then

$$\left.\frac{\partial V}{\partial z_i}\right|_{z_i = v_i} = 0 \ . \quad (18)$$

These are a set of six nonlinear equations that should be solved for the VEVs for each set of parameters. It is important to realise that it is not enough to find a solution of these equations but it is necessary to show that it is a minimum of the potential. To find the solutions we did not directly solve eq. (18) but rather follow the following three step procedure:

- *Finding solutions of the extremum equations*

We start by taking random values for h, h_0, h_ν, A, ε^2 and $\tilde{m}_0$ in the following interesting ranges

$$\begin{aligned} -3 \leq A &\leq -3 \\ 10^{-6} \leq |h_\nu| &\leq 10^{-1} \\ 10^{-2} \leq |h|, |h_0| &\leq 1 \\ 10^3 \leq |\varepsilon^2/GeV^2| &\leq 10^6 \\ 250\ GeV \leq \tilde{m}_0 &\leq 1500\ GeV \end{aligned} \quad (19)$$

We also take random values for v_R, v_S in the range

$$10\ GeV \leq |v_R|, |v_S| \leq 1000\ GeV \quad (20)$$

[2]Notice that for $\langle H_u \rangle \neq \langle H_d \rangle$ we must have $\tilde{m}_u^2 \neq \tilde{m}_d^2$ even in MSSM.

Then we choose $\tan\beta = \frac{v_u}{v_d}$ and fix v_u, v_d by the W mass relation (put $v_L = 0$ in first approximation)

$$m_W^2 = \frac{1}{2}g^2(v_u^2 + v_d^2 + v_L^2) \tag{21}$$

Now we solve eq. (18) for v_L and v_F neglecting v_L^3 terms and in the approximation

$$\tilde{m}_u^2 = \tilde{m}_d^2 = \cdots = \tilde{m}_0^2 \tag{22}$$

Finally we solve the extremum equations *exactly* for $\tilde{m}_u^2, \tilde{m}_d^2, \ldots, \tilde{m}_0^2$. This is possible because they are linear equations on the mass squared terms. In conclusion we get a set of parameters (h_ν, h, ..., $\tilde{m}_d$, ...) and VEVs (v_d, v_u, ..., v_F) that exactly solve eq. (18).

•*Showing that the solution is a minimum*

To show that the solution is a true minimum we begin by calculating the squared mass matrices. To do this we write the weak neutal scalar fields as a vector $z_i \equiv (H_d, H_u, \tilde{\nu}, \tilde{\nu}^c, S, \Phi)$ and we set

$$z_i = v_i + \frac{1}{\sqrt{2}}(x_i + iy_i) \tag{23}$$

Then the mass squared matrices are given by

$$\begin{aligned} M^2_{Rij} &= \left[\frac{1}{2}\left(\frac{\partial^2 V}{\partial z_i \partial z_j} + c.c.\right) + \frac{\partial^2 V}{\partial z_i \partial z_j^*}\right]_{z_i=v_i} \\ M^2_{Iij} &= \left[-\frac{1}{2}\left(\frac{\partial^2 V}{\partial z_i \partial z_j} + c.c.\right) + \frac{\partial^2 V}{\partial z_i \partial z_j^*}\right]_{z_i=v_i} \end{aligned} \tag{24}$$

which obey the sum rule

$$\mathrm{Tr}\, M_R^2 = \mathrm{Tr}\, M_I^2 + m_Z^2 \tag{25}$$

After obtaining the mass matrices we evaluate numerically their eigenvalues. The solution is a minimum if all nonzero eigenvalues are positive. A consistency check is that we should get two zero eigenvalues for M_I^2 corresponding to the Goldstone boson of the Z^0 and to the majoron J.

•*Comparing with other minima*

There are three kinds of minima to which we should compare our solution. These correspond to the cases

- $v_u = v_d = v_L = v_R = v_S = 0$
  $v_F \neq 0$ (26)
- $v_L = v_R = v_S = 0$
  $v_u, v_d, v_F \neq 0$ (27)
- $v_u = v_d = v_L = 0$
  $v_R, v_S, v_F \neq 0$ (28)

These cases can be solved explicitly for each set of parameters. As a final result we found[13] a large region in parameter space where our solution that breaks R_p and $SU_2 \otimes U(1)$ is an absolute minimum.

3.3. The Higgs sector

The phenomenological implications of this model for accelerator and laboratory physics have by now been studied quite extensively []. In this review we just concentrate on the Higgs sector. First of all the charged Higgs are exactly as in the *MSSM*. This is because we also have just two Higgs doublets. On the contrary, the neutral Higgs sector is far more complicated. Here we are going to make the simplification of considering just a 1-generation model. As we have said in the previous section a phenomenologically consistent model requires the presence of flavour nondiagonal couplings in order to assure that the ν_τ decays fast enough. In this approximation the M_R^2 and M_I^2 mass matrices eq. (24) are 6×6. This gives 6 CP even states and 5 CP odd, including the massless majoron. Radiative corrections to the mass matrices have been also studied[15]. The consequences of their inclusion is similar to *MSSM* and we will not discuss them further here.

The two main points regarding the Higgs sector have to do with different production and decay rates. The production mechanisms are the Bjorken process

$$Z \to H_i f\bar{f} \tag{29}$$

and

$$Z \to H_i A_j \tag{30}$$

In the *MSSM* where $H_1 = h$, $H_2 = H$ and $A_2 = A$ (there is no majoron in this limit) there a complementarity between the two types of decay and this has been used as the basis of the experimental analysis used to place limits on the SUSY Higgs spectrum. In the present model there is a breakdown in the sum rule involving the coupling strenghts characterizing the processes in eq. (29) and eq. (30) with respect to *MSSM* expectations. As a result the overall Higgs production rates are weaker than in *MSSM* by a factor $\epsilon^2 < 1$ where

$$\epsilon = \frac{g_{ZZH}}{g^{SM}_{ZZH}} \tag{31}$$

is the ratio of the ZZH coupling of this model compared to the standard model coupling.

Regarding the decays of the Higgs, there is in the models with spontaneously broken R_p, a new Higgs decay channel with majoron emision

$$H_i \rightarrow JJ \tag{32}$$

In order to evaluate the rates we need to evaluate the corresponding trilinear couplings. This was done in[15] where the the decay channel eq. (32) was compared with the other channels. The conclusion was that normally the dominant decay mode is the invisible channel given in eq. (32) over most of the kinematical range available. Since the majoron is weakly coupled, it will escape detection. Thus the decay eq. (32) will lead to events with missing energy carried by the majorons. This has to taken in account in putting limits on the Higgs mass.

In summary, the invisible Higgs decay mode is expected to have quite important implications if there exists, as suggested by neutrino physics, a global symmetry that gets broken around the weak scale. From this point of view it is desirable to obtain limits on Higgs bosons that are not vitiated by detailed assumptions on its mode of decay. This can be done from the existing Z sample at LEP, as we will briefly review below following Ref. [18].

4. Experimental bounds from LEP I

The production and subsequent decay of any Higgs boson, which may decay visibly or invisibly, involves three independent parameters: the Higgs boson mass M_H, its coupling strength to the Z, normalized by that of the standard model, call this factor ϵ^2, and the invisible Higgs boson decay branching ratio.

One can use the results published by the LEP experiments on the searches for various exotic channels in order to deduce the regions in the parameter space of the model that is already ruled out. Here we briefly summarize the procedure used in Ref. [18] in order to obtain these limits. For each value of the Higgs mass, one calculates the lower bound on ϵ^2, as a function of the branching ratio $BR(H \rightarrow \text{visible})$. By taking the highest of such bounds for $BR(H \rightarrow \text{visible})$ in the range between 0 and 1, one obtains the absolute bound on ϵ^2 as a function of M_H.

For a Higgs of low mass (below 30 GeV) decaying to invisible particles one considers the process $Z \rightarrow HZ^*$, with $Z^* \rightarrow e^+e^-$ or $Z^* \rightarrow \mu^+\mu^-$ and combines the results of the LEP experiments on the search for acoplanar lepton pairs [19–21] which found no candidates in a total sample corresponding to 780.000 hadronic Z decays. The efficiencies for the detection of the signal range from 20% at very low Higgs masses to almost 50% for $M_H = 25$ GeV.

For higher Higgs masses the rate of the process used above is too small, and one considers instead the channel $Z \rightarrow HZ^*$, $Z^* \rightarrow q\bar{q}$. Here the results of the searches for the standard model Higgs in the channel $Z \rightarrow Z^* H_{SM}$ with $H_{SM} \rightarrow q\bar{q}$ and $Z \rightarrow \nu\bar{\nu}$ can be translated, following Ref. [22]. The efficiency of these searches for an invisible Higgs increases from 25% at $M_H = 30$ GeV to about 50% at $M_H = 50$ GeV.

For visible decays of the Higgs boson its signature is the same as that of the standard model one, and the searches for this particle can be applied directly. For masses below 12 GeV one takes the results of a model independent analysis made by the L3 collaboration (Ref. [23]). For masses between 12 and 35 GeV the results in Ref. [19, 22, 23] can be combined; finally for masses up to 60 GeV the combined result of all the four LEP experiments given in Ref. [22] can be used. In all cases the bound on the ratio $BR(Z \rightarrow ZH)/BR(Z \rightarrow ZH_{SM})$ was calculated from the quoted sensitivity, taking into account

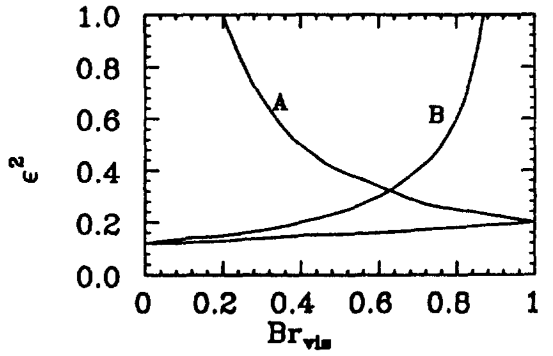

Figure 1. Exclusion contours in the plane ϵ^2 vs. $BR(H \to \text{visible})$ for the particular choice $m_H = 50$ GeV. The two curves corresponding to the searches for visible (curve A) and invisible (curve B) decays are combined to give the final bound, which holds irrespective of the value of $BR(H \to \text{visible})$.

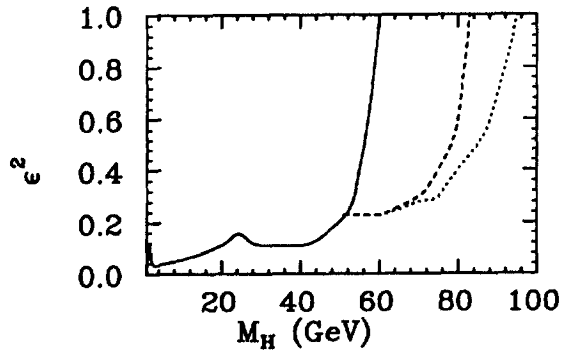

Figure 2. The solid curve shows the region in the ϵ^2 vs. m_H that can be excluded by the present LEP analyses, independent of the mode of Higgs decay, visible or invisible. The dashed and dotted curves show the region that can be explored at LEP II, for the given centre-of-mass energies.

the background events where they existed.

As an illustration, we show in Fig. 1 (from Ref. [18]) the exclusion contours in the plane ϵ^2 vs. $BR(H \to \text{visible})$ for the particular choice for the Higgs mass $M_H = 50$ GeV. The two curves corresponding to the searches for visible and invisible decays are combined to give the final bound; values of ϵ^2 above 0.2 are ruled out independently of the value of $BR(H \to \text{visible})$. The solid line in Fig. 2 shows the region in the ϵ^2 vs. M_H that can be excluded by the present LEP analyses, independent of the mode of Higgs decay, visible or invisible.

5. Prospects for LEP II

One can also estimate the additional range of parameters that can be covered by LEP II, assuming that the total luminosity collected will be 500 pb^{-1}, and for two possible values of the centre-of-mass energy: 175 GeV and 190 GeV.

The results on the visible decays of the Higgs are based on the study of efficiencies and backgrounds in the search for the standard model Higgs described in Ref. [24]. For the invisible decays of the Higgs one has considered only the channel HZ with $Z \to e^+e^-$ or $Z \to \mu^+\mu^-$, giving a signature of two leptons plus missing transverse momentum. The requirement that the invariant mass of the two leptons must be close to the Z mass can kill most of the background from WW and $\gamma\gamma$ events; the background from ZZ events with one of the Z decaying to neutrinos is small and the measurement of the mass recoiling against the two leptons allows to further reduce it, at least for M_H not too close to M_Z. Hadronic decays of the Z were not considered in Ref. [18], since the background from WW and $We\nu$ events is very large, and b-tagging is much less useful than in the search for ZH_{SM} with $Z \to \nu\bar{\nu}$, since the $Zb\bar{b}$ branching ratio is much smaller than $Hb\bar{b}$ in the standard model.

The dashed and dotted curves on Fig. 2 show the exclusion contours in the ϵ^2 vs. M_H plane that can be explored at LEP II, for the given centre-of-mass energies. Again, these contours are valid irrespective of whether the Higgs decays visibly, as in the standard model, or invisibly.

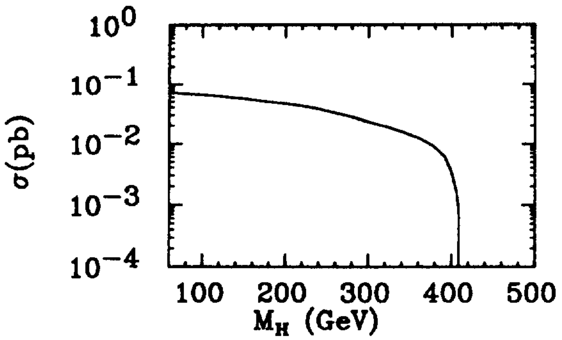

Figure 3. Total cross section for the Higgs bremstrahlung process at 500 GeV.

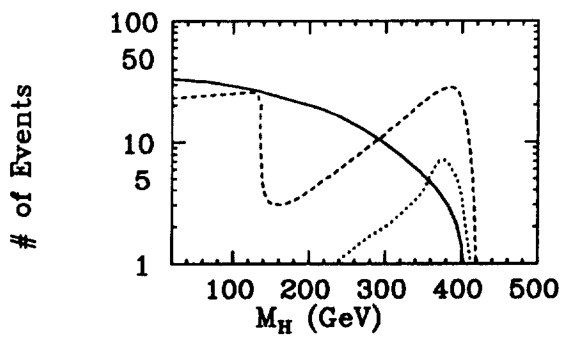

Figure 4. Final number of events for the signal for $\epsilon^2 \times Br_{invis} = 1$ (solid) and backgrounds $e^+e^- \to \nu\bar{\nu}Z$ (dashed) and WW background (dotted)

6. Invisible Higgses at the NLC

At the NLC there are two production mechanisms for Higgs particles: the Higgs bremstrahlung off the Z boson line

$$e^+e^- \to Z^* \to ZH \tag{33}$$

and the fusion process

$$e^+e^- \to \nu\bar{\nu}W^*W^* \to \nu\bar{\nu}H \tag{34}$$

For Higgses decaying invisibly, this second mechanism becomes irrelevant since it would lead to no visible signature. In Fig. 3 we plot the cross section for the Higgs bremstrahlung process as a function of the Higgs mass at $\sqrt{s} = 500$ GeV.

The main sources of background for the invisible decays of the Higgs are the processes

$$\begin{array}{lll} e^+e^- \to \nu\bar{\nu}Z & & (A) \\ e^+e^- \to WW & \to (q\bar{q}')\,[l]\,\nu & (B) \\ & \to (l\bar{l})\,\nu & (C) \\ e^+e^- \to e\nu_e W & \to (q\bar{q}')\,[e]\,\nu_e & (D) \\ & \to (e^+e^-)\,\nu_e & (E) \end{array} \tag{35}$$

where the particles in square brackets escape undetected and the fermion pairs have an invariant mass close to the Z mass. The large values of the last two total cross sections makes the WW and $e\nu_e W$ backgrounds very large for the hadronic decays of the Z even after imposing the Z invariant mass reconstruction. For this reason we will consider in our study only the leptonic decay modes of the Z, $Z \to e^+e^-$ or $Z \to \mu^+\mu^-$. The signature will be therefore two leptons with invariant mass compatible with the Z mass plus missing transverse momentum. The requirement of missing transverse momentum eliminates the $\gamma\gamma$ and $ll(\gamma)$ events as well. In this case the process $e\nu_e W$ (E) becomes irrelevant. The most dangerous background we are left with is the process A [25]. To suppress it we impose the reconstruction of the Z energy

$$E_Z(m_H) = (s + m_Z^2 - m_H^2)/(2\sqrt{s}) \pm \Delta E \tag{36}$$

We assume an energy resolution $\Delta E = 10$ GeV. Further suppression can be obtained from the fact that the Z's in the signal are produced to larger polar angles than in the background [17]. We impose an angular cut $|\cos\theta_Z| < 0.7$. After imposing these cuts the WW background (C) becomes very small (see Fig. 4.). In Fig. 4 we show the number of events we are left with for the signal and backgrounds A and C for a luminosity $\mathcal{L} = 10^4$ pb^{-1} and for $\epsilon^2 \times Br_{invis} = 1$.

In Fig. 5 we show the exclusion contours (at 95% CL) in the $\epsilon^2 \times Br_{invis}$ vs. M_H plane that can be explored at the NLC. Invisible Higgses with masses below 200 GeV can be detected if their coupling to the Z is higher than 1/3 of the

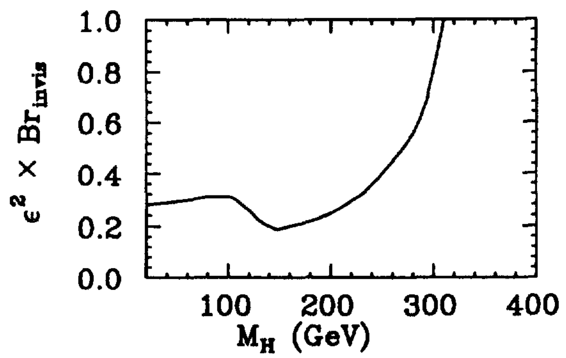

Figure 5. Accessible region at the NLC in the plane $\epsilon^2 \times Br_{invis}$, M_H at 95 % CL.

"standard" Higgs coupling; fully coupled Higgs bosons can be detected up to masses of almost 300 GeV.

Finally, we point out that at the NLC it will be possible to transform an electron beam into a photon one through the laser backscattering mechanism. This kind of process will allow the NLC to operate also in the $e\gamma$, and $\gamma\gamma$ and will provide us with new mechanisms for production and detection of an invisibly decaying Higgs particle.

7. Discussion

The Higgs boson can decay to invisible Goldstone bosons in a wide class of models in which a global symmetry, such as lepton number, is broken spontaneously around or below the weak scale. These models are attractive from the point of view of neutrino physics and suggest the need to search for the invisble mode of the Higgs boson.

We have reviewed the model-independent limits on the Higgs boson mass and Z coupling strength that can be deduced from the present LEP samples. The limits are summarized in Fig. 1 and 2 and do not depend on the mode of Higgs boson decay. They are probably conservative and could be somewhat improved with more data and/or more refined analysis.

Moreover we have investigated the reach of a high energy linear e^+e^- collider to discover a Higgs boson in the invisible mode. In Fig. 6 we show the exclusion contours (at 95% CL) in the $\epsilon^2 \times Br_{invis}$ vs. M_H plane that can be explored at the NLC. Invisible Higgses with masses below 200 GeV can be detected if their coupling to the Z is higher than 1/3 of the "standard" Higgs coupling; fully coupled Higgs bosons can be detected up to masses of almost 300 GeV.

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