Abstract

We present and study the results for the standard model process $e^+e^- \rightarrow \nu\bar{\nu}b\bar{b}$ at c.m. energies $150 \leq \sqrt{s}(GeV) \leq 240$ and for Higgs boson masses $80 GeV \leq m_H \leq 120 GeV$, obtained from all tree-level diagrams and including the most important radiative corrections. The matrix elements have been calculated by the 'spinor bracket' method without neglecting masses, while the phase space integrals by an importance sampling Monte Carlo numerical integrator. The $\sqrt{s}$ dependence and the interference properties of the Higgs boson contribution and of various coherent background contributions to the total cross section are examined. The effect of the QED initial state radiative corrections is estimated. The important differential distributions for the Higgs boson and the background components are studied, providing information useful for choosing cuts in Higgs searches. We also examine the effect of a minimal set of cuts and evaluate the importance of the WW fusion for detecting a higher mass Higgs boson at LEPII.
1 Introduction

Increasingly precise LEP I data \[1\], combined with an evidence for the top quark of mass $176 \pm 15$ GeV \[2\], have made feasible a very important—and impressively successful—test of the Standard Model (SM) to an accuracy level of one-loop quantum corrections \[3\], \[4\]. These data are also compatible with the Minimal Supersymmetric extension of the SM (MSSM \[4\], which (at a price of introducing many more particles) provides perturbative stability against radiative corrections in a grand unification scheme.

Yet, comparison to presently available data and requirements of theoretical consistency do not restrict sharply the mass of the SM Higgs boson, and even more so the parameters of the Higgs sector of the MSSM \[5\]. On the experimental side, LEP I collaborations converge slowly towards a combined lower experimental bound $m_H \geq 65$ GeV for the mass of the SM Higgs boson \[1\]: given this bound, the other results of LEP I are well-known to be rather insensitive on $m_H$. As for the theoretical constraints, the radiatively corrected SM Higgs potential $V[\phi]$ develops an instability \[3\], and depending on the assumptions about the cure of this quantum catastrophe, some limits on $m_H$ are implied. E.g. assuming that the instability is avoided perturbatively by onset of new physics at $\Lambda \approx 10^3$ GeV (respectively $\Lambda \approx 10^4$ GeV) implies, for $m_t = 174$ GeV and $\alpha_s(m_z) = 0.118$, a lower limit $m_H > 75$ GeV (respectively $m_H > 100$ GeV) \[7\]. A change $\Delta m_t$ of the top quark mass induces a change $\Delta m_H \approx 2\Delta m_t$ of such a limit \[6\]. On the other hand, assuming the MSSM with $m_t \approx 170$ GeV and supersymmetry breaking scale $m_s \approx 10^3$ GeV implies an upper limit $m_H < 125$ GeV on the lightest mass Higgs boson, the limit changing with $m_t$ according to $\Delta m_H \approx \Delta m_t$ \[8\]. The same upper bound (for $m_t \approx 170$ GeV) was found from a low energy renormalization group analysis of the Higgs sector of the ‘next-to-minimal’ supersymmetric SM \[9\]. Actually, because of the proximity of a fixed point in the low energy evolution of $m_t$, one would expect that $m_H$ is well below this upper bound, expecting e.g. $m_H \leq 105$ GeV for the above values of $m_t$ and $m_s$ \[10\].

Detection of at least the lightest Higgs boson (HB) and experimental study of its properties is presently among the highest priority goals in particle physics. Yet, LEP II will provide the only opportunity in the near future to search directly for a HB of mass $m_H > 65$ GeV. It is thus very important to have complete and carefully analysed results for all $e^+e^-$-collision SM and MSSM processes with an ‘uncoverable’ HB signal at 170-210 GeV c.m. energies. For the first time many of these processes will have 4 fermion final states, arising partly through 2 intermediate heavy bosons, and we should learn to optimize the search for the HB in such processes and to push the explorable $m_H$ range as high as possible. A precise knowledge of the experimental background to these processes is, of course, also much needed.

In this work we assume the validity of the SM at LEP II energies. Since the SM Higgs boson detectable at LEP II decays predominantly to $b\bar{b}$, the most promising SM processes are

$$e^+e^- \rightarrow \ell^+\ell^-b\bar{b}, \ \nu\bar{\nu}b\bar{b}, \ q\bar{q}b\bar{b}$$
where $\ell = e$ or $\mu$, summation over the neutrino species is understood, and $q\overline{q}$ is a quark-antiquark pair generating jets just like the $b\overline{b}$ pair does. In the present study we concentrate on the complete tree level calculation, including approximatively the most important radiative corrections, of the process

$$e^+e^- \rightarrow \nu\overline{\nu}b\overline{b} \quad (1)$$

and examine some problems of extracting the HB signal from it at LEP II despite a large background. We should mention that process (1) is of interest at LEP II for another reason: it will constitute an important background in the search for the process $e^+e^- \rightarrow \chi\chi b\overline{b}$, where $\chi$ would be the lightest neutralino, stable when R-Parity is respected. Both $\nu$'s and $\chi$'s (if they exist) will be ‘felt’ only through energy-momentum conservation.

An intermediate SM Higgs boson contributes to process (1) through the so-called Higgsstrahlung and $WW$ fusion diagrams of Fig. 1. This HB signal is embedded in a coherent electroweak background arising from the 20 diagrams in Fig. 2. These can be classified as

- the $e$-exchange diagrams of Fig. 2a, where $ZZ$ or $Z\gamma$ coexist and cascade on the $s$-channel
- the $WW$ multiperipheral and fusion diagrams of Fig. 2b
- the $W$- and $e$-exchange diagrams of Fig. 2c, with a radiation of $\gamma$ or $Z$
- the $s$-channel $e^+e^-$ annihilation diagrams of Fig. 2d, with $Z$ radiation from a final fermion.

For reasons of brevity, the background processes of Fig. 2a, b, c, d will be called respectively $ZZ/\gamma$ cascade, $WW$ fusion, $W$ exchange and $Z$ radiation processes. We work in the t’ Hooft-Feynman gauge, but have ignored diagrams with a coupling of the Higgs boson or of Goldstone bosons to $e^\pm, \nu_e, \overline{\nu}_e$ leptons, since they contribute insignificantly.

Experimentally, one must extract the HB signal from all final states consisting of a pair of imperfectly $b$-tagged jets + missing products. The important components of the incoherent background, denoted $B_{QCD}^\gamma, B_{WW}, B_W, B_{\ell\ell}$ and $B_{\mu\mu}$ for brevity, correspond to the following processes:

- $B_{QCD}^\gamma$: $s$-channel (virtual $\gamma/Z$ mediated) production of $b\overline{b}$ or $c\overline{c}$ jets (there is the possibility of misidentification of $c$ jets in $b$-tagging), accompanied by soft undetected hadrons or by beam radiated photons escaping detection.
- $B_{WW}$: $WW$ production, with one $W$ decaying to $cs$ jets and the other decaying undetected through a chain $W \rightarrow \nu\tau \rightarrow \nu\overline{\nu}+\text{soft charged particles}$.
- $B_W$: $We\nu$ events, with $e$ escaping undetected and $W$ decaying into $cs$ jets.
- $B_{\ell\ell}$, $\ell = e$ or $\mu$: $\ell^+\ell^-b\overline{b}, \ell^+\ell^-c\overline{c}$ final states, with the leptons escaping in the forward and/or backward direction.
It is evident that very efficient criteria of event acceptance must be invoked in order to isolate (if possible) the Higgs component at a detectable level.

In order to put the present study in a proper context, we recall briefly previous work on $e^+e^- \rightarrow \nu \nu b\bar{b}$. Starting from [11], there have been several ambitious simulations of this process by experimental groups, aiming at a realistic estimate of the SM Higgs boson signal to background ratio at LEP II (e.g. [12],[13]). Yet, they ignored the WW fusion diagram of Fig. 1, assuming that it was important only at higher energies. But as shown by our results, WW fusion is already important at LEP II energies for two reasons: (a) it increases the signal by about 10% around its maximum and (b) it increases by large factors the weak signal around the kinematical threshold for HZ production. They also ignored parts of the electroweak background and evaluated the rest approximately, invoking only 2 $\rightarrow$ 2 body processes followed by appropriate decays. Theorists, on the other hand, were contended to study $e^+e^- \rightarrow \nu \nu H$ without quantitative studies of the background. They usually included the WW fusion contribution to this 2 $\rightarrow$ 3 body process [14],[15],[16], but stressed its importance only for $\sqrt{s} \geq 500$ GeV [17],[18]. It was only recently [19],[20] when our present work was well under way, that the process $e^+e^- \rightarrow \nu \nu b\bar{b}$ was calculated including all signal and background contributions, with emphasis on results for $\sqrt{s}$ above 200 GeV.

Despite some overlap between [19],[20] and the present study, we decided to complete and publish our results for several reasons. We shall not dwell on a trivial one, the need to check independently calculations involving 22 diagrams and multidimensional phase space integrations, since we find good agreement when calculating same things with same parameters. More importantly, in view of the necessity of event selection criteria for improving the signal/background ratio, we put emphasis on studying the detailed characteristics of the signal and the coherent background, including some differential cross sections not considered in [19],[20]. This information can be beneficially combined with experience about $B_{QCD}^\nu$, $B_{WW}$, $B_W$ and $B_{ll}$ to choose a set of experimental cuts aiming to extract a measurable HB signal. We also introduce several quantitatively important technical improvements over previous work (approximative inclusion of the QED initial state radiative (ISR) corrections and of important $QCD$ radiative corrections to the adopted HB width and the Higgs–$b\bar{b}$ coupling, approximative employment of an Improved Born Approximation), and examine the effect of a set of minimal cuts. These results allow estimates of the SM Higgs boson detectability at LEP II more realistic than previously. As a by-product, we reevaluate the importance of the $W-W$ fusion for detecting a higher mass HB at LEP II, already noticed in [19],[20]. We should mention once more that, for reasons of experimental accessibility, all cross section presented below are summed over neutrino species, (in [20] cross sections with electronic neutrinos are emphasized).
2 II. Method of calculation and computation

The process $e^+e^- \rightarrow \nu\bar{\nu}b\bar{b}$ interrelates six external spin $\frac{1}{2}$ fermions. For fixed momenta of these fermions, the initial-spin averaged and final-spin summed matrix element squared is given by

$$\frac{1}{4} \sum_{\lambda} \left| \sum_F iM^F_{\lambda} \right|^2$$

where $\lambda$ runs over a basis of spin states for external fermions and $F$ runs over all Feynman diagrams contributing to the process. Being interested in the Higgs signal, we must avoid approximating the $b$ quark as massless: so we decided not to use this approximation for the electrons either. We then choose a convenient spinor basis in accordance with [21], [22] and calculate analytic expressions for all $iM^F_{\lambda}$: the remaining operations in (2) are performed, for each point of phase space, by a computer program. (The advantages of this procedure in comparison to the traditional ‘trace technique’ are discussed in [21], [22]). Before describing the structural form of $iM^F_{\lambda}$, we would like to briefly motivate (in a novel way) and present the spinor basis used, since it deserves —due to its usefulness— to be more widely understood and used. Its basic advantage is to extend smoothly to massive fermions the natural spin basis for massless fermions.

A well-known covariant way of characterizing spin states is by invoking the translation— invariant Pauli-Lubanski axial 4-vector operator [23].

$$W^\mu = -\frac{1}{2} \varepsilon^{\mu\nu\lambda\rho} \Sigma_{\nu\lambda} \hat{P}_\rho, \quad (\hat{P}_\mu = i\partial_\mu)$$

$$[W^\mu, \hat{P}_\mu] = 0, \quad [W^\mu, \hat{P}^\nu] = 0 \quad \text{(3)}$$

$$[W^\mu, W^\nu] = i\varepsilon^{\mu\nu\lambda\rho} \hat{P}_\lambda W_\rho, \quad [W^2, W^\nu] = 0 \quad \text{(4)}$$

where $\Sigma_{\mu\nu}$ are the spin-parts of Lorentz generators, $W^2 = W_\mu W^\mu$, and we use the conventions $\varepsilon^{0123} = 1$ and $g_{00} = 1, g_{11} = g_{22} = g_{33} = -1$. $\hat{P}_\mu$ gives $\pm P_\mu$ when acting on $e^{\pm iP_\mu,\chi}$, i.e. on particle (antiparticle) solutions of 4-momentum $P_\mu$. For $P_\mu$ timelike (massive particle), $W^\mu$ is spacelike and reduces in the particle rest-frame to the spin 3-vector

$$W^\mu = (0, \pm m\Sigma) \quad \text{for} \quad P^\mu = (m, 0)$$

while for $P_\mu$ lightlike (massless particle) $W^\mu$ is lightlike ‘parallel’ to $P^\mu$, $W^\mu = \pm h P^\mu$, where the helicity $h = \hat{P} \cdot \hat{W} / (P^0)^2$ is a Lorentz-invariant characteristic of a massless particle. According to (3), (4) we can choose simultaneous eigenstates of $\hat{P}_\mu$, $W^2$ and of the projection of $W^\mu$ along some (eventually $P_\mu$-dependent) 4-vector $S^\mu$: a spin basis is characterized by the choice of $S^\mu$. This choice is best understood if we complement $P^\mu$ by three independent 4-vectors so as to form a complete basis of Lorentz 4-vectors convenient for expressing $W^\mu$.

A $P^\mu$ either lightlike or timelike can be complemented by two unit spacelike 4-vectors $e^\mu_{(1)}$, $e^\mu_{(2)}$ such that they form an orthogonal sub-basis ($a \cdot b$ and $a^2$ mean $a_\mu b^\mu$ and $a_\mu a^\mu$ respectively)

$$P \cdot e_{(1)} = P \cdot e_{(2)} = e_{(1)} \cdot e_{(2)} = 0, \quad e_{(1)}^2 = e_{(2)}^2 = -1 \quad \text{(5)}$$
The choice of the fourth basic 4-vector useful for both \( P^2 = 0 \) and \( P^2 > 0 \) is motivated as follows.

Since for massless particles the helicity \( h \) (identical for such particles to the chirality) is a Poincaré invariant property conserved by vector and axial-vector interactions, we choose eigenstates of \( h \) as a spinor basis for massless fermions. Thus we choose \( S^\mu = P^\mu \) for such particles. Complementing now \( P^\mu \), \( e_{(1)}^\mu \), \( e_{(2)}^\mu \) by a lightlike 4-vector \( k^\mu \) such that

\[
k \cdot e_{(1)} = k \cdot e_{(2)} = k^2 = 0, \quad P \cdot k \neq 0
\]

we can write

\[
W^\mu = \pm h P^\mu = \frac{W \cdot k}{P \cdot k} P^\mu
\]

a form convenient for generalization to massive particles.

Turning now to massive particles, \( P^\mu \) becomes timelike and \( W^\mu \) spacelike \((W^2 = -m^2 \Sigma^2)\). Expressing such a \( W^\mu \) in our Lorentz basis we find (using \( W \cdot P = 0 \) and (5), (6))

\[
W^\mu = \frac{W \cdot k}{P \cdot k} \left( P^\mu - \frac{m^2}{P \cdot k} k^\mu \right) - (W \cdot e_{(1)}) e_{(1)}^\mu - (W \cdot e_{(2)}) e_{(2)}^\mu
\]

Thus choosing for massive particles

\[
S^\mu = P^\mu - \frac{m^2}{P \cdot k} k^\mu
\]

we obtain a spin-characterization which connects smoothly the massive to the massless case: as \( P^2 = m^2 \rightarrow 0 \), \( W^\mu \rightarrow \pm h P^\mu \), \( S^\mu \rightarrow P^\mu \) and (8) \( \rightarrow \) (7). This \( S^\mu \) satisfies

\[
S^2 = -m^2, \quad S \cdot P = 0.
\]

If \( k^\mu \) transforms to \( \ell^\mu \) when Lorentz-boosting to the rest frame of the particle, \( S^\mu \) transforms simultaneously to \((0, -m \hat{\ell})\), where \( \hat{\ell} \) is the unit 3-vector along \( \ell \), while \( W^\mu \) transforms to \((0, \pm m \hat{\Sigma})\). Hence

\[
\frac{W \cdot k}{P \cdot k} = -\frac{W \cdot S}{m^2} = \pm \hat{\Sigma} \cdot \hat{\ell}
\]

i.e. our spinor basis for a massive particle corresponds to eigenstates of the projection of the spin \( \hat{\Sigma} \) along the unit vector \( \hat{\ell} \) in the rest frame of the particle.

The final result of this procedure is to use, for both massless and massive fermions, as spinor basis the eigenstates of \( \frac{W \cdot k}{P \cdot k} \), where \( k^\mu \) is some fixed lightlike 4-vector such that \( P \cdot k \neq 0 \), and \( P^\mu \), \( W^\mu \) are the 4-momentum and the Pauli-Lubanski 4-vector of each particle. It is useful to notice that for spin \( \frac{1}{2} \) fermions

\[
-\frac{1}{2} \epsilon^{\mu\nu\lambda\rho} \Sigma_{\nu\lambda} = i \gamma_5 \Sigma^{\mu\rho}
\]
(with the convention $\gamma_5 = i\gamma^0\gamma^1\gamma^2\gamma^3$), so that

$$W \cdot k = -\frac{1}{2}\gamma_5(\not{k} P - \not{P} \not{k}) = \gamma_5(P \not{k} - \not{P} \cdot k)$$

(13)

using $\Sigma^{\mu\nu} = \frac{i}{2}[\gamma^\mu, \gamma^\nu]$ and the notation $\phi = a_\mu \gamma^\mu$ for any 4-vector $a_\mu$. In this case our spin-characterizing operator can be written

$$\frac{W \cdot k}{P \cdot k} = \pm \left(\frac{P \not{k}}{P \cdot k} - 1\right) \gamma_5$$

(14)

where $+$ ($-$) corresponds to positive (negative) frequency solutions, i.e. to particle (antiparticle) wavefunctions $u(p)$ ($v(p)$) respectively.

Denoting the eigenvalues of $\frac{W \cdot k}{P \cdot k}$ by $\lambda/2$ ($\lambda = \pm 1$), let $u_\lambda(p, m)$, $v_\lambda(p, m)$ be the corresponding fermion and antifermion eigenstates, solutions of the Dirac equation with 4-momentum $p_\mu$. It is easy to see that, for both massive ($m > 0$) and massless ($m = 0$) particles, all the above requirements and the normalization-fixing relations

$$\sum_\lambda u_\lambda(p, m) \overline{u}_\lambda(p, m) = \hat{p} + m, \quad \sum_\lambda v_\lambda(p, m) \overline{v}_\lambda(p, m) = \hat{p} - m$$

are satisfied by the choice \[21\]

$$u_\lambda(p, m) = \frac{1}{\eta}(\hat{p} + m) \nu_\lambda(k), \quad v_\lambda(p, m) = u_{-\lambda}(p, -m)$$

(16)

where $\eta = \sqrt{2p \cdot k}$ and

$$\not{k} \nu_\lambda(k) = 0, \quad \nu_\lambda(k) \overline{\nu}_\lambda(k) = \frac{1}{2}(1 + \lambda \gamma_5)$$

(17)

i.e. $\nu_-(k), \nu_+(k)$ are the helicity eigenstates of a massless spin $\frac{1}{2}$ fermion having the reference 4-momentum $k_\mu$. We can take e.g. $\nu_+(k) = \not{e} \nu_-(k)$ where $e^\mu$ is any 4-vector satisfying $e^2 = -1$, $e \cdot k = 0$. One should notice that (17) corresponds to a normalization

$$\overline{\nu}_\lambda(k) \gamma^\mu \nu_\lambda(k) = 2k^\mu$$

(18)

which, in collaboration with the $1/\eta$ factor in (16), leads to

$$\overline{\nu}_\lambda(p, m) u_\lambda(p, m) = 2m, \quad \overline{\nu}_\lambda(p, m) v_\lambda(p, m) = -2m.$$ 

(19)

According to our spin specification

$$u_\lambda(p, m) \overline{\nu}_\lambda(p, m) = \frac{1}{2} [1 + \lambda \gamma_5 \hat{e} (3)] (\hat{p} + m)$$

$$= \frac{1}{2} \left[ 1 + \lambda \gamma_5 \left(1 - \frac{m}{p \cdot k} k\right) \right] (\hat{p} + m)$$

(20)
where $e^\mu_{(3)}$ is the dimensionless unit spacelike 4-vector $S^\mu/m$. The corresponding expression for $\nu_\lambda(p, m)$ is obtained by changing $\lambda \to -\lambda$, $m \to -m$ in the r.h.s. of (20). Formulae (15), (16), (19) and (20) have a smooth limit for $m \to 0$, and give then the correct relations for massless fermions.

Any tree Feynman diagram is characterized by a number of uninterrupted fermion lines, each streaming a fermion nb. flow which enters and leaves the diagram. Along such a line there is a contraction of Dirac indices of a sequence of spinors, interaction vertices and numerators of fermion propagators. Vector or scalar boson propagators provide contractions of Lorentz indices of vertices belonging to different fermion lines. The essence of the ‘spinor bracket’ or ‘pseudohelicity method’ of [21], [22] is to express each such fermion line sequence in terms of few basic blocks. To exhibit the structure of each $iM^F_\lambda$ in terms of the building blocks of the present approach, substitute in the corresponding expression all numerators $\not{p} \pm m$ of fermion propagators according to (15) and write all coupling constants as $C_L P_L + C_R P_R$ in terms of the chiral projectors $P_L = \frac{1}{2}(1 - \gamma_5)$, $P_R = \frac{1}{2}(1 + \gamma_5)$. Then each $iM^F_\lambda$ becomes a sum of products, these last having as factors two types of complex scalar functions: the spinor brackets

$$\overline{\nu}_\lambda(p_1, m_1)(C_L P_L + C_R P_R)u_{\lambda_2}(p_2, m_2)$$

(whose form motivates the choice of building units, see below), and the so-called $Z$-functions

$$Z(1, 2, 3, 4) = [\overline{\nu}_\lambda(p_1, m_1)\Gamma^\mu u_{\lambda_2}(p_2, m_2)][\overline{\nu}_\lambda(p_3, m_3)\Gamma'_\mu u_{\lambda_4}(p_4, m_4)]$$

$$\Gamma^\mu = \gamma^\mu(C_L P_L + C_R P_R), \quad \Gamma'_\mu = \gamma^\mu(C'_L P_L + C'_R P_R)$$

Decomposing next $u_{\lambda}(p, \pm m)$ into chiral spinors $\alpha_{\lambda}(p)$, $\nu_{\lambda}(k)$

$$u_{\lambda}(p, \pm m) = \frac{1}{\eta} \rho_{\nu-}(k) \pm \frac{m}{\eta} \nu_{\lambda}(k) \equiv \alpha_{\lambda}(p) \pm \mu \nu_{\lambda}(k)$$

and using the so-called Chisholm identities for such spinors

$$[\overline{\alpha}_\lambda(q)\gamma^\mu \alpha_{\lambda'}(q_2)]\gamma_{\mu} = 2\delta_{\lambda\lambda'}[\alpha_{\lambda}(q_2)\overline{\alpha}_\lambda(q_1) + \alpha_{\lambda}(q_1)\overline{\alpha}_\lambda(q_2)]$$

$$[\overline{\alpha}_\lambda(q)\gamma_{\mu} \alpha_{\lambda'}(k)]\gamma_{\mu} = 2\delta_{\lambda\lambda'}[\nu_{\lambda}(k)\overline{\alpha}_\lambda(q) - \alpha_{\lambda}(q)\overline{\alpha}_{\lambda'}(k)]$$

$$[\overline{\nu}_{\lambda}(k)\gamma_{\mu} \alpha_{\lambda'}(q)]\gamma_{\mu} = 2\delta_{\lambda\lambda'}[\alpha_{\lambda}(q)\overline{\nu}_\lambda(k) - \nu_{\lambda}(k)\overline{\nu}_{\lambda'}(q)]$$

one can express straightforwardly, and once for ever [21], [22], also the $Z$-functions in terms of the following basic units (here we slightly refine previous work)

$$Y^L_\lambda(p_1, p_2) \equiv \overline{\alpha}_\lambda(p_1)P_{L,R}u_{\lambda}(p_2)$$

$$S^L_{\lambda}(p_1, p_2) \equiv \overline{\nu}_\lambda(p_1)P_{L,R}u_{\lambda}(p_2) = \overline{\alpha}_\lambda(p_1)P_{L,R}\alpha_{-\lambda}(p_2)$$

$$\mu = \frac{m}{\eta}, \quad \overline{\alpha}_\lambda(p)\nu_{-\lambda}(k) = \overline{\nu}_\lambda(k)\alpha_{-\lambda}(p) = \sqrt{2p \cdot k} \equiv \eta$$

The main properties and the explicit expressions (after adopting $\nu_+(k) = \not{\rho}_\nu(k)$) of the $Y$- and $S$-functions are

$$Y^L_\lambda(p_1, p_2) = Y^R_{-\lambda}(p_1, p_2) = Y^L_{-\lambda}(p_2, p_1) = Y^R_{\lambda}(p_2, p_1)$$
\[ Y^{L}(p_1, p_2) = \mu_1 \eta_2, \quad Y^{R}(p_1, p_2) = \mu_2 \eta_1 \]
\[ S^L_\lambda(p_1, p_2) = -S^R_\lambda(p_2, p_1), \quad S^R_\lambda(p_1, p_2) = -S^L_\lambda(p_1, p_2)^* \]
\[ S^L_+(p_1, p_2) = \frac{2}{\eta_1 \eta_2}[(p_1 \cdot k)(p_2 \cdot e) - (p_2 \cdot k)(p_1 \cdot e) + i \varepsilon(k, e, p_1, p_2)] \]
\[ S^R_+(p_1, p_2) = -S^L_+(p_1, p_2)^*, \quad S^R_-+_+(p_1, p_2) = S^L_{-+}(p_1, p_2) = 0 \]

where \( \varepsilon(k, e, p_1, p_2) = \varepsilon_{\mu\nu\lambda\rho}k^\mu e^\nu p_1^\lambda p_2^\rho, \quad \varepsilon_{0123} = -1 \). We shall denote the nonzero \( S \)-functions simply \( S_\pm(p_1, p_2) \), neglecting the upper \( L, R \) index. For completeness of presentation, the expressions for \( Z \)-functions are tabulated below. For \( k^\mu \) and \( e^\mu \) we used the convenient choice \( k^\mu = (1, 1, 0, 0), \quad e^\mu = (0, 0, 1, 0) : k^\mu \) should be non-orthogonal to all external \( p^\mu \).

\[
Z(\ +, \ +, \ +, \ +) = 2[S_+^L(p_1, p_3)S_-^L(p_4, p_2)C_R^L C_R^R + \mu_1 \mu_2 \eta_3 \eta_4 \eta_1 \eta_2 \mu_3 \mu_4 C_R^L C_L^L]
\]
\[
Z(\ +, \ +, \ -, \ -) = 2[S_+^R(p_1, p_3)S_-^R(p_4, p_2)C_R^L C_L^R + \mu_1 \mu_2 \eta_3 \eta_4 \eta_1 \eta_2 \mu_3 \mu_4 C_L^L C_R^R]
\]
\[
Z(\ +, \ -, \ -, \ +) = 2[\mu_1 \eta_2 \mu_3 \mu_4 C_R^L C_R^R + \mu_1 \mu_2 \eta_3 \mu_4 C_R^L C_L^L - \mu_1 \mu_2 \eta_3 \mu_4 C_L^L C_R^R - \mu_1 \mu_2 \mu_3 \mu_4 C_R^R C_L^R]
\]
\[
Z(\ +, \ +, \ +, \ -) = 0
\]
\[
Z(\ +, \ +, \ -, \ +) = 2\eta_2 C_R^L [\mu_3 S_+^L(p_1, p_4) C_L^L + \mu_4 S_+^R(p_3, p_1) C_R^R]
\]
\[
Z(\ +, \ -, \ +, \ +) = 2\eta_1 C_R^L [\mu_2 S_-^L(p_3, p_2) C_L^L + \mu_3 S_-^R(p_2, p_4) C_R^R]
\]
\[
Z(\ +, \ -, \ -, \ +) = 2\eta_4 C_R^R [\mu_2 S_+^L(p_1, p_3) C_R^R + \mu_1 S_+^L(p_3, p_2) C_L^L]
\]
\[
Z(\ +, \ -, \ -, \ -) = 2\eta_3 C_L^L [\mu_1 S_+^L(p_4, p_2) C_L^L + \mu_2 S_+(p_1, p_4) C_R^R]
\]

The remaining \( Z \) are obtained by exchanging simultaneously \(+ \leftrightarrow -\) and \( L \leftrightarrow R \) in the above expressions.

To recapitulate, using the described spinor basis and following the outlined calculational procedure, each Feynman amplitude \( iM^E \) was expressed as a sum of products of the above defined \( Y \)-, \( S \)- and \( Z \)-functions. The \( Z \)-functions are expressed in turn in terms of the truly basic building units, the \( S \)- and \( Y \)-functions, having simple properties and explicit expressions. The whole scheme is very convenient and allows compact programming.

We use the t’Hooft-Feynman gauge and Breit-Wigner propagators for the \( Z, W \) and \( H \) bosons. The mass of \( W \) \((m_W = 80.23 \text{ GeV})\) and the mass \((m_Z = 91.1888 \text{ GeV})\) and width \((\Gamma_Z = 2.4974 \text{ GeV})\) of \( Z \) have been taken as inputs; the width of \( W \) being calculated from the SM tree-level formula.

The HB mass was varied in the range 80-120 GeV. For each \( m_H \) the total \( \Gamma_H \) is estimated using the program kindly provided by the authors of [24], which includes the significant QCD and electroweak radiative corrections to the tree-level width formula. The widths \( \Gamma_H \) and \( \Gamma_{bb} \) found for various \( m_H \) are as follows:

In order to include approximatively the corresponding QCD correction to the \( Hb\bar{b} \) vertex in the Feynman amplitudes, we employ in the \( Hb\bar{b} \) coupling a running mass \( \bar{m}_b(m_H^2) \) which, when used in the tree-level formula for \( \Gamma_{bb} \), reproduces the above stated widths. In the range
| $m_H$ (GeV) | 80 | 90 | 100 | 110 |
|------------|----|----|-----|-----|
| $\Gamma_H$ (MeV) | 1.734 | 1.911 | 2.106 | 2.393 |
| $\Gamma_{bb}$ (MeV) | 1.513 | 1.661 | 1.806 | 1.948 |

of masses considered here, $\bar{m}_b(m_H^2)$ differs from the QCD running mass at $Q^2 = m_H^2$ by a factor 1.12 ±0.01. The value of $a_S(m_Z^2)$ was taken 0.123. The kinematical quark masses needed in various parts of our calculations are taken to be $m_b = 4.7$, $m_c = 1.45$ and $m_t = 175$ GeV, while the KM matrix element $U_{tb} = 1$. The gauge couplings were fixed by adopting $\alpha_{\text{eff}} = 1/128$, $(1/137.0359895)$ at $Z(\gamma)$ vertices and using the effective $g^2 = 4\sqrt{2} G_F m_W^2$ ($G_F=1.16639565 \times 10^{-5}$). These effective parameters partly include EW radiative corrections in our tree-level amplitudes, in the spirit of the ‘Improved Born Approximation’ [26].

The QED ISR corrections were calculated approximatively using the universal radiation factor, incorporating virtual and soft contributions up to second order in $\alpha$, with infrared photon exponentiation, and a hard bremsstrahlung contribution [27]. Our cross sections should not be sensitive to final state radiation, because $\nu\bar{\nu}$ do not radiate and $b\bar{b}$ are massive enough.

Denoting final $\nu, \bar{\nu}, b, \bar{b}$ as 3, 4, 5, 6 respectively, their Lorentz-invariant phase space volume element $d\text{Lips} (s; p_3 p_4 p_5 p_6)$ was decomposed into a product of effective 2-body elements

$$d\text{Lips} (s; p_3 p_4 p_5 p_6) = d\text{Lips} (s_{34}; p_3 p_4) d\text{Lips} (s_{56}; p_5 p_6) \frac{ds_{34}}{2\pi} \frac{ds_{56}}{2\pi}$$

where

$$p_{ij}^\mu = p_i^\mu + p_j^\mu, \quad s_{ij} = E_{ij}^2 - \vec{p}_{ij}^2$$

and each 2-body $d\text{Lips}$ was evaluated in its c.m. frame. Smoothing the integrand by well-known changes of variables [28], the phase space integration was performed by the event generation and Monte Carlo integration program VEGAS [29], with estimated accuracy better than 1%.

The fortran program exists also in a MC generator form, where it is interfaced to JETSET [34] and so after hadronization produces final state particles. The MC generator is controlled by a set of cards where one can determine the HB mass, the c.m energy, the inclusion or not of radiative corrections, possible cuts on $m_{\nu\bar{\nu}}$ and $m_{bb}$, the number of requested events, and whether one wishes to write the events in LUND format on some separate file. One can restrict, through switches in the same cards, the calculation to a subclass of diagrams and to the production of $\nu_\mu$ or $\nu_\tau$ only. Finally, all input constants used in the calculations: $m_Z, m_W, \Gamma_Z, G_F, \alpha$ (at 0 and $M_Z$), $\alpha_S$ and the quark kinematical masses, are given from the outside on the same set of cards, providing to the user complete control of the input variables for comparison purposes.

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1The code is available from the authors upon request
3 Results and comments

Considering SM processes $e^+e^- \to b\bar{b} + f\bar{f}$ at LEPII, it is well known that about 20% of the Higgs signal belongs to the final state $\nu\bar{\nu}b\bar{b}$, imbedded in a large background. Using the calculational procedure and the parameters presented in the previous section, we now proceed to investigate the characteristic properties of the HB-mediated (Fig. 1) and the coherent background (Fig. 2) contributions to $e^+e^- \to \nu\bar{\nu}b\bar{b}$ (sum over neutrino species), in order to provide useful information for Higgs searches at LEPII.

To gain an orientation we start with Fig. 3, which shows, for $m_H = 90$ GeV and over c.m. energies $150 < \sqrt{s}$ (GeV) < 240, the total expected number of $e^+e^- \to \nu\bar{\nu}b\bar{b}$ events for integrated luminosity 1000 pb$^{-1}$, together with the expected signal-only events from only the HB-mediated diagrams of Fig. 1 and the expected background events from the diagrams of Fig. 2. The expected signal events when one includes also the interference term between diagrams 1 and 2 are also plotted, but are indistinguishable (differences below 0.2 %) from the signal-only curve. In the same figure the two lower curves provide the signal-to-background ratio and the signal-to-Higgsstrahlung ratio, the second indicating the relative importance of the WW fusion contribution to the signal. It is interesting to note that the last contribution provides an important extension of the HB detectability limit at LEPII. We see from the figure that for $\sqrt{s}$ clearly above the HZ threshold the WW fusion contributes about 10% of the signal, but for $\sqrt{s}$ around and below $m_H + m_Z$ its relative importance grows rapidly with decreasing c.m energy. We will examine further down a realistic numerical example.

Fig. 4 shows the Higgs signal for $m_H = 80, 90, 100$ and 110 GeV and the background cross sections at LEPII energies, before and after taking into account ISR corrections. These corrections depend strongly on $\sqrt{s}$, ranging from 5% to 25% and being largest (as expected) at those energies where the cross sections increases most rapidly with $\sqrt{s}$.

We next examine the relative importance of various background contributions and their energy dependence, paying particular attention to $E_{CM} = 175, 192$ and 205 GeV, which are the nominal energies studied in [33]. Fig. 5 shows the separate contributions due to the background processes of Figures 2a, 2b, 2c and 2d, without and with the ISR corrections respectively. The $ZZ/\gamma$ cascade contribution (Fig. 2a) dominates. It grows rapidly at the $ZZ$ production threshold, while the other contributions increase more gradually with $\sqrt{s}$. At $\sqrt{s} = 175, 192$ and 205 GeV the total coherent background cross section with ISR corrections is about 0.006 pb, 0.045 pb and 0.065 pb respectively. It is clear that the diagrams of Fig. 2d can be practically ignored at LEP II energies, and those of Fig. 2b contribute too weakly to provide some sensitivity on the exact value of $m_t$. An interesting effect is a very strong destructive interference between the WW fusion and W exchange diagrams of the background, as clearly demonstrated on the same figure. This is a reflection of the unitarity (through gauge symmetry) of the SM, which prevents these background components from violating asymptotically the log($s$) bound.

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2The smallness of this interference term was checked to hold for all Higgs masses including $m_H = m_Z$
To decide about acceptance cuts, it is important to know the more accessible differential cross sections of the background and the signal contributions. As such we studied the differential cross sections versus the invisible invariant mass $m_{\nu\bar{\nu}}$, the visible invariant mass $m_{bb}$, the $\cos\theta_{bb}$ of the total visible 3-momentum $\vec{p}_{bb}$ and the $\cos\theta_b$ of the $b$-quark 3-momentum in the collision c.m. system. All differential cross sections have been calculated with ISR corrections. The above differential cross sections for the signal ($m_H=90$ GeV) and the total background at $\sqrt{s} = 192$ GeV are presented on figure 6. The next figure 7 shows how the background changes when decreasing $\sqrt{s}$ from 192 to 175 GeV. And Fig. 8 shows how the background differential cross sections split into approximatively gauge independent components $ZZ/\gamma$, WW fusion +W exchange and Z radiation.

The gross features of the background can be understood from the structure of the $ZZ/\gamma$ cascade diagrams, which dominates it. The sharp peak around $m_{\nu\bar{\nu}} \approx m_Z$ is evidently due to all $\nu\bar{\nu}$ pairs being produced through an intermediate $Z$. The two peaks at $m_{bb} \geq 0$ and $m_{bb} \approx m_Z$, are due to the $\gamma \to b\bar{b}$ and $Z \to b\bar{b}$ cascades in the diagrams of Fig. 2a. The peak at $m_{bb} \approx m_Z$ dominates as the energy increases and we pass the ZZ production threshold. One should also notice that the $ZZ/\gamma$ cascade and the $W$ exchange contributions peak along the beam axis. This is due to the $t$- and $u$-channel exchange of a light fermion in diagrams of Fig. 2a and 2c, leading in the differential cross section to factors of the form $(\sin^2\theta_{bb} + (\text{mass})^2/s)^{-1}$, where $(\text{mass})^2/s$ is very small for $Z\gamma$ cascade and small for $ZZ$ cascade (except for $\sqrt{s}$ just above the ZZ production threshold): for a discussion see e.g. [30].

The peak of the signal around $m_{\nu\bar{\nu}} \approx m_Z$ is due to the dominance of the Higgsstrahlung diagrams above threshold, while the WW fusion produces a broad $m_{\nu\bar{\nu}}$ spectrum. The narrow width of the HB is responsible for the spectacular peak at $m_{bb}$. One should also notice that, in contrast to the background, the signal angular distributions are remarkably isotropic. For the Higgsstrahlung contribution this is in agreement with the analytic result [31] that the differential cross section of $e^+e^- \to HZ$ has angular distribution (in the c.m. system)

$$1 + \frac{Q^2\sin^2\theta}{2m_Z^2}$$

where $Q = \lambda^{1/2}(s,m_Z^2,m_H^2)/2\sqrt{s}$ is the 3-momentum of $Z$ or $H$: for $\sqrt{s}$ just above $m_H + m_Z$, $Q^2/2m_Z^2$ is very small and the distribution is flat, but this will change sufficiently above the $HZ$ production threshold. The angular flatness of the $WW$ fusion signal is in accordance with the expectation of an isotropic distribution of the 3-momenta of the Higgs bosons produced by $WW$ fusion at energies not far beyond the WW threshold. It is also clear that there should be considerable constructive interference between the two signal components for $\sqrt{s}$ around and below the $HZ$ threshold. Indeed an examination shows that at these c.m. energies about half the increase of the signal cross section associated with the presence of WW fusion comes from the interference with the Higgsstrahlung, while at higher $\sqrt{s}$ this interference dies out. Similar conclusions have been reached in the very recent study of [32].

Knowing the main properties of the signal and the coherent background, one should next proceed to cope with the crucial problem of the very large incoherent background, whose main
components, as discussed in the introduction, are $B_{QC}^\gamma$, $B_{WW}$, $B_W$ and $B_{\ell\ell}$ ($\ell = e$ and $\mu$). This problem is clearly beyond the scope of this work: it has been studied by the LEP experimental groups as reported in [33]. We shall simply adopt the following event selection criteria, which were prescribed by the LEP200 event generator group, and study how they affect the signal and the coherent background:

(i) $50 \text{ GeV} < m_{bb}$

(ii) $(m_Z - 25 \text{ GeV}) \leq m_{\nu\bar{\nu}} \leq (m_Z + 25 \text{ GeV})$

Experimental suppression of all background to the HB signal in $e^+ e^- \rightarrow bb\bar{b}$ missing products requires of course several further event acceptance criteria, such as allowing events with only two jets, both $b$-tagged and of constrained morphology (acoplanarity and acollinearity cuts), forbidding events with too small number of charged tracks or with large missing energy in the beam pipe, etc [33].

The total signal and background cross sections after cuts (i) and (ii), for various HB masses and without or with ISR corrections, are shown on Figs. [10]. Fig 9 shows how these cuts affect the background components. We see that the coherent background is reduced greatly (by about 65%) at $\sqrt{s} = 175\text{GeV}$, but the reduction is much smaller (10-15%) at $\sqrt{s} = 192$ and $205 \text{ GeV}$. This can be attributed mostly to cut (i) which suppressed the photon mediated $bb\bar{b}$ production, whose relative importance in the total background cross-section decreases with increasing $\sqrt{s}$. As regards the signal, it is mostly influenced by cut (ii), which reduces the contribution induced by the WW fusion and characterised by a broad $m_{\nu\bar{\nu}}$ spectrum. So cuts (i)+(ii) lower the signal by only 5-10% for $\sqrt{s} \geq m_H + m_Z$, but the reduction grows rapidly as $\sqrt{s}$ decreases below the HZ threshold.

It is interesting to estimate carefully, including the effect of our cuts, how the HB detectability limit at LEPII is extended by the WW fusion process. To this end, we show on Fig 11 the $\sqrt{s}$ dependence of the (ISR corrected) signal cross section before and after cuts, for $m_H = 100$ GeV, and the corresponding total signal/Higgsstrahlung ratio. We see that for $\sqrt{s}$ clearly above $m_H + m_Z$, the WW fusion contribution (about 10% of the signal) is reduced by about 30-40% by our cuts. In the region 180-200 GeV, where the WW fusion relative importance increases rapidly with decreasing $\sqrt{s}$, our cuts reduce by only about 20% the signal-to-Higgsstrahlung ratio. E.g. at $\sqrt{s}=192 \text{ GeV}$ about 50% of the events are due to WW fusion, this reducing to about 40% after the cuts. The final number of expected events for an integrated luminosity (e.g. by summing 4 experiments over 2 years) of 1000 pb$^{-1}$ is 8.

Finally Fig. [12] presents information which allows an immediate estimate of the reduction of the number of background and signal events due to the limited acceptance of micro-vertex detectors used to tag the $b$ jets. It is clear that signal and coherent background are about equally lowered by this effect. Micro-vertex detectors extending to $\cos \theta=0.9$ permit the loss of only about 10-15% of the events when $b$-tagging of both jets is required, while one has a negligible loss (around 2%) if contented with only one $b$-tagged jet.

In order to allow comparisons with other existing Higgs codes, we followed the LEP200 Higgs event generator prescriptions concerning the input values and obtained the cross-sections.
requested for comparison purposes. We present them in the tables below. As can be seen in the agreement is satisfactory, differences are of the order of 1%.

\[
\begin{array}{|c|c|c|c|c|}
\hline
m_H \text{ (GeV)} & 65 & 90 & 115 & \infty \\
\hline
175 \text{ GeV } (\nu_\mu \bar{\nu}_\mu) & 63.447(0.289) & 2.316(0.011) & 1.257(0.006) & 1.257(0.006) \\
192 \text{ GeV } (\nu_\mu \bar{\nu}_\mu) & 72.01(0.317) & 46.840(0.093) & 19.575(0.087) & 19.453(0.087) \\
205 \text{ GeV } (\nu_\mu \bar{\nu}_\mu) & 67.704(0.301) & 53.569(0.240) & 25.531(0.114) & 23.646(0.105) \\
175 \text{ GeV } (\nu_e \bar{\nu}_e) & 70.103(0.302) & 4.996(0.023) & 1.051(0.006) & 1.051(0.006) \\
192 \text{ GeV } (\nu_e \bar{\nu}_e) & 79.158(0.348) & 52.744(0.140) & 20.605(0.102) & 19.678(0.098) \\
205 \text{ GeV } (\nu_e \bar{\nu}_e) & 76.165(0.359) & 61.429(0.280) & 31.048(0.146) & 25.737(0.122) \\
\hline
\end{array}
\]

Table 1: The process $e^+e^- \rightarrow \nu\bar{\nu}b\bar{b}$ at Born level, following the LEP200 generator Higgs group input values and prescriptions. Cross sections in fb. The numerical integration errors are included in parentheses.

\[
\begin{array}{|c|c|c|c|c|}
\hline
m_H \text{ (GeV)} & 65 & 90 & 115 & \infty \\
\hline
175 \text{ GeV } (\nu_\mu \bar{\nu}_\mu) & 54.186(0.283) & 1.660(0.008) & 0.917(0.006) & 0.917(0.006) \\
192 \text{ GeV } (\nu_\mu \bar{\nu}_\mu) & 64.405(0.246) & 35.367(0.109) & 14.68(0.087) & 14.60(0.087) \\
205 \text{ GeV } (\nu_\mu \bar{\nu}_\mu) & 63.826(0.222) & 45.736(0.171) & 21.275(0.105) & 20.058(0.105) \\
175 \text{ GeV } (\nu_e \bar{\nu}_e) & 60.806(0.294) & 3.663(0.018) & 0.775(0.006) & 0.775(0.006) \\
192 \text{ GeV } (\nu_e \bar{\nu}_e) & 71.151(0.268) & 40.212(0.164) & 15.222(0.098) & 14.61(0.098) \\
205 \text{ GeV } (\nu_e \bar{\nu}_e) & 71.794(0.266) & 52.543(0.199) & 25.428(0.124) & 21.76(0.122) \\
\hline
\end{array}
\]

Table 2: The process $e^+e^- \rightarrow \nu\bar{\nu}b\bar{b}$ including ISR corrections, following the LEP200 generator Higgs group input values and prescriptions. Cross sections in fb. The numerical integration errors are included in parentheses.
4 Summary and conclusions

In this work we studied the SM process $e^+e^- \rightarrow \nu\bar{\nu}b\bar{b}$ (sum over neutrino species) at c.m. energies $\sqrt{s} = 150-240$ GeV and for Higgs boson masses $m_H = 80-120$ GeV. The results were obtained from all tree level diagrams and included approximatively the most important radiative corrections, i.e. the QED ISR corrections, the QCD and EW corrections to the Higgs width and a QCD correction to the $Hb\bar{b}$ coupling. The matrix elements were calculated by the ‘spinor bracket’ method without neglecting masses, and we tried to make the method more accessible by describing it in detail and motivating its spinor basis in a novel way. The phase space integrals were calculated by an importance sampling Monte Carlo numerical integrator.

The signal is composed of the Higgsstrahlung and the WW fusion amplitudes. The first dominates for $\sqrt{s}$ above the HZ threshold, but the second becomes increasingly important as $\sqrt{s}$ decreases around and below that threshold, extending probably (depending on luminosity) to $m_H \geq \sqrt{s} - m_Z$ the Higgs detectability limit at LEPII. The two amplitudes behave similarly over phase space, providing in particular isotropic angular distributions of visible momenta, but differ in their $m_{\nu\bar{\nu}}$ spectrum at higher $\sqrt{s}$: the one of the WW fusion is always broad, while that of the Higgsstrahlung peaks strongly around $m_{\nu\bar{\nu}} = m_Z$ as soon as $\sqrt{s} \geq m_H + m_Z$. In fact the two amplitudes interfere semi-constructively for $\sqrt{s}$ below and around the HZ threshold, but hardly so well above it. About 50% of the signal cross section around that threshold is associated with the WW fusion, this becoming about 10% at higher energies.

The background consists of 4 classes of amplitudes, describing ZZ/$\gamma$ cascade, WW fusion, W exchange and Z radiation processes. The ZZ and Z$\gamma$ cascades dominate, the first leading beyond the ZZ threshold. Between the WW fusion and W exchange components there is a strong destructive interference and their coherent contribution is more than an order of magnitude smaller than the ZZ/$\gamma$ one, while the Z radiation contribution is practically negligible. The overall behaviour over phase space is characterised by a peak around $m_{\nu\bar{\nu}} = m_Z$, two peaks in $m_{b\bar{b}}$ due to the $\gamma \rightarrow b\bar{b}$ and $Z \rightarrow b\bar{b}$ cascades, and a clear angular peak along the beam axis. The total background cross section (before cuts) lies between the $m_H = 90$ GeV and $m_H = 100$ GeV signal cross sections throughout LEPII energies. The signal-background interference is negligible, even around $m_H = m_Z$.

Applying a set of minimal cuts prescribed by the LEPII groups, $m_{b\bar{b}} \geq 50$ GeV and $m_Z - 25$ GeV $\leq m_{\nu\bar{\nu}} \leq m_Z + 25$ GeV, has a very different effect for $\sqrt{s}$ below and above the HZ threshold. At lower energies both the signal and the background cross sections are greatly reduced, although the first mostly due to the $m_{\nu\bar{\nu}}$ cut while the second to the $m_{b\bar{b}}$ cut. At higher energies both the signal and the background are reduced weakly (\(\sim 10\%\)), the signal-background ratio improving only slightly. The effect of an angular acceptance of a b-tagging micro-vertex on the signal and the background were found roughly equal, rejecting about 10-15% of events for present type extended microvertices. The code exists also in the form of a MC generator.
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Figure 1: Signal Feynman diagrams
Figure 2: Background Feynman diagrams

Figure 3: Number of expected $\nu\bar{\nu}b\bar{b}$ events from the background, the Higgs signal ($m_H=90$ GeV) and their sum, as a function of $E_{CM}$ for 1000 pb$^{-1}$. The signal+interference is also plotted, but is indistinguishable from the signal-only curve. The ratios of the signal-to-background and the total signal-to-Higgstrahlung contribution are also shown. QED ISR corrections are included.

Figure 4: Higgs signal $\nu\bar{\nu}b\bar{b}$ cross section as a function of $E_{CM}$ for different HB masses and the background, without and with ISR corrections.

Figure 5: Total background $\nu\bar{\nu}b\bar{b}$ cross section and subprocess contributions as functions of $E_{CM}$ without and with ISR corrections.

Figure 6: Differential cross sections for the Higgs signal with $m_H$ 90 GeV (solid line) and for the background (dashed line) at $E_{CM}$=192 GeV. (ISR corrections included).

Figure 7: Differential cross sections for the background at 192 GeV (solid line) and 175 GeV (dashed line) (ISR corrections included).

Figure 8: Differential cross sections, at $E_{CM}$=192 GeV for the total background (solid line) and the subprocesses a) ZZ/$\gamma$ production (dashed line) b) W fusion and exchange (dotted dashed line) and c) Z radiation (dotted line) (ISR corrections included).
Figure 9: Total background $\bar{\nu}\nu\bar{b}b$ cross section and subprocess contributions as functions of $E_{CM}$, without and with ISR corrections, after cuts.

Figure 10: Higgs signal $\bar{\nu}\nu\bar{b}b$ cross section as a function of $E_{CM}$ for different Higgs masses and the background, without and with ISR corrections, after cuts.

Figure 11: The number of expected Higgs $\nu\bar{\nu}\bar{b}b$ events for 1000 pb$^{-1}$ as a function of $E_{CM}$ for $m_H = 100$ GeV, before and after cuts. The ratios of the signal-to-background and the total signal-to-Higgstrahlung contribution are also shown (ISR corrections are included).

Figure 12: Percentage of events for Higgs (dashed and dotted dashed lines) and background (solid and dotted) outside the fiducial acceptance of a microvertex extending to $\cos\theta$ when demanding a) at least one tagged jet or b) both jets tagged.