Neutrino Mass and Missing Momentum Higgs Boson Signals

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In the simplest scheme for neutrino masses invoking a triplet of Higgs scalars there are two CP-even neutral Higgs bosons H_i (i=1,2) and one massive pseudoscalar A. For some choices of parameters, the lightest H_1 may be lighter than the Standard Model Higgs boson. If the smallness of neutrino mass is due to the small value of the triplet expectation value, as expected in a seesaw scheme, the Higgs bosons may decay dominantly to the invisible neutrino channel. We derive limits on Higgs masses and couplings that follow from LEP I precision measurements of the invisible Z width.

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Neutrino mass constitutes one of the deepest open issues in the Standard Model (SM) of particle physics, which now finds some observational support [1]. Neutrino masses in the few eV range may also be crucial in explaining the large scale structure of the universe. In many $SU(2)\otimes U(1)$ extensions of the SM neutrino masses are generated from the spontaneous violation of a global lepton number symmetry leading to the existence of a physical Goldstone boson - the majoron [2]. In such models the majoron acts as the tracer of the neutrino mass generation mechanism and may have, depending on the details of the model, many interesting phenomenological implications [3]. If the breaking of lepton number occurs at the weak scale the lightest Higgs boson can decay dominantly into the weakly interacting majorons [4]. Since these escape detection, this decay is called invisible and has as signature missing momentum in the reaction. Here we consider a more direct way in which neutrino mass physics can show up as a missing momentum Higgs boson signature. As our illustrative model we consider the simplest triplet model for generating neutrino masses as first proposed, in the majoron-less form [5]. The model contains a complex $SU(2)$ triplet of scalar bosons Δ, in addition to the standard Higgs doublet ϕ

$$\phi = \begin{bmatrix} \phi^0 \\ \phi^- \end{bmatrix}, \; \Delta = \begin{bmatrix} \Delta^0 & \Delta^+/\sqrt{2} \\ \Delta^+/\sqrt{2} & \Delta^{++} \end{bmatrix}, \qquad (1)$$

where we have used the 2×2 matrix notation for the Higgs triplet. Apart from obvious terms in ϕ and Δ, the scalar potential contains two mixed quartic terms and one tri-linear term given as

$$\lambda_3 \phi^\dagger\phi \, \mathrm{tr}\,(\Delta^\dagger\Delta) + \lambda_5(\phi^\dagger\Delta^\dagger\Delta\phi) - \frac{\kappa}{\sqrt{2}}(\phi^T\Delta\phi + \mathrm{h.c.}) \qquad (2)$$

where the λ's are dimensional-less couplings and κ has mass dimension and should lie at the weak scale or less. Note that the smallness κ is natural according to 't Hooft's criterion, since the symmetry of the model increases as $\kappa \to 0$. We assume that the fields ϕ and Δ acquire nonzero real vacuum expectation values (VEVS) v_2 and v_3, respectively. It is easy to verify explicitly that, for many choices of its parameters, the potential has indeed minima for nonzero values of v_2 and v_3. According to this, we shift the fields in the following way

$$\phi^0 = \frac{v_2}{\sqrt{2}} + \frac{R_2 + iI_2}{\sqrt{2}}$$
$$\Delta^0 = \frac{v_3}{\sqrt{2}} + \frac{R_3 + iI_3}{\sqrt{2}} \qquad (3)$$

Notice that the existence of a cubic term in the scalar potential breaks explicitly the lepton number symmetry, avoiding the Gelmini–Roncadelli triplet majoron [6,7] - now ruled out experimentally by the measurement of the invisible Z width at LEP I [8].

As a result of the Yukawa coupling of the Higgs triplet to the leptons $\ell\Delta\ell$ a Majorana neutrino mass matrix will be generated, as follows

$$m_\nu = h_\nu v_3 \qquad (4)$$

In order to comply with cosmological limits the magnitude of the Yukawa couplings and/or v_3 is substantially restricted. We are interested in the limit of small values of the triplet VEV v_3. In this limit the smallness of neutrino mass in eq. (4) can be ascribed to the smallness of v_3 without having to invoke tiny Yukawa couplings. This is exactly what naturally happens in a seesaw scheme [9].

Taking into account the fact that this model contains one doubly-charged and one singly-charged scalar boson, in addition to the one charged unphysical $SU(2)\otimes U(1)$ Goldstone mode (longitudinal $W^\pm$), it follows that the neutral Higgs sector of this model is composed by four real fields. Due to CP invariance they split into two unmixed sectors. For example, for $v_3 \ll \kappa$ and $v_3 \ll v_2$ the CP–even Higgs mass squared matrix may be written as

$$\boldsymbol{M}_R^2 = \begin{bmatrix} m_H^2 & \left(\gamma m_H^2 - 2m_A^2\right)\frac{v_3}{v_2} \\ \left(\gamma m_H^2 - 2m_A^2\right)\frac{v_3}{v_2} & m_A^2 + \mathcal{O}(v_3^2) \end{bmatrix} \qquad (5)$$

where γ is a ratio of λ's, m_H^2 is the SM Higgs mass, and m_A^2 is the physical pseudoscalar boson mass given by

$$m_A^2 = \tfrac{1}{2}\kappa \frac{v_2^2 + 4v_3^2}{v_3} \qquad (6)$$

This mass results from diagonalizing the 2×2 CP–odd Higgs mass squared matrix via a rotation matrix O_I defined by a small angle

$$\sin\beta = \frac{2v_3}{\sqrt{(v_2^2 + 2v_3^2)}} \qquad (7)$$

and obeying $O_I \boldsymbol{M}_I^2 O_I^T = \mathrm{diag}(0, m_A^2)$, so that the first state is the unphysical Goldstone boson. On the other hand $O_R \boldsymbol{M}_R^2 O_R^T = \mathrm{diag}(m_{H_1}^2, m_{H_2}^2)$ where we take, by definition, $m_{H_1} \leq m_{H_2}$, so that the first state corresponds to the lightest CP–even Higgs boson. The corresponding matrix O_R is given in terms of an angle α. The parameter $\sin\alpha$ determining the projection of H_1 along the triplet can be large when $m_A < m_H$, as we will see below.

Note that none of the CP-even or CP-odd Higgs boson masses obtained in our model lie at the scale v_3. Consequently, even though the model suffers from the usual hierarchy problem, this may be avoided by supersymmetrization, as in the SM. This would estabilize all Higgs masses at the weak scale.

The W and Z masses come from the kinetic part of the scalar Lagrangian

$$\mathcal{L}_0 = (\mathcal{D}_\mu \phi)^\dagger \mathcal{D}^\mu \phi + \mathrm{tr}\,[(\mathcal{D}_\mu \Delta)^\dagger \mathcal{D}^\mu \Delta] \qquad (8)$$

where the covariant derivative is defined by

$$\mathcal{D} = \partial^\mu + ig\mathbf{T}\cdot\mathbf{W}^\mu + \frac{i}{2}g' Y V^\mu \qquad (9)$$

where g and g' are the $SU(2)$ and $U(1)$ gauge couplings respectively. The generators act on the scalars fields as

$$\begin{aligned} \mathbf{T}\phi &= \tfrac{1}{2}\vec{\tau}\phi, \qquad & \mathbf{T}\Delta &= -\frac{1}{2}\vec{\tau}\Delta - \frac{1}{2}\Delta\vec{\tau}^* \\ Y\phi &= -1\phi, \qquad & Y\Delta &= 2\Delta\,, \end{aligned} \qquad (10)$$

With these definitions we have $T_3\phi^0 = \frac{1}{2}\phi^0$ and $T_3\Delta^0 = -1\Delta^0$. The W mass is given by [5]

$$m_W^2 = \frac{1}{4}g^2(v_2^2 + 2v_3^2) \qquad (11)$$

so that $\sqrt{v_2^2 + 2v_3^2} \simeq 246$ GeV. From the measurement of the ρ parameter one has [10]

$$\rho = 1 + \frac{2v_3^2}{v_2^2 + 2v_3^2} = 1.001 \pm 0.002\,. \qquad (12)$$

which implies in practice that $v_3 \leq 9.5$GeV and v_2 almost fixed, leading to $\sin\beta \lesssim 10^{-2}$. This restriction on v_3 is automatically fulfilled for the values we are dealing with.

In order to have an idea of the expectations of the model for the various Higgs boson masses we diagonalize the exact mass matrices and impose the potential minimisation conditions, checking the positivity of the physical CP-even and CP-odd eigenvalues. In Fig. (1) we show the CP–odd Higgs boson mass in our model versus κ for different v_3 values. The allowed region lies above the curve corresponding to $v_3 \approx 9.5$ GeV.

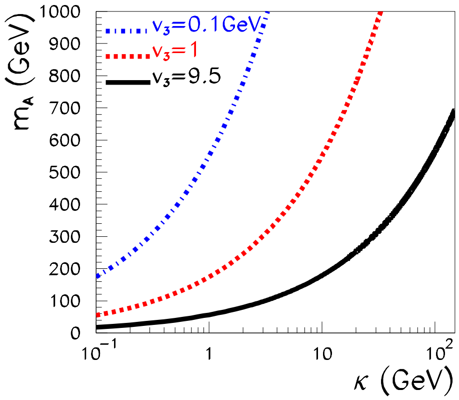

FIG. 1. Lightest CP–odd Higgs boson mass versus κ for different v_3 values.

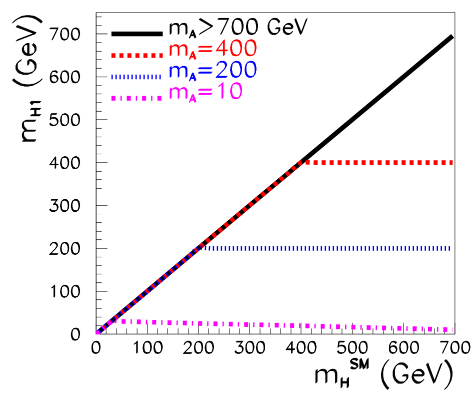

FIG. 2. Lightest CP–even Higgs mass vs the SM Higgs boson mass for different m_A values.

Similarly in Fig. (2) we show m_{H_1}, the mass of the lightest CP–even Higgs boson in our model, as a function of the SM Higgs mass m_H for different m_A values. For example, if we fix m_A at 200 GeV we find that H_1 lies along the solid diagonal line and is almost the same as the SM doublet Higgs boson up to $m_H \sim 200$ GeV. Past this value the H_1 mass lies on the horizontal line at $m_{H_1} =$ 200 GeV and H_1 is therefore mostly triplet. As we will see later, it will decay mostly to neutrinos instead of to b-quarks. Another feature worth noticing is that there can be a substantial mixing in the CP–even sector and, as a result, the H_1 mass may also be lower than the SM Higgs mass m_H, especially for small m_A values. Finally, we can also see that $m_{H_1} < m_A$, so that the decay $H_1 \to AA$

does not occur.

The mechanisms for Higgs boson production at e^+e^- colliders are the emission of a CP–even Higgs by a Z–boson, and the associated production of a CP–even Higgs and a CP–odd Higgs A. The relevant couplings are given by

$$\mathcal{L}_{H_a AZ} = \frac{g}{2c_w} Z^\mu \left[R_2 \overleftrightarrow{\partial^\mu} I_2 - 2\, R_3 \overleftrightarrow{\partial^\mu} I_3 \right]$$
$$= \frac{g}{2c_w} Z^\mu \left[\sin\beta O^R_{a2} - 2\cos\beta O^R_{a3} \right] H_a \overleftrightarrow{\partial^\mu} A. \qquad (13)$$

where $c_W \equiv \cos\theta_W$ and H_a is any of the two CP–even neutral Higgs bosons. The parameter defined by

$$\epsilon_A = \sin\beta\cos\alpha + 2\cos\beta\sin\alpha \approx 2\sin\alpha \qquad (14)$$

will determine the strength of the H_1AZ coupling. From eq. (8) we find that the H_aZZ couplings are

$$\mathcal{L}_{H_a ZZ} = \frac{g m_Z}{4c_W^2} Z^\mu Z_\mu \left[\cos\beta O^R_{a2} + 2\sin\beta O^R_{a3} \right] H_a, \quad (15)$$

and correspondingly we define

$$\epsilon_B = \cos\beta\cos\alpha - 2\sin\beta\sin\alpha \approx \cos\alpha \qquad (16)$$

as a measure of the strength of the H_1ZZ coupling.

Notice the factor of 2 in the expressions for the parameters ϵ_A and ϵ_B which determine the Bjorken and the associated production cross sections, respectively. It comes from the hypercharge of the triplet.

We now turn to the couplings relevant for the invisible decay of the lightest Higgs bosons. The pseudoscalar A may decay into stable neutrinos, via the triplet Yukawa coupling h_ν of eq. (4), or into $b\bar{b}$, via its projection onto the doublet. In order to evaluate the relative importance of the two branchs we need information on the neutrino mass. In the presence of the cubic lepton–number–breaking term κ there is no efficient means of reducing the relic neutrino number density, as a result of which neutrinos in this model must obey the limit from cosmology on stable neutrino masses

$$m_\nu \lesssim 92\,\Omega h^2 \mathrm{eV} \qquad (17)$$

where $\Omega h^2 \leq 1$ is a basic cosmological parameter [11]. Present determinations give h is $.65 \pm 0.1$, while the total Ω may still be as large as one, as suggested by inflationary models.

Without adding singlet scalar bosons there is, on the other hand, no way to introduce the majoron in a phenomenologically acceptable way, which does not conflict with LEP I data on the observed width of the Z into invisible channels. From eq. (4) we see that

$$\frac{\Gamma(A \to b\bar{b})}{\Gamma(A \to \nu\bar{\nu})} = \left(\frac{3h_b \sin\beta}{h_\nu \cos\beta} \right)^2 \approx \left(\frac{6m_b}{m_\nu} \right)^2 \frac{v_3^4}{v_2^4} \qquad (18)$$

where h_ν here denotes the Yukawa coupling of the most massive of the neutrinos with the triplet. One sees that for $m_\nu \sim 10$ eV and $v_3 \lesssim$ few MeV, the decay of A to neutrinos will be dominant.

An analogous calculation for the CP–even sector gives

$$\frac{\Gamma(H_1 \to b\bar{b})}{\Gamma(H_1 \to \nu\bar{\nu})} = \left(\frac{3h_b \cos\alpha}{h_\nu \sin\alpha} \right)^2 \lesssim \left(\frac{6m_b}{m_\nu} \right)^2 \frac{v_3^4}{v_2^4} \gamma^2 \quad (19)$$

Clearly for $\cos\alpha \to 0$ H_1 is mostly triplet and therefore decays mainly to neutrinos. This corresponds to the horizontal lines in Fig. 2. In the opposite situation $\cos\alpha \to 1$ H_1 is mostly doublet and decays mainly to $b\bar{b}$. The price we must pay in order to have $\cos\alpha \approx 0$ is to have again a small v_3 as seen from eq. (19). In the last step in eq. (19) we assumed $m_A \ll m_H$ in order to obtain a conservative upper bound on v_3 as a function of γ. Again, one sees that for $m_\nu \sim 10$ eV and $\gamma = 1$ we find that for v_3 values in the MeV range the lightest CP–even Higgs boson H_1 will also decay invisibly. For smaller γ values the upper bound on v_3 is relaxed correspondingly.

Note that in the simplest scheme presented above the smallness of v_3 is put in by hand. However, the simplest model may be regarded as an effective parametrization of a more complete left-right symmetric see-saw scheme [9] in which lepton number is a local symmetry violated spontaneously at a large scale v_R. The smallness of v_3 would account for the smallness of neutrino mass and would arise naturally from a minimization condition of the scalar boson potential leading to $v_3 \sim m_W^2/v_R$.

We now perform a model independent study of the limits that can be set based on Higgs boson production in e^+e^- colliders at the Z peak and its subsequent invisible decay. Consider the massive pseudoscalar A and the lightest CP–even scalar H_1. As we have seen H_1 and A may decay invisibly when v_3 is small. As seen above, for small γ the bound on v_3 for H_1 to decay invisibly may be somewhat relaxed. The basic parameters needed to describe the implications of the production of Higgs bosons at the Z peak in this model are the masses m_A and m_{H_1}, the coupling parameters ϵ_A and ϵ_B which determine the Bjorken and associated production cross sections and the product of the visible and invisible decay branching ratios.

The Bjorken process contribution to the invisible Z width $Z \to Z^* H_1$ is, for most of the parameter space, very small compared to that of the associated production. Thus, in order to get a conservative bound, we only consider the associated process, bearing in mind that the inclusion of the Bjorken contribution would only improve our results, i.e., would exclude a slightly wider region of parameter space.

Therefore we consider in what follows the limits that can be set on associated Higgs boson production at the Z peak, $e^+e^- \to Z \to H_1 A$ when both CP-even (H_1) as well as CP-odd Higgs bosons (A) decay invisibly. One can write the contribution to the invisible Z width as:

$$\Delta\Gamma_{inv} = \frac{{\epsilon_A}^2}{2}\lambda^{\frac{3}{2}}\left(1, \frac{m_{H_1}^2}{m_Z^2}, \frac{m_A^2}{m_Z^2}\right)\Gamma(Z \to \nu_e\bar{\nu}_e)B_{inv}A_{inv}$$

where B_{inv} and A_{inv} denote the invisible branching ratios of $BR(H_1 \to \nu\bar{\nu})$ and $BR(A \to \nu\bar{\nu})$, respectively, and λ is the usual Källen function $\lambda(a,b,c) = (a-b-c)^2 - 4bc$.

Taking into account the experimental error in the determination of the invisible Z width, given by $\sigma = 2$ MeV [8], we have determined 95 % CL bounds on ${\epsilon_A}^2$ in the m_{H_1}-m_A plane for fixed values of the product $B_{inv}A_{inv}$. As it is well known, the integration of a Gaussian probability distribution from $-\infty$ to $+1.64\sigma$ gives 0.95, so that the 95 % CL exclusion region is defined by imposing $\Delta\Gamma_{inv} > 1.64\sigma$. In Fig. (3) we show these results for a fixed value of the product $B_{inv}A_{inv} = 1$. This corresponds to both scalar and pseudoscalar decaying totally to neutrinos. This choice is meant for definiteness. The constraints corresponding to any other value of $B_{inv}A_{inv}$ may be simply obtained by rescaling the results for our reference value given in Fig. (3). In the plot we have five curves labeled by a value of ϵ_A^2.

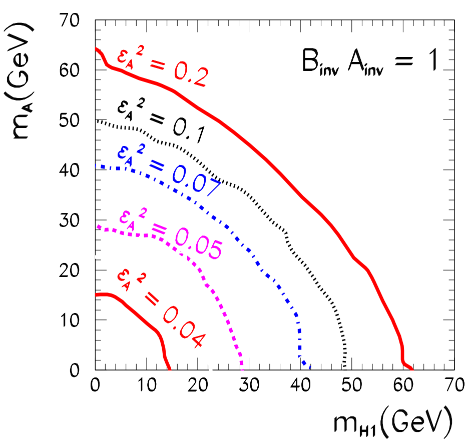

FIG. 3. 95 % CL bounds on ${\epsilon_A}^2$ in m_H-m_A plane when both H_1 and A decay to neutrinos.

No points below each of these curves are allowed with ϵ_A^2 larger than that value. We see from this plot that simply by using the neutrino counting at the Z peak one can already impose important constraints on the parameters of the model. For example, for H_1 and A masses around 20 GeV the upper bound on ${\epsilon_A}^2$ is a few times 10^{-2}. Going beyond this requires a dedicated analysis of the various event topologies that are possible in this model, for example, di-jet plus missing momentum. Fortunately the same topologies are also the ones present in other models where the invisible Higgs boson decay involves majorons, considered both for LEP I as well as LEP II data [12]. The results of that analysis can easily be adapted to the present model. In any case an updated analysis of the present LEP II data by the LEP colaborations themselves would be welcome.

In short, we have illustrated, with a very simple model, how LEP experiments may shed information on the Higgs boson sectors of models of neutrino mass. In contrast, all previously considered models with invisibly–decaying Higgs bosons invoked the existence of majorons and employed only physics at the weak scale. We have shown how the invisibly decaying Higgs boson signal can arise from the existence of a small scale (v_3 in our model) associated to neutrino masses in a model with explicit violation of lepton number. In a more complete see-saw left-right scheme, the smallness of v_3 would arise naturally as $v_3 \sim m_W^2/v_R$.

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