Higgs-strahlung in Abelian Extended
Supersymmetric Standard Model

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Abstract

We analyze the Higgs production via the Higgs-strahlung process $e^+e^- \rightarrow h_k l^+l^-$ in an Abelian Extended Supersymmetric SM. We work in the large Higgs trilinear coupling driven minimum of the potential, and find that the next-to-lightest Higgs cannot be produced by this process. Other Higgs scalars, namely the lightest and the heaviest, have cross sections comparable to that in the pure SM. It is found that the present model has observable differences with the other popular model, NMSSM, in the same type of minimum.

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1 Introduction

The main mechanism for Higgs production in $e^+e^-$ collisions at LEP2 energies is the Higgs-strahlung process $e^+e^- \rightarrow ZH$ in which $H$ is radiated from a virtual $Z$ line. In the literature, this process has already been analyzed in SM \cite{1}, MSSM \cite{2} and NMSSM \cite{3,4}. Hence, it would be convenient to discuss this collision process in other models, such as $Z'$ models, to confront the experimental results with all possible theoretical predictions. The $Z'$ models in which the SM gauge group is extended by an extra $U(1)$ follow directly from SUSY-GUT’s or String compactifications. In addition this sound theoretical foundation, they have proven to be phenomenologically viable as well. For instance, like NMSSM, the $Z'$ models can solve the MSSM $\mu$ problem when SUSY is broken around the weak scale. However, they differ from MSSM and NMSSM in that they are further constrained the phenomenologically required smallness of the $Z - Z'$ mixing angle. In \cite{5}, a general string-inspired model was analyzed and phenomenologically viable portions of the parameter space have been identified. In this work we shall be working in one possible minimum constructed in \cite{5}, that is, we shall consider that portion of the parameter space where the Higgs soft trilinear coupling is large compared to other soft masses in the Higgs sector.

In a recent work \cite{6} we have analyzed another important LEP2 process $e^+e^- \rightarrow (hA) \rightarrow b\bar{b}b\bar{b}$ in the Abelian Extended Supersymmetric Standard Model which was already introduced in \cite{5}. In \cite{6} we found that it was essentially the second neutral Higgs scalar $h_2$ contributing to the process in the limit of small $Z - Z'$ mixing angle. Now, using the model in \cite{6}, we shall analyze the Higgs-strahlung type process $e^+e^- \rightarrow h_k l^+ l^-$, where $l = (e, \mu)$ and $h_k$ ($k=1,2,3$) be the CP-even Higgs scalars of the model. As is well known from MSSM \cite{7} Higgs-strahlung and pair production processes are complementary to each other and we expect similar behaviour to occur in the present model as well.

In Sec. 2 we shall reproduce the formulae relevant to the aim of the work, at the tree level. In Sec.3 we shall first analyze the cross section numerically, and discuss the results. Next, we consider the contribution of the radiative corrections at the one-loop order, and show that they do not alter the tree-level results significantly.
2 Formalism

The coupling of CP-even Higgs scalars $h_k$ ($k=1,2,3$) to vector bosons $Z_i$ ($i=1,2$) can be expressed as

$$L = K_{kij} g_{\mu\nu} Z_i^\mu Z_j^\nu h_k,$$  \hspace{1cm} (1)

where the couplings $K_{ijk}$ are model dependent [7]. In giving the expressions for $K_{ijk}$ in the present model it is convenient to work with two dimensionless quantities, that is, Higgs Yukawa coupling $h_S$, and normalized $U(1)_Y'$ coupling constant $\rho_S$ defined by

$$\rho_S = \frac{g_{Y'} Q_S}{G},$$  \hspace{1cm} (2)

where $g_{Y'}$ and $Q_S$ are the $U(1)_Y'$ coupling constant and $U(1)_Y'$ charge of the SM singlet, respectively.

We shall first present the expressions for all the masses and couplings using the tree-level potential, deferring the discussion of effects of the radiative corrections to the next section. In the large trilinear coupling minimum [4], CP-even Higgs masses have, in the increasing order, the following expressions:

$$m_{h_1} \sim \frac{\sqrt{2} h_S M_Z}{G}, \quad m_{h_2} \sim \sqrt{1 + \frac{2 h_S^2}{G^2} M_Z}, \quad m_{h_3} \sim \sqrt{3 \rho_S + \frac{2 h_S^2}{G^2} M_Z},$$  \hspace{1cm} (3)

where $G = \sqrt{g_2^2 + 2 g_{Y'}^2}$, and $g_2$ and $g_{Y'}$ are $SU(2)$ and $U(1)_Y$ coupling constants. Moreover, neutral gauge bosons are almost diagonal with the masses $M_{Z_1} = M_Z$, and $M_{Z_2} \sim \sqrt{3} \rho_S M_Z$. It is clear that $\rho_S$ measures the $Z'$ mass in units of $\sqrt{3}$. Using the elements of the diagonalizing matrix $R$ defined in [4], one can easily obtain the coupling constants $K_{kij}$ in (1) as listed below:

$$K_{111} = \frac{G M_Z}{\sqrt{6}} (\cos^2 \alpha + 3 \rho_S^2 \sin^2 \alpha),$$

$$K_{211} = \frac{G M_Z}{2} \rho_S \sin 2\alpha,$$

$$K_{311} = -\frac{G M_Z}{\sqrt{12}} (\cos^2 \alpha - 3 \rho_S^2 \sin^2 \alpha).$$
\[ K_{122} = \frac{GMZ}{\sqrt{6}}(\sin^2 \alpha + 3\rho_S^2 \cos^2 \alpha) \]
\[ K_{222} = -\frac{GMZ}{2} \rho_S \sin 2\alpha \]  
(4)
\[ K_{322} = -\frac{GMZ}{\sqrt{12}}(\sin^2 \alpha - 3\rho_S^2 \cos^2 \alpha) \]
\[ K_{112} = -\frac{2GMZ}{\sqrt{6}} \rho_S^2 \sin 2\alpha \]
\[ K_{212} = -GMZ \rho_S \cos 2\alpha \]
\[ K_{312} = -\frac{4GMZ}{\sqrt{12}} \rho_S^2 \sin 2\alpha \]

where \( \alpha \) is the \( Z - Z' \) mixing angle, and is well below the phenomenological upper bound in the present minimum of the potential. The couplings \( K_{kij} \) listed in (4) are interesting enough to have a detailed discussion of them here. If \( \rho_S \sim O(1) \) (which should be the case in phenomenologically acceptable \( Z' \) models [8]), one can safely neglect \( \sin \alpha \) dependent terms all together. In this case the couplings \( K_{ijk} \) get simplified and possess the following properties:

- \( h_1 \) and \( h_3 \) do have only \( Z_i Z_i \) type couplings, that is, they have diagonal couplings in \( (Z_1, Z_2) \) space.

- Coupling costant of \( h_3 \) is always \( -1/\sqrt{2} \) times the coupling constant of \( h_1 \), which is an immediate consequence of the diagonalizing matrix \( R \) given in [8]. In this sense \( h_3 \) is a replica of \( h_1 \) with a larger mass.

- Coupling of \( h_2 \) requires always a transition from \( Z_1 \) to \( Z_2 \) and vice versa, that is, it does have only off-diagonal couplings in \( (Z_1, Z_2) \) space. \( Z_1 \), for example, yields \( Z_2 \) after radiating \( h_2 \); however, it yields \( Z_1 \) in the final state, after radiating \( h_{1,3} \). After discussing the coupling of Higgs scalars to fermions, the status of \( h_2 \) will be clearer.

The vector coupling \( v_f^{(i)} \) and axial coupling \( a_f^{(i)} \) of Higgs scalars to the fermion \( f \) and vector boson \( Z_i \) are given by
\[
\begin{pmatrix}
  v_f^{(1)} \\
  v_f^{(2)}
\end{pmatrix} =
\begin{pmatrix}
  \cos \alpha & \sin \alpha \\
  \sin \alpha & -\cos \alpha
\end{pmatrix}
\begin{pmatrix}
  v_f \\
  v_f'
\end{pmatrix} \tag{5}
\]
\[
\begin{pmatrix}
  a_f^{(1)} \\
  a_f^{(2)}
\end{pmatrix} =
\begin{pmatrix}
  \cos \alpha & \sin \alpha \\
  \sin \alpha & -\cos \alpha
\end{pmatrix}
\begin{pmatrix}
  a_f \\
  a_f'
\end{pmatrix} \tag{6}
\]
where

\[ v_f = \frac{g_2}{4 \cos \theta_W} (1 - 4 \sin^2 \theta_W) \]  \hspace{1cm} (7)

\[ a_f = \frac{g_2}{4 \cos \theta_W} \]  \hspace{1cm} (8)

\[ v'_f = \frac{g_{Y'}}{2} (Q_{fL} + Q_{fE}) \]  \hspace{1cm} (9)

\[ a'_f = \frac{g_{Y'}}{2} (Q_{fL} - Q_{fE}) \]  \hspace{1cm} (10)

where \( Q_{fL} \) and \( Q_{fE} \) are the \( U(1)_Y \) charges of lepton doublet \( L \) and lepton singlet \( E^c \), respectively.

**3 \quad e^+ e^- \rightarrow h_k l^+ l^- Cross Section**

We shall analyze the cross section for the \( e^+ e^- \rightarrow h_k l^+ l^- \) scattering, assuming that it proceeds via the following chain

\[ e^+ e^- \rightarrow Z_i^* \rightarrow h_k Z_j^* \rightarrow h_k l^+ l^- . \]  \hspace{1cm} (11)

The possibility that the vector bosons might come to mass-shell will be taken into account by including their total widths to their propagators. A straightforward calculation yields the following expression for the differential cross section

\[ \frac{d\sigma}{dx} = \frac{1}{36(4\pi)^3} \sqrt{(1 + r_k - x)^2 - 4r_k} \left( (1 - r_k)^2 - 2x(r_k - 5) + x^2 \right) \]

\[ \sum_{ijmn} K_{kij} K_{kmn} (a_e^a a_e^i + v_e^a v_e^i) (a_l^n a_l^j + v_l^n v_l^j) \]

\[ D_{Z_i}(s) D_{Z_a}(s) D_{Z_j}(x) D_{Z_m}(x) \]  \hspace{1cm} (12)

where \( x = m_{l^+l^-}^2 / s \) is the invariant dilepton mass normalized to \( s \), \( r_k = m_{h_k}^2 / s \), and

\[ D_{Z_i}(s) = \frac{1}{s - M_{Z_i}^2 + iM_{Z_i}\Gamma_{Z_i}} \]  \hspace{1cm} (13)
is the propagator of the vector boson $Z_i$ and $\Gamma_{Z_i}$ is its total width. In the cross section differential in (12), $x$ has the range $0 \leq x \leq (1 - \sqrt{r_k})^2$ as follows from the kinematics of the decay.

For each value of $k$, using the coupling constants $K_{kij}$ listed in (4), we can evaluate $\frac{d\sigma}{dx}$ from (12). In the limit of small mixing angle, equations (5) and (6) yield the following vector and axial couplings:

\begin{align*}
v_f^{(1)} &= v_f ; & a_f^{(1)} &= a_f \\
v_f^{(2)} &= v'_f ; & a_f^{(2)} &= a'_f.
\end{align*}

(14)

As is seen from (9) and (10), vector and axial couplings of $Z'$ boson depends on the unknown $U(1)_{Y'}$ charges of the lepton superfields, as expected. The family independent $Z'$ models are severely constrained by the Z-pole observables. However, if $U(1)_{Y'}$ charges are family- dependent, then $Z'$ models are less constrained by the LEP data. Indeed, such family dependent extra $U(1)$’s are realizable in string models [9, 5]. Alternatively, one can assume leptophobic $Z'$ [10] and neglect these $U(1)_{Y'}$ lepton charges. Based on these reasonings, we shall assume vanishing axial and vector leptonic couplings for $Z'$:

\begin{align*}
v_f^{(2)} &= 0 ; & a_f^{(2)} &= 0,
\end{align*}

(15)

which implies that $\frac{d\sigma_k}{dx}$ in (12) gets no contribution from $Z_2$ boson. Thus, the production of $h_1$ and $h_3$, with the diagonal couplings mentioned before, proceed by the contribution of the $Z_1$ boson only. The difference between $h_1$ and $h_3$ comes mainly from the dependence of their masses on $h_S$ and $\rho_S$ (See equation (3)). As opposed to $h_1$ and $h_3$, $h_2$ production is impossible as its production cross section identically vanishes due to its vanishing couplings to vector bosons in the present minimum of the potential.

That only $h_2$ contributes to pair production process $e^+e^- \to hA$ [6], and no contribution comes to Higgs-strahlung $e^+e^- \to h_kZ_i$ from $h_2$ is a special case of the complementarity between the two processes, as already noted in MSSM [2]. In the scattering process under concern, a non-zero $h_2$ signal would show up when the lepton superfields attain non-zero $U(1)_{Y'}$ charges and mixing angle is large enough. However, when these two conditions are satisfied one expects non-negligible shifts in the Z-pole observables in such a light $Z'$ model [11].

5
In analyzing $e^+e^- \rightarrow h_kl^+l^-$ scattering process we first integrate the differential cross section given in (12) over the entire range of $x$ to obtain the total cross section. Next we depict the dependence of the total cross section $\sigma$ for each Higgs on $h_S$ and $\rho_S$. In this way we can investigate the dependence of the cross section on $\mu$ parameter and $Z'$ mass. In the numerical analysis we take a lower bound of 60 $GeV$ for $m_{h_1}$ which implies a lower bound of $h_{s_{\text{min}}} \sim 0.345$ via (3).

Depicted in Fig. 1 is the $h_S$ dependence of the $h_1$ production cross section, $\sigma_{h_1}$ for $\sqrt{s} = 130, 136, 161, 172, 183, 205$ $GeV$ from left to right. The first five points have been probed at LEP2 and the last one is planned to be under search soon. In forming Fig. 1 we let $h_S$ vary from 0.345 (corresponding to $m_{h_1} \sim 60$ $GeV$) up to the kinematical threshold $\sqrt{G^2/(2M_Z)}$ (corresponding to $m_{h_1} \sim \sqrt{s}$). For each $\sqrt{s}$ value the cross section vanishes at some value of $h_S$ at which the Higgs production stops. As is seen from the figure, for $\sqrt{s} = 205$ $GeV$ the cross section remains around 10 $fb$ up to $h_S \sim 0.7$, so that $m_{h_1}$ values up to 130 $GeV$ can be probed at this CM energy depending on the total integrated luminosity of the collider under concern. Similar behaviour occurs for $\sqrt{s} = 161, 172, and 183$ $GeV$ for $h_S$ less than 0.4, 0.45 and 0.55, respectively.

Being much heavier than $h_1$, production of $h_3$ is very hard at these CM energies. To demonstrate the behaviour of the cross section we plot $\sigma_{h_3}$ as a function of $h_S$ for $\sqrt{s} = 183$ $GeV$ and various values of $\rho_S$. Using $r_3 \leq 1$, $M_{Z_3} \geq M_Z$, and $h_S \geq 0.345$, it is easy to find the ranges of $h_S$ and $\rho_S$: $0.58 \leq \rho_S \leq 1.17$ and $0.87 \geq h_S \geq 0.345$. As is seen from the figure, the cross section is much smaller than $\sigma_{h_1}$, and vanishes for certain $h_S$ values as in Fig. 1. For small $h_S$ the largest cross section occurs for $\rho_S = 0.58$ at $h_S = 0.345$. For higher values of $\rho_S$, cross section gets maximized at the higher end of the $h_S$ range. It is clear that with such a small cross section, it is necessary to have a large luminosity to obtain sufficient number of events differing hopefully from the predictions of the SM.

As is generally the case, unless one resorts to some sophisticated numerical techniques (like the PYHTIA and JETSET packages), the theoretical results obtained in this way can hardly be confronted with the experiment directly, because of the various cuts applied in event selection process to suppress the unwanted background. Indeed, as one of the latest LEP2 releases [12] shows, the Bjorken process under concern has a large SM background and Higgs search in MSSM and 2HDM for $\sqrt{s} = 130 - 172$ $GeV$ has a negative result.
However, the results of the present work may serve as a guide for the design of near-future colliders like NLC.

In the above we have mainly focused on unpolarized electron-positron beams in which cross section differential \( d\sigma/d\Omega \) involves only one angular variable: the angle between Higgs particle and the electron beam. When the initial beams are transversely polarized, however, \( d\sigma/d\Omega \) picks up an explicit dependence on the Higgs boson azimuthal angle. Therefore, in this case, in addition to the cross section itself, one can work out the azimuthal asymmetry of the process, which provides further information on the appropriate portion of the parameter space of the model under concern. However, let us note that the total cross section for transversely polarized beams is the same as the one for unpolarized beams. If the electron-positron beams are longitudinally polarized, angular dependence of \( d\sigma/d\Omega \) is the same as in the unpolarized case. However, the difference between the two beam polarizations enables one to define spin asymmetry of the process which can be important in constraining the parameter space of the model. Unlike the case of the transverse polarization of the beams, here the total cross section differs from that of the unpolarized case. The SM-counterpart of the process under concern was analyzed by [13] with a detailed discussion of the polarization effects and it seems unnecessary to have a detailed discussion of the polarization effects here because results of [13] applies equally to the \( Z' \) model under concern.

Until now our analysis has been based entirely on the tree-level potential in the limit of large trilinear Higgs coupling, which has already been shown to follow from string-inspired models in [8] through the solution of RGE’s. However, it is known that the results obtained using the tree-level potential are very sensitive to the choice of the renormalization scale \( Q \) of the RGE’s, and hence one should take into account the effects of the loop corrections on the scalar potential to get more reliable results. This is what we will do next, up to one-loop order, in the framework of the effective potential formalism[14].

In the framework of the effective potential formalism, the contribution of the radiative corrections to the scalar potential is given by

\[
\Delta V = \frac{1}{64\pi^2} \text{Str} \mathcal{M}^4 \ln \left[ \frac{\mathcal{M}^2}{Q^2} - C \right],
\]

where \( \text{Str} \) means the usual supertrace, \( \mathcal{M}^2 \) is the field dependent mass-
squared matrix of the fields, and $C$ is a constant that depends on the renormalization scheme $[14]$. The dominant contribution to the effective potential comes from the top quark and stops, due to relatively large value of the top Yukawa coupling. In the limit of degenerate stops, the most significant one-loop corrections arise for Higgs trilinear coupling $A_s$, $H_2$ quartic coupling $\lambda_2$, and $H_2$ soft mass-squared $m^2_2$, namely

$$
\hat{A}_s = A_s + \beta_{h_t} S_{t\tilde{t}} A_t \\
\hat{m}^2_2 = m^2_2 + \beta_{h_t} [(A^2 + A^2_t) S_{t\tilde{t}} - A^2] \\
\hat{\lambda}_2 = \lambda_2 + \beta_{h_t} S_{t\tilde{t}} h_t^2
$$

(17)

where $\beta_{h_t} = \frac{3}{(4\pi)^2} h_t^2$, $S_{t\tilde{t}} = \ln \frac{m_{t1} m_{t2}}{m_t^2}$ is the top-stop splitting function, and $A^2 = m^2_Q + m^2_U$, $m^2_Q$ and $m^2_U$ being the soft mass-squareds of squark doublet $\tilde{Q}$ and up-type $SU(2)$ squark singlet $U^c$, respectively. As was already discussed in [5], one can realize the large Higgs trilinear coupling minimum at low energies by integrating RGE’s with non-universal boundary conditions from the string scale down to the weak scale. The resulting parameter space satisfies $m^2_Q \sim m^2_U \sim (h_t A_t)^2 \sim (h_s A_s)^2$ that are much larger than the Higgs soft mass-squareds. For this parameter set, it is easily seen that $\hat{m}^2_{t1,2} \sim m^2_Q + m^2_t$, proving the consistency of assuming degenerate stops. Consequently $S_{t\tilde{t}} \sim \ln(1 + m^2_Q/m^2_t)$, and it remains close to unity as the squark soft masses are around the weak scale.

In the light of these results, one can now discuss the stability of the large trilinear coupling minimum under the contributions of the radiative corrections in (16). From the first equality in (16) one observes that $A_s$ increases in proportion with $A_t$. Similarly, from the second equation in (16) one infers that $m^2_2$ increases with $A^2_t$. Letting $m^2 = m^2_Q + m^2 + m^2_2$ be the sum of the soft mass-squareds of the Higgs fields, and $\lambda$ be the sum of the quartic coefficients in the potential, one can show that the large trilinear coupling minimum occurs after $A_s$ exceeds the critical point

$$
A^\text{crit}_s = \sqrt{\frac{8 \lambda m^2}{9 h^2_s}}.
$$

(18)

This formula is valid when $m^2 > 0$, which is easy to realize at low energies if non-universality is permitted at the string level [5]. Actually, such a condition is necessary for the need to an $A_s$ dominated minimum to arise, as otherwise
one can create an appropriate low energy minimum by using merely the Higgs soft mass-squareds. At the tree-level \( \lambda = 3h_s^2 \), and at the loop-level it increases by \( \beta_{ht} S_t h_t^2 \) due to the contribution of \( \lambda_2 \) in (16). Similarly, due to the contribution of \( \tilde{m}_t^2 \) in (16), \( m^2 \) increases by \( \beta_{ht} A_t^2 S_t \tilde{t} \). Inserting these loop contributions to (17), and using the tree-level relation \( A_s \sim A_t \) required by the solution of the RGE’s, one gets the following expression for the radiatively corrected \( A_s^{\text{crit}} \):

\[
A_s^{\text{crit}} = \frac{A_s^{\text{crit}}}{1 - \frac{1}{3} \beta_{ht} S_t \tilde{t}}.
\] (19)

Therefore radiative corrections magnify the critical point that \( A_s \) must exceed to develop a large trilinear coupling minimum. If the tree-level value of \( A_s \) is close to the critical point, radiative corrections can destabilize the vacuum. For example, for \( m_1^2 = m_2^2 = m_s^2 = m_0 > 0 \) one gets \( A_s^{\text{crit}} \sim 2.82 m_0 \). While for \( A_s < A_s^{\text{crit}} \), all VEV’s identically vanish, for \( A_s > A_s^{\text{crit}} \) all VEV’s are equal and get the value of \( A_s/(h_s \sqrt{2}) \). Even if \( A_s > A_s^{\text{crit}} \), when it is in the vicinity of \( A_s^{\text{crit}} \), radiative corrections cause a transition from the broken phase to the symmetric phase, and completely restores the symmetry. However, for large enough \( A_s \), there occurs no problem in the implementation of the minimum of the potential under concern. Thus, by choosing \( A_s \) large enough one can analyze the physical processes with reliable accuracy using the tree-level potential.

Recently, in [13], electroweak breaking in E(6)-based \( Z' \) models together with radiative corrections to scalar masses was discussed. Here we consider a general string-inspired family-dependent \( Z' \) model without exotics [5] in which we do not restrict ourselves to E(6) charge assignments. One notes that the minimum value of \( h_s \), \( h_s^{\text{min}} = 0.345 \), which we have used in forming the graphs turns out to be the ‘typical’ value for \( h_s \) in such models.

In all extensions of the SM namely, 2HDM, MSSM, NMSSM, and \( Z' \) models, \( hZZ \) and \( hAZ \) couplings are smaller than those of the SM due to the mixings in the Higgs sector (and also in the vector boson sector in \( Z' \) models) of the model, and this cause a reduction in the associated cross sections compared to those in the SM. For example, in MSSM [16], lightest Higgs production cross section is suppressed by the existence of the factor \( \sin^2(\beta - \alpha) \), where \( \alpha \) is the CP-even Higgs mixing angle. In addition to the mixing-induced reduction in the cross sections, one has a multitude of scalar particles which make it hard to differentiate between the underlying models.
Let us first summarize the situation in the $Z'$ models and then discuss the other models briefly. As was discussed in [6], the Higgs pair production process allows only $h_2$ production in the limit of large Higgs trilinear coupling. Unlike this, however, the Bjorken process discussed here forbids the $h_2$ production, which make two processes complementary to each other. Hence in the Higgs-strahlung type processes, as the collider energy sweeps a certain range of values, one can get at most two CP-even scalars at the final state.

In MSSM, there does not exist a minimum of the potential like the one discussed here. However, in the limit of large Higgs bilinear mass term, $\mu B$ (decoupling limit), there is a single light scalar in the spectrum, all other scalars being much heavier. In this sense, MSSM mimics the SM at low energies concerning their Higgs sectors. In the opposite limit, one has a relatively light Higgs spectrum consisting of two CP-even and one CP-odd scalar, whose existence are to be probed at the colliders. In the absence of similarities between $Z'$ models and MSSM in the large Higgs trilinear coupling minimum of the former, the only observable difference between the two models relies on the number of scalars they predict.

In NMSSM, Higgs sector is larger than that of the MSSM, due to the existence of the SM-singlet superfield $S$. Indeed, NMSSM has three CP-even, and two CP-odd scalars which make it completely different than MSSM and $Z'$ models (due to one extra CP-odd scalar). However, in discussing the Higgs-strahlung process (where CP-odd Higgs effects are absent) NMSSM and $Z'$ models are expected to behave in a similar manner as they have the same number of CP-even scalars. Given the experimental signature of the Higgs-strahlung process under concern over a certain energy range, is it possible to differentiate between NMSSM and $Z'$ models? In the limit of large Higgs trilinear coupling, it is possible to have a parallel analysis of these two models from which one may conclude on the dissimilarities between the two. It is clear that to have a proper large Higgs trilinear coupling minimum, in which all scalars have the VEV’s of similar magnitude, it is necessary to suppress the contribution of $S^3$ term to the soft supersymmetry breaking part of the scalar potential. With this restriction in mind and assuming a small but non-vanishing Yukawa coupling, $k$, for the $S^3$ term (to forbid the creation of an axion) in the superpotential, one finds the following $hZZ$ couplings,

$$K_{111} = \frac{G M_Z}{2}, \quad K_{211} = 0, \quad K_{311} = 0 \quad (20)$$
corresponding to the first three couplings in (4). The mass spectrum now is given by

\[ m_{h_1} \sim \sqrt{2hS}M_Z \frac{G}{G}, \quad m_{h_2} \sim \sqrt{2h^2 + 4k^2}M_Z \frac{G}{G}, \quad m_{h_3} \sim \sqrt{G^2 + 2h^2}M_Z, \quad (21) \]

so that the lightest Higgs has the same mass in both models. However, as dictated by (20), in the NMSSM only the lightest Higgs contributes to the Higgs-strahlung process under discussion. This forms a spectacular difference between the two models: while one expects two CP-even Higgs particles at the final state in \( Z' \) models, in NMSSM there is just one Higgs particle.

Despite the negative Higgs search at LEP2, the Bjorken process discussed here will continue to be one of the main Higgs search tools in future colliders. Therefore it seems necessary to have the predictions of various models in hand for comparison with the near-future collider data. In conclusion, we have analyzed Higgs production through the Higgs-strahlung type processes in \( Z' \) models in that minimum of the potential for which Higgs trilinear coupling is large compared to the other soft mass parameters. At the same type of minimum, NMSSM is found to have observable differences with the \( Z' \) models.

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Figure 1: Dependence of $\sigma_{h_1}$ on $h_S$ for $\sqrt{s} =$ 130, 136, 161, 172, 183 and 205 GeV from left to right.
Figure 2: Dependence of $\sigma_{h_{3}}$ on $h_{S}$ for $\sqrt{s} = 183$ GeV and $\rho_{S} = 0.58$ (solid), 0.8 (dashed), 1.0 (dotted) and 1.17 (dot-dashed).