Charged Higgs Sector with and without R–Parity

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Abstract

The simplest way of studying systematically R–parity violating phenomena is by introducing a bilinear term in the superpotential of the type $\epsilon \hat{L}\hat{H}_2$, which violates R–parity and lepton number but keep barion number conserved. In its simplest version, this “$\epsilon$–model” is a two parameter extension of the MSSM and a one parameter extension of the MSSM–SUGRA. Here we study the charged Higgs sector of the model, which mixes with the stau sector, and compare it with the charged Higgs sector of the MSSM. We demonstrate that $m_{H^\pm}$ can be lower than $m_W$ already at tree level, and calculate the production cross section of pairs of charged Higgs and staus. In this model it is possible the mixed production of a charged Higgs and a stau and its production cross section can be sizable. We finally comment about the new R–parity violating decay modes of the charged scalars.

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1 Introduction

After the discovery of the top quark at Fermilab [1], the Higgs boson is the only particle predicted by the Standard Model (SM) still undetected. Theory does not predict the mass $m_H$ of the Higgs boson, but it should satisfy $m_H < \sim 1$ TeV, otherwise the Higgs sector becomes non-perturbative. On the other hand, an experimental lower bound of $m_H > 77$ GeV has been set by LEP [2] due to its non-observation in $e^+e^-$ collisions.

Despite its simplicity, the SM Higgs sector is not attractive theoretically because the Higgs boson mass is unstable under radiative corrections. The most popular extension of the SM that solves this problem is the Minimal Supersymmetric Standard Model (MSSM), whose Higgs sector contains two CP–even neutral Higgs bosons $h$ and $H$, a CP–odd Higgs boson $A$, and a pair of charged Higgs bosons $H^\pm$.

Here we are interested in the phenomenology of the charged Higgs boson [3] in a model which violates R–parity through a bilinear term of the form $\varepsilon_{ab}\hat{L}_a^i\hat{H}^b_i$ in the superpotential [4]. The superpotential we consider is

$$W = \varepsilon_{ab}\left[h^{ij}_U\hat{Q}^i_j\hat{H}^b_2 + h^{ij}_D\hat{D}^i_j\hat{H}^a_1 + h^{ij}_E\hat{L}^i_j\hat{H}^a_1 - \mu\hat{H}^a_1\hat{H}^b_2 + \varepsilon_3\hat{L}_1^i\hat{H}^a_2\right]$$

(1)

where the first four terms correspond to the MSSM and the last one is the bilinear term which violates R–parity [4]. From now on, the model described by the superpotential in eq. (1) will be called here the $\varepsilon$–model. For simplicity we consider the case where $\varepsilon_1 = \varepsilon_2 = 0$ and $\varepsilon_3 \neq 0$. This model is a truncated version of a more complete model in which a vacuum expectation value of a right–handed sneutrino field induces the $\varepsilon$–term [6].

In the $\varepsilon$–model, a vacuum expectation value $v_3$ of the left–handed sneutrino–tau field is induced by the $\varepsilon$–term and, therefore, the lepton number is not conserved. In addition, the tau–neutrino mixes with the neutralinos and the $\nu_\tau$ acquires a mass which satisfy $m_{\nu_\tau} \sim (\mu v_3 + \varepsilon_3 v_1)^2$, where $v_1$ is the vev of $H_1$ and $v_3$ is the vev of $\tilde{\nu}_\tau$. We need a small value of the neutrino–tau mass in order to satisfy the experimental upper bound on $m_{\nu_\tau}$, and this is achieved without a fine tuning on the combination $(\mu v_3 + \varepsilon_3 v_1)$. The reason is that with universality of scalar masses and universality of bilinear soft terms, the smallness of the combination $(\mu v_3 + \varepsilon_3 v_1)$ is natural and, in fact, proportional to $h_\nu^2/8\pi^2$, where $h_\nu$ is the bottom quark Yukawa coupling. When embedded into supergravity, this $\varepsilon$–model is a one parameter extension of the MSSM–SUGRA and, therefore, the simplest way to study systematically R–parity violating phenomena [7].

The $\varepsilon$–term in eq. (1) cannot be rotated away by a redefinition of the superfields $\hat{H}_1^i = (\mu\hat{H}_1^i - \varepsilon_3\hat{L}_3^i)/\sqrt{\mu^2 + \varepsilon_3^2}$. If this rotation is performed, the bilinear term disappears from the superpotential, but trilinear R–parity violating terms are re–introduced. Furthermore, if supersymmetry is broken, bilinear terms which induce a vacuum expectation value of the sneutrino–tau are re–introduced in the scalar sector. This occurs even if universality
of scalar masses and universality of bilinear soft breaking terms hold at the unification scale.

Also important in the \( \epsilon \)-model is the fact that the charginos mix with the tau lepton, the CP–even Higgs bosons mix with the real part of the \( \tilde{\nu}_\tau \) field, the CP–odd Higgs bosons mix with the imaginary part of \( \tilde{\nu}_\tau \), and the charged Higgs bosons mix with the staus. In this note we analyze some aspects of the charged Higgs sector, first within the MSSM, and later in the \( \epsilon \)-model.

## 2 Charged Higgs in the MSSM

The MSSM charged Higgs boson has a one–loop corrected mass given by \cite{8, 9}

\[
m_{H^\pm}^2 = m_W^2 + m_A^2 + \text{Re}[A_{H^+H^-}(m_{H^\pm}^2) - A_{WW}(m_W^2) - A_{AA}(m_A^2)],
\]

where the charged Higgs mass \( m_{H^\pm} \) is the solution to an implicit equation. The tree level approximation is adequate in a large region of parameter space, specially if \( m_A \gg m_Z \). Nevertheless, there are regions where one–loop contributions are crucial. Furthermore, it is possible that the charged Higgs mass is smaller than the tree level lower bound given
by $m_{H^\pm} > m_W$. As an example, we consider the case in Fig. 1. In this case we take as an input the charged Higgs mass $m_{H^\pm} = 70$ GeV and calculate the one–loop corrected CP–odd Higgs mass $m_A$. This is done simply by solving the implicit equation for $m_A$ given in eq. (2). Note that this value of the charged Higgs mass is not possible at the classical level, because in that case $(m_A^2)_{\text{tree}}$ is negative. But quantum corrections are large enough to lift $m_A^2$ to positive values, as we can see for three different choices of the parameters that control the squark sector. Two of the curves are truncated after taking into account the constraints from color breaking minima [10].

In the MSSM the charged Higgs pair production cross section depends only on the charged Higgs mass. Nevertheless, radiative corrections can be important [11, 12]. For the same choice of parameters in the previous figure, we have plotted the one–loop corrected cross section $\sigma(e^+e^- \rightarrow H^+H^-)$ in Fig 2. Corrections from $-15\%$ up to $+50\%$ can be appreciated, and may have an impact on charged Higgs searches at LEP2.

Charged Higgs decay modes in the MSSM are well known. If $\tan \beta$ is small $H^+ \rightarrow \ell \overline{\nu}_\ell$ dominates, and if $\tan \beta$ is large $H^+ \rightarrow \tau^+ \nu_\tau$ is the dominant decay mode. The decay mode $H^+ \rightarrow \ell \overline{\nu}_\ell$ should also be considered. QCD corrections to the decay width can be implemented with the running quark masses [13], and electroweak radiative corrections are small [14].

Figure 2: Tree level (dashes) and radiatively corrected (solid) total production cross section of a pair of charged Higgs bosons at LEP2 as a function of $\tan \beta$ in the MSSM.
3 Charged Higgs and Staus in the $\epsilon$–Model

In the presence of a bilinear R–parity violating term in the superpotential, as given in eq. (1), the phenomenology of the charged Higgs sector changes dramatically. First of all, the charged Higgs boson fields mix with the stau sector. The mass terms in the lagrangian are

$$V_{\text{quadratic}} = [H_1^+, H_2^-, \tilde{\tau}_L^-, \tilde{\tau}_R^-] M_{S\pm}^2 [H_1^+, H_2^+, \tilde{\tau}_L^+, \tilde{\tau}_R^+] + ...$$  (3)

where, for simplicity we decompose the charged scalar mass matrix into three terms

$$M_{S\pm}^2 = M_{H\pm}^2 + M_{\tau\pm}^2 + M_{\epsilon_3}^2.$$  (4)

In the first matrix $M_{H\pm}^2$ we include the MSSM terms relevant to the charged Higgs sector, and it is given by ($g_Z^2 \equiv g^2 + g'^2$):

$$M_{H\pm}^2 = \begin{pmatrix}
    m_{H_1}^2 + \mu^2 + \frac{1}{4}g^2v_2^2 + \frac{1}{8}g_Z^2(v_1^2 - v_2^2) & B\mu + \frac{1}{4}g^2v_1v_2 & 0 & 0 \\
    B\mu + \frac{1}{4}g^2v_1v_2 & m_{H_2}^2 + \mu^2 + \frac{1}{4}g^2v_2^2 - \frac{1}{8}g_Z^2(v_1^2 - v_2^2) & 0 & 0 \\
    0 & 0 & m_{\tau\pm}^2 + \mu^2 + \frac{1}{4}g^2v_1^2 - \frac{1}{8}g_Z^2(v_1^2 - v_2^2) & 0 \\
    0 & 0 & 0 & m_{\epsilon_3}^2 + \frac{1}{4}g^2v_1^2 - \frac{1}{8}g_Z^2(v_1^2 - v_2^2)
\end{pmatrix}$$  (5)

Similarly, we include in $M_{\tau\pm}^2$ the MSSM terms relevant for the stau sector, and they are

$$M_{\tau\pm}^2 = \begin{pmatrix}
    0 & 0 & 0 & 0 \\
    0 & 0 & 0 & 0 \\
    0 & m_{L\pm}^2 + \frac{1}{4}h^2v_1^2 - \frac{1}{8}(g^2 - g'^2)(v_1^2 - v_2^2) & \frac{1}{\sqrt{2}}h_\tau(A_\tau v_1 - \mu v_2) & \frac{1}{\sqrt{2}}h_\tau(A_\tau v_1 - \mu v_2) \\
    0 & m_{L\pm}^2 + \frac{1}{4}h^2v_1^2 - \frac{1}{8}(g^2 - g'^2)(v_1^2 - v_2^2) & m_{R\pm}^2 + \frac{1}{4}h^2v_1^2 - \frac{1}{8}g_Z^2v_1^2 & \frac{1}{\sqrt{2}}h_\tau(A_\tau v_1 - \mu v_2)
\end{pmatrix}$$  (6)

Finally, the terms not present in the MSSM and induced by the $\epsilon_3$–term in the superpotential are given by

$$M_{\epsilon_3}^2 = \begin{pmatrix}
    \frac{1}{4}(g^2 - 2h_\tau^2)v_1v_3 - \mu \epsilon_3 & 0 & \frac{1}{4}(g^2 - 2h_\tau^2)v_1v_3 - \mu \epsilon_3 & -\frac{1}{\sqrt{2}}h_\tau(\epsilon_3 v_2 + A_\tau v_3) \\
    0 & \epsilon_3^2 + \frac{1}{8}(g^2 - g'^2)v_3^2 & -B_2\epsilon_3 + \frac{1}{4}g^2v_3^2 & -\frac{1}{\sqrt{2}}h_\tau(\mu v_3 + \epsilon_3 v_1) \\
    \frac{1}{4}(g^2 - 2h_\tau^2)v_1v_3 - \mu \epsilon_3 & -B_2\epsilon_3 + \frac{1}{4}g^2v_3^2 & \epsilon_3^2 + \frac{1}{8}(g^2 + g'^2)v_3^2 & 0 \\
    -\frac{1}{\sqrt{2}}h_\tau(\epsilon_3 v_2 + A_\tau v_3) & -\frac{1}{\sqrt{2}}h_\tau(\mu v_3 + \epsilon_3 v_1) & 0 & \frac{1}{4}h^2v_3^2 - \frac{1}{8}g_Z^2v_3^2
\end{pmatrix}$$

which is identical to zero in the MSSM limit where $\epsilon_3 = \nu_3 = 0$. The charged scalar mass matrix $M_{S\pm}^2$ in eq. (1) is diagonalized by a rotation defined by diag($0, m_{H\pm}^2, m_{\tau\pm}^2, m_{\epsilon_3}^2$) = $R_{S\pm}M_{S\pm}^2R_{S\pm}^T$, and the mass eigenstates are

$$\begin{pmatrix}
    S_1^+ \\
    S_2^+ \\
    S_3^+ \\
    S_4^+
\end{pmatrix} \equiv \begin{pmatrix}
    G^+ \\
    H^+ \\
    \tilde{\tau}_1^+ \\
    \tilde{\tau}_2^+
\end{pmatrix} = R_{S\pm} \begin{pmatrix}
    H_1^+ \\
    H_2^+ \\
    \tilde{\tau}_L^+ \\
    \tilde{\tau}_R^+
\end{pmatrix}$$  (8)
where $G^+$ is the massless Goldstone boson, $\tilde{\tau}_1^+$ and $\tilde{\tau}_2^+$ are the two mass eigenstates with the largest stau component, and the remaining mass eigenstate is the charged Higgs $H^+$. The determinant of the matrix $M^2_{S\pm}$ is explicitly zero only after imposing the minimization conditions, or tadpole equations, which in the $\epsilon_3$-model are equal to

$$
\begin{align*}
\tau_0^1 &= (m_{H_1}^2 + \mu^2)v_1 - B\mu v_2 - \mu\epsilon_3 v_3 + \frac{1}{8}(g^2 + g'^2)v_1(v_1^2 - v_2^2 + v_3^2), \\
\tau_0^2 &= (m_{H_2}^2 + \mu^2 + \epsilon_3^2)v_2 - B\mu v_1 + B_2\epsilon_3 v_3 - \frac{1}{8}(g^2 + g'^2)v_2(v_1^2 - v_2^2 + v_3^2), \\
\tau_0^3 &= (m_{L_3}^2 + \epsilon_3^2)v_3 - \mu\epsilon_3 v_1 + B_2\epsilon_3 v_2 + \frac{1}{8}(g^2 + g'^2)v_3(v_1^2 - v_2^2 + v_3^2),
\end{align*}
$$

at tree level. At the minimum we must impose $\tau_0^i = 0$, $i = 1, 2, 3$.

We calculate the charged Higgs mass $m_{H^\pm}$ and we find that already at tree level this mass can be lighter than $m_W$, as opposed to the MSSM case. This can be understood if we work in the approximation $|v_3| \ll |\epsilon_3| \ll m_W$ and $b_\tau = 0$, where we find

$$
m_{H^\pm}^2 - m_A^2 \approx m_W^2 + (\mu^2 + B_2^2) \left[ \frac{1}{m_{H^\pm}^2 - m_{\tilde{\tau}_L}^2} - \frac{1}{m_{A^0}^2 - m_{\tilde{\nu}_L}^2} \right] \epsilon_3^2.
$$

Here the masses $m_{H^\pm}$, $m_{\tilde{\tau}_L}$, $m_{A^0}$, and $m_{\tilde{\nu}_L}$ are calculated in the limit $v_3 = \epsilon_3 = 0$, i.e., in the MSSM. Using standard relations it can be shown that $m_{H^\pm}^2 - m_A^2 < m_W^2$ already at tree level. This can be appreciated in Fig. 3, where we make a scan of $10^3$ points in

Figure 3: Tree level and one–loop charged Higgs mass as a function of $m_A$ in the $\epsilon_3$–model.
Figure 4: Possible values of the one–loop corrected charged Higgs mass as a function of tan β for different upper bound for $\epsilon_3$ and $v_3$.

Parameter space: some of the tree level points are below the dashed line corresponding to the MSSM lower limit $m_{H^\pm} = m_W$. We impose that the induced tau–neutrino mass is smaller than 30 MeV.

In the same Fig. 3 we plot the one–loop renormalized charged Higgs mass. To calculate it, we work in the MSSM approximation where we add to the tree level mass the correction given in eq. (2). The number of points below the MSSM lower limit increases after adding radiative corrections because these corrections are negative if tan β is small. There are also many points where quantum corrections are large and positive. This happens when tan β is large.

The effect that the $\epsilon_3$–term has on the charged Higgs mass can be appreciated also in Fig. 4. Here we make a scan over parameter space with $10^4$ points, and the curves are the boundary below which no points are found (“forbidden” region). In solid line we take $\epsilon_3 < 0.5$ and $v_3 < 0.5$ GeV and the situation is pretty close to the MSSM case. The minimum value of the charged Higgs mass is around $m_W$, but radiative corrections decrease this minimum value when tan β is small, and increase it if tan β is very large. As soon as the values of $\epsilon_3$ and $v_3$ are increased, the charged Higgs mass can take lower and lower values, as indicated by the other curves. We stress the fact that the “forbidden” regions are in reality regions where the scan did not find any solution.
Figure 5: Total production cross section of a pair of (a) charged Higgs and (b) light staus, as a function of their mass.

The production cross sections of pairs of charged scalars are also modified by the $\epsilon_3$ term. In Fig. 5, we plot the charged Higgs pair production cross section $\sigma(e^+e^- \rightarrow H^+H^-)$ as a function of the charged Higgs mass $m_{H^\pm}$, made with a scan of $4 \times 10^4$ points in parameter space. Most of the points fall into the MSSM curve, but there are some deviations due to charged Higgs fields mixing with right–stau. This can be easily understood if we look at the $Z$ coupling to a pair of charged scalars. The $ZS_i^+S_j^-$ Feynman rule is equal to $i\lambda^{ij}_{ZS+i}(p + p')^\mu$ where $p$ and $p'$ are the momenta in the direction of the positive electric charge flow. The $\lambda$ couplings are equal to $\lambda_{ZS+i} = R_{S+i} R_{S-i}^T$ where the couplings in the unrotated basis are

$$\lambda'_{ZS+i} = \frac{g}{2c_W} \begin{pmatrix} -c_{2W} & 0 & 0 & 0 \\ 0 & -c_{2W} & 0 & 0 \\ 0 & 0 & -c_{2W} & 0 \\ 0 & 0 & 0 & 2s_{2W}^2 \end{pmatrix}$$

(11)

If the right stau were decoupled from the rest of the charged scalars, the charged Higgs pair production cross section would be identical to the MSSM for any value of $\epsilon_3$ or $v_3$. The reason is that the upper–left $3 \times 3$ relevant sub-matrix of $\lambda'_{ZS+i}$ in eq. (11) is proportional to the identity.

In Fig. 5b, we plot the total production cross section of a pair of light staus $\tilde{\tau}_1^\pm$. 
The points concentrate around the two MSSM curves corresponding to the production of left–staus (upper curve) and right–staus (lower curve). The cross sections, calculated for a center of mass $\sqrt{s} = 192$ GeV, have a maximum value of 0.6 pb for any of the charged scalars.

It is very interesting to notice that, contrary to what happens in the MSSM, the mixed production of a charged Higgs and a stau is allowed in the $\epsilon$–model. In Fig. 6 we plot the mixed cross section $\sigma(e^+e^- \rightarrow H^\pm\tilde{\tau}_1^\mp) \equiv \sigma(e^+e^- \rightarrow H^+\tilde{\tau}_1^-) + \sigma(e^+e^- \rightarrow H^-\tilde{\tau}_1^+)\)$ as a function of the total mass of the products $m_{H^\pm} + m_{\tilde{\tau}_1^\pm}$ for a center of mass energy $\sqrt{s} = 192$ GeV. The mixed cross section has a sizable maximum value of 0.125 pb, and as expected it is smaller than the maximum value of the pair production cross section.

Decay modes of charged scalar particles are modified with respect to the MSSM for two reasons. First, there are new R–parity violating channels like $H^+ \rightarrow \tilde{\chi}_1^0\tau^+$ and $H^+ \rightarrow \tilde{\chi}_1^+\nu_\tau$ for the charged Higgs, and $\tilde{\tau}_i^+ \rightarrow \nu_\tau\tau^+$ and $\tilde{\tau}_i^+ \rightarrow c\bar{s}$ for the staus. And second, the lightest supersymmetric particle is not stable. In the case of the neutralino as the LSP, its decay modes are $\tilde{\chi}_1^0 \rightarrow \nu_\tau Z^* \rightarrow \nu_\tau q\bar{q}(l\bar{l})$ and $\tilde{\chi}_1^0 \rightarrow \tau W^* \rightarrow \tau q\bar{q}(l\bar{l})$. Furthermore, the LSP need not to be the lightest neutralino, and if the LSP is the lightest stau, it can have R–parity violating decays $\tilde{\tau}_1^+ \rightarrow \nu_\tau\tau^+(c\bar{s})$ with a 100% branching ratio.
4 Conclusions

In the “$\epsilon$–model”, where R–parity and lepton number are violated by the introduction of a bilinear term of the type $\epsilon_{\alpha\beta}\bar{L}_i\tilde{H}_2^b$ in the superpotential, the charged Higgs sector mixes with the stau sector. The mass eigenstates are a set of four charged scalars: the unphysical massless Goldstone boson $G^\pm$, the charged Higgs boson $H^\pm$, and two staus $\tilde{\tau}_i^\pm$, $i = 1, 2$. We showed that the charged Higgs mass $m_{H^\pm}$ can be lower than $m_W$ even at tree level, contrary to what happens in the MSSM. Including radiative corrections, the charged Higgs mass can be as light as 45 GeV. Values of $\epsilon_3$ and $v_3$ of the order of 100 GeV are compatible with the induced tau–neutrino mass. New R–parity violating decay modes are allowed, where the charged Higgs can decay into charginos or neutralinos and staus can decay into quarks or leptons. The lightest supersymmetric particle need not to be $\tilde{\chi}_1^0$ because it is not stable. If $\tilde{\tau}_1^\pm$ is the LSP, then it can decay into the R–parity violating decay modes $\tilde{\tau}_1^+ \rightarrow \nu_\tau \tau^+ (c\bar{s})$ with 100% branching fraction, no matter how small the R–parity violating parameters $\epsilon_3$ and $v_3$ are. We claim that the “$\epsilon$–model” is the simplest way of studying systematically R–parity violating phenomena, and in its simplest version, it is a two parameter extension of the MSSM and a one parameter extension of the MSSM–SUGRA.

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