Physics Letters B 292 (1992) 329-336
North-Holland

PHYSICS LETTERS B

New Higgs signatures in supersymmetry with spontaneous broken R parity

J.C. Romão [1]

Centro de Física da Matéria Condensada, INIC, Av. Prof. Gama Pinto 2, P-1699 Lisbon Codex, Portugal

F. de Campos [2,3] and J.W.F. Valle [2,4]

Theoretical Physics Division, CERN, CH-1211 Geneva 23, Switzerland

Received 31 July 1992

Higgs production from Z decays in supersymmetry with spontaneous broken R parity proceeds mostly by the Bjorken process as in the standard model. However, the corresponding production rates can be weaker than in the standard model (SM), especially in the low mass region. This will substantially weaken the Higgs boson mass limits derived from LEP1. More strikingly, the main Higgs decay channel is "invisible", over most of the mass range accessible to LEP1, leading to events with large missing energy carried by majorons. This possibility should be taken into account in the planning of Higgs boson search strategies not only at LEP but also at high energy supercolliders.

1. Introduction

The problem of mass generation is one of the main puzzles in electroweak physics. Although the Higgs mechanism [1] has been suggested more than two decades ago, it was only very recently with the LEP experiments that one has started to probe sensitively the nature of the scalar sector [2]. The existence of supersymmetry (SUSY) at the electroweak scale is desirable as it can act as a stabilizing symmetry that naturally protects this scale against quantum corrections associated with superhigh scales. Although as yet there is no experimental support for SUSY the above argument has been taken as a strong motivation to carry out searches at higher energies. Unfortunately there is no clue as to how SUSY is realized. The most popular ansatz – called the minimal supersymmetric standard model (MSSM) [3] – assumes that SUSY is realized in the presence of a discrete R parity (R_p) symmetry under which all standard model particles are even while their SUSY partners are odd. This implies that SUSY particles must always be pair-produced, the lightest of them being stable. There is great interest in investigating theories without R parity [4]. There are two ways to break R_p: explicitly [5] and spontaneously [6,7]. The second provides a more systematic way to include R parity violating effects, that moreover automatically respects low energy *baryon number conservation* and evades restrictions based on cosmological baryogenesis arguments [8]. Although these could be avoided even in the explicit R_p breaking scenario [9], they are certainly avoided in the present model, inasmuch as the breaking of R parity sets in spontaneously only as an electroweak scale phenomenon. Here we concentrate entirely on this second possibility. There are two cases to consider, depending on whether lepton number is part of the gauge symmetry [7] or not [6]. In the first case there is a Z' gauge boson which acquires mass via the Higgs mechanism at a scale related to that which characterizes R parity violation. Here we focus on the simplest case where R_p is violated in the

[1] Bitnet address: ROMAO@PTIFM.
[2] Also at Instituto de Física Corpuscular – CSIC, Departament de Física Teòrica, Universitat de València, E-46100 Burjassot, Valencia, Spain.
[3] Bitnet address: CAMPOSC@EVALUN11; Decnet address: 16444::CAMPOSC.
[4] Bitnet address: VALLE@EVALUN11; Decnet address: 16444::VALLE.

absence of an additional gauge symmetry beyond the minimal SU(2)⊗U(1) structure. This possibility has been demonstrated to occur [10] in the model suggested in ref. [6] for many suitable values of the low energy parameters, consistent with observation. R parity breaking is driven by *isosinglet* slepton vacuum expectation values (VEVs) [6], so that the Goldstone boson (majoron) associated with spontaneous R parity breaking is mostly singlet and as a result the Z does not decay by majoron emission, in agreement with observation [2]. The R parity breaking scale typically lies in the phenomenologically interesting range ~ 10 GeV–1 TeV, leading to large rates for the associated R_p violating effects, which may well be accessible to experimental test [11–13]. It also leads to an interesting way to explain the solar neutrino data [14].

In this letter we focus on the neutral Higgs sector of these models and the corresponding implications for Higgs searches at LEP1. We find that both Higgs production rates and decay mechanisms may differ substantially from the SM and MSSM predictions. First, Higgs production rates can be substantially lower than in SM and MSSM and, second, its main decay mode can be invisible,

$$h \to J + J, \qquad (1)$$

where J denotes the majoron. This holds for a wide region of parameters and should also have important implications for the Higgs search strategies at high energy hadron supercolliders.

2. The model and the scalar potential

We consider the SU(2)⊗U(1) model defined by the superpotential [#1] terms [6]

$$\begin{aligned} &h_u u^c Q H_u + h_d d^c Q H_d + h_e e^c l H_d + \hat{\mu} H_u H_d \\ &+ (h_0 H_u H_d - \epsilon^2)\Phi + h_\nu \nu^c l H_u + h \Phi \nu^c S + M \nu^c S \\ &+ M_\Phi \Phi\Phi + \lambda \Phi^3 . \qquad (2) \end{aligned}$$

The first five terms are the usual ones that define the R_p conserving MSSM. The fifth term ensures that electroweak symmetry breaking can take place at the tree level, as in ref. [15]. The last five terms involve SU(2)⊗U(1) superfields (Φ, ν_i^c, S_i) carrying a conserved lepton number assigned as (0, −1, 1) respectively. Such singlets arise in several extensions of the standard model and may lead to interesting phenomenological signatures of their own [16–18]. Their presence here is essential in order to drive the spontaneous violation of R parity and electroweak symmetries in a phenomenologically consistent way [6]. The bilinear $H_u H_d$ term plays an important role in giving more flexibility in the minimization of the Higgs potential while at the same time obeying all experimental constraints especially the chargino mass limit from LEP. The terms Φ^2, Φ^3 and $\nu^c S$ do not play any important role for our present considerations and will be ignored. We also assume that the coupling matrices $h_{\nu ij}$ and h_{ij} are nonzero only for the third generation and set $h_\nu \equiv h_{\nu 33}$ and $h \equiv h_{33}$. With this assumption we are studying effectively a one generation model. We are well aware that a phenomenologically consistent model requires the presence of flavour nondiagonal couplings such as $h_{\nu 23}$, needed in order to ensure that the massive ν_τ decays fast enough [4] to obey cosmological limits. This has been shown to be the case due to the existence of the majoron emission decay channel $\nu_\tau \to \nu_\mu + J$. However for our present purposes the effective one-generation model approach will be enough.

To complete the specification of the model we give the form of the full scalar potential along neutral directions

$$\begin{aligned} V_{\text{total}} = &|h\Phi\tilde{S} + h_\nu \tilde{\nu} H_u|^2 + |h_0 \Phi H_u + \hat{\mu} H_u|^2 \\ &+ |h\Phi\tilde{\nu}^c|^2 + |-h_0 \Phi H_d - \hat{\mu} H_d + h_\nu \tilde{\nu}\tilde{\nu}^c|^2 \\ &+ |-h_0 H_u H_d + h\tilde{\nu}^c\tilde{S} - \epsilon^2|^2 + |h_\nu \tilde{\nu}^c H_u|^2 \\ &+ \tilde{m}_0[-A(-h\Phi\tilde{\nu}^c\tilde{S} + h_0 \Phi H_u H_d - h_\nu \tilde{\nu} H_u \tilde{\nu}^c) \\ &+ (1-A)\hat{\mu} H_u H_d + (2-A)\epsilon^2 \Phi + \text{h.c.}] \\ &+ \sum_i \tilde{m}_i^2 |z_i|^2 + \alpha(|H_u|^2 - |H_d|^2 - |\tilde{\nu}|^2)^2 , \qquad (3) \end{aligned}$$

where $\alpha = \frac{1}{8}(g^2 + g'^2)$ and z_i denotes any neutral scalar field in the theory.

The pattern of spontaneous symmetry breaking of both electroweak and R parity symmetries has been studied in ref. [10]. There it was demonstrated ex-

[#1] All couplings h_u, h_d, h_e, h_ν, h are described by arbitrary matrices in generation space. Note that we have added some new terms that were not included in ref. [6] because they are allowed by our symmetries.

plicitly that, for suitable values of the low energy parameters consistent with observation, the energy is minimum when both R parity and electroweak symmetries are spontaneously broken. Electroweak breaking is driven by the isodoublet VEVs $v_u = \langle H_u \rangle$ and $v_d = \langle H_d \rangle$, assisted by the VEV v_F of the scalar in the singlet superfield Φ. The combination $v^2 = v_u^2 + v_d^2$ is fixed by the W mass,

$$m_W^2 \simeq \tfrac{1}{2} g^2 (v_u^2 + v_d^2) , \tag{4}$$

while the ratio of isodoublet VEVs determines the important parameter

$$\tan\beta = \frac{v_u}{v_d} . \tag{5}$$

This basically recovers the standard tree level spontaneous breaking of the electroweak symmetry in the MSSM. On the other hand the spontaneous breaking of R parity is driven by nonzero VEVs for the scalar neutrinos. The scale characterizing R parity breaking is set by the isosinglet VEVs

$$v_R = \langle \tilde{\nu}_\tau^c \rangle , \quad v_S = \langle \tilde{S}_\tau \rangle , \tag{6,7}$$

where $V = \sqrt{v_R^2 + v_S^2}$ and can lie anywhere in the range ~ 10 GeV–1 TeV. A necessary ingredient for the consistency of this model is the presence of a small seed of R parity breaking in the SU(2) doublet sector,

$$v_L = \langle \tilde{\nu}_{L\tau} \rangle . \tag{8}$$

The spontaneous R parity breaking above also entails the spontaneous violation of *total* lepton number (conserved by eq. (2)) leading to the existence of the majoron, given by

$$\frac{v_L^2}{V v^2} (v_u H_u - v_d H_d) + \frac{v_L}{V} \tilde{\nu}_\tau - \frac{v_R}{V} \tilde{\nu}_\tau^c + \frac{v_S}{V} \tilde{S}_\tau . \tag{9}$$

Astrophysical considerations [19] related to stellar cooling by majoron emission require a small value of $v_L \sim 100$ MeV, which is also naturally obtained from the minimization of the Higgs potential [#2].

Before entering the discussion of Higgs production and decay mechanisms we need to complete the specification of the scalar potential in eq. (3) by including the relevant radiative corrections. As has been widely pointed out [20,21] there are potentially important radiative corrections to the Higgs scalar mass matrices coming from loops involving heavy quarks and squarks. To take them into account we followed the effective potential method used by Ellis, Ridolfi and Zwirner [21] for the case of the MSSM. The one loop effective potential at a scale Q (which we take to be M_Z) is given by

$$V_{\text{eff}}(Q) = V(Q) + \Delta V(Q) , \tag{10}$$

where $V(Q)$ is the tree level potential given in eq. (3) and

$$\Delta V(Q) = \frac{1}{64\pi^2} \operatorname{Str} \mathcal{M}^4 \left(\ln \frac{\mathcal{M}^2}{Q^2} - \frac{3}{2} \right). \tag{11}$$

In the last equation Str denotes the Supertrace and $\mathcal{M}^2$ is the field-dependent generalized squared mass matrix for the model. For our purposes it is enough to consider only the contribution from the bottom and top quarks and their squarks. In the numerical calculations we also took the same mass for all the squarks. In this way we calculated $\Delta V(Q)$ as a function of the scalar fields of our model. The radiatively corrected effective potential obtained in this way was then minimized following the same procedure used in ref. [10], namely we first solved the extremum equations

$$\frac{\partial V_{\text{eff}}}{\partial z_i} = 0 , \tag{12}$$

where $z_i = (H_d, H_u, \tilde{\nu}, \tilde{\nu}^c, S, \phi)$. Then we evaluated the mass squared matrices for the real and imaginary parts of the fields. A point in parameter space is a minimum of the potential if all the eigenvalues of these matrices are positive (except for the Goldstone modes).

After including radiative corrections we found out that the basic conclusions of ref. [10] remain true in this case, namely that R_p violation may be energetically favored over R_p conserving minima. Our results also indicate that the mass of the lightest CP even state grows with m_{top} as in the MSSM if this eigenstate is mostly a doublet, as one would expect. However in our model this state is often a singlet where the above considerations do not apply. These points will be fur-

[#2] The marked hierarchy in the values of v_R and v_L follows because v_L is related to a Yukawa coupling h_ν and vanishes as $h_\nu \to 0$. This was shown to have important consequences for the propagation of solar neutrinos [14].

ther discussed in sections 3 and 4 where we present our results.

3. Higgs production

The neutral Higgs boson couplings to the Z boson arise in our model from two sources. In addition to the usual coupling to H_u and H_d there are couplings to $\tilde{\nu}$. They come from

$$\mathcal{L}_1 = -\frac{ig}{2\cos\theta_w} Z_\mu (H_d^{1*}\overleftrightarrow{\partial}^\mu H_d^1 - H_u^{2*}\overleftrightarrow{\partial}^\mu H_u^2)$$

$$-\frac{ig}{2\cos\theta_w} Z_\mu (\tilde{\nu}^*\overleftrightarrow{\partial}^\mu \tilde{\nu}) , \quad (13)$$

$$\mathcal{L}_2 = \frac{g^2}{4\cos^2\theta_w} Z_\mu Z^\mu (H_d^{1*}H_d^1 + H_u^{2*}H_u^2 + \tilde{\nu}^*\tilde{\nu}) . \quad (14)$$

Now we shift the fields according to

$$H_d = v_d + \frac{1}{\sqrt{2}} [\mathrm{Re}(H_d^0) + \mathrm{i}\,\mathrm{Im}(H_d^0)] , \quad (15)$$

$$H_u = v_u + \frac{1}{\sqrt{2}} [\mathrm{Re}(H_u^0) + \mathrm{i}\,\mathrm{Im}(H_u^0)] , \quad (16)$$

$$\tilde{\nu} = v_L + \frac{1}{\sqrt{2}} [\mathrm{Re}(\tilde{\nu}) + \mathrm{i}\,\mathrm{Im}(\tilde{\nu})] . \quad (17)$$

The cubic interactions arising from the $\mathcal{L}_2$ terms are rewritten as

$$\mathcal{L}_2 = \frac{g}{2\cos\theta_w} M_Z Z_\mu Z_\mu \cos\gamma$$
$$\times [\cos\beta\,\mathrm{Re}(H_d^0) + \sin\beta\,\mathrm{Re}(H_u^0) + \tan\gamma\,\mathrm{Re}(\tilde{\nu})] , \quad (18)$$

where $\tan\gamma = v_L/\sqrt{v_u^2+v_d^2}$ and β is the usual parameter in eq. (5).

We now rewrite eq. (18) in the basis of mass eigenstate scalar bosons. To do this we write the weak scalar fields as a vector $z_i \equiv (H_d, H_u, \tilde{\nu}, \tilde{\nu}^c, \tilde{S}, \phi)$. Then we set

$$z_i = v_i + \frac{1}{\sqrt{2}} (x_i + \mathrm{i}y_i) , \quad (19)$$

where $i=1, ..., 6$ are the fields defined in the multiplet. The mass eigenstates are then defined by

$$H_i = P_{ij}x_j , \quad A_i = Q_{ij}y_j , \quad (20)$$

where P and Q are orthogonal matrices [#3].

Here $i=1, ..., 6$ denotes any of the scalar bosons involved. In the mass eigenstates H_i and A_i the fields are ordered according to their increasing masses. In the 0^- sector, two of them are massless, one is the majoron, denoted by J, (A_2), the other (A_1) is the field eaten up by the Z. Substituting eq. (20) into eq. (18) and eq. (13) we find, respectively,

$$\mathcal{L}_2 = \frac{g}{2\cos\theta_w} M_Z Z_\mu Z_\mu H_i$$
$$\times [\cos\beta\, P_{i1} + \sin\beta\, P_{i2} + \tan\gamma\, P_{i3}]\cos\gamma , \quad (21)$$

$$\mathcal{L}_1 = \frac{g}{2\cos\theta_w} Z_\mu H_i \overleftrightarrow{\partial}^\mu A_j C_{ij} , \quad (22)$$

where $C_{ij} = P_{i1}Q_{j1} - P_{i2}Q_{j2} + P_{i3}Q_{j3}$.

In the MSSM limit these expressions reduce to the familiar form where $P_{11} = -\sin\alpha = -P_{22}$, $P_{12} = \cos\alpha = P_{21}$, $Q_{21} = \sin\beta$, $Q_{22} = \cos\beta$ and $P_{i3} = 0 = Q_{i3}$. As a consequence one recovers in that case the simple result $C_{12} \to -\cos(\alpha-\beta)$, $C_{22} \to -\sin(\alpha-\beta)$ and $\cos\beta\, P_{11} + \sin\beta\, P_{12} = \sin(\beta-\alpha)$.

The interactions in eq. (21) and eq. (22) above lead to the following Z decay channels:

$$Z \to H_i f\bar{f} \quad (23)$$

(Bjorken process), where $f = u, c, d, s, b, e, \mu, \tau$ and

$$Z \to H_i A_j . \quad (24)$$

In the MSSM we have $H_1 = h$, $H_2 = H$ and there is only one physical pseudoscalar boson $A_2 = A$ (there is no majoron in this limit).

It is instructive to compare the predictions on Higgs production of the present spontaneously broken R parity model with those of the MSSM. In the latter case there is a complementarity between the two types of decays in eq. (23) and eq. (24), and this has been taken as the basis of the experimental analysis used to place limits on the supersymmetric Higgs spectrum [22]. In the present model there is a breakdown in the sum rule involving the coupling strengths characterizing the processes in eq. (23) and eq. (24) with respect to MSSM expectations. As a result overall Higgs production rates can be weaker than in MSSM. In addition, we have checked that whenever the "associated" production in eq. (24) is kinemati-

[#3] For simplicity we assume CP conservation.

cally possible, it is dynamically suppressed because the pseudoscalar Higgs boson A is mostly an $SU(2)\otimes U(1)$ singlet and therefore too weakly coupled to the Z. As a result, Higgs production at LEP proceeds mostly by the Bjorken process as in the standard model, except for the possibly suppressed coupling.

Our results for Higgs production via the Bjorken process in the spontaneously broken R parity model are summarized in figs. 1 and 2. For definiteness we have fixed the value of the R parity breaking scale at $v_R = v_S = 100$ GeV and $\tan\beta = 3$. The SUSY and R_p parameters have been randomly varied in the following range: $-250 < \mu < 250$ GeV, $30 < M_2 < 250$ GeV, $M_1/M_2 = \frac{5}{3}\tan^2\theta_w$, $m_{\tilde{q}} = 1$ TeV, $m_\varphi = 10$ TeV, $10^{-5} < h_\nu < 10^{-1}$. The remaining parameters have been taken as in ref. [10]. In all our present analysis we have randomly varied the relevant unknown parameters over physically allowed regions only. For example, we have imposed the lower limit on the lightest chargino mass $m_{\chi^+} > 45$ GeV. Moreover, we have accepted only those points that are solutions which minimize the scalar potential of the theory. In fig. 1 we show the maximum allowed couplings characterizing the Bjorken mechanism, versus the lightest scalar Higgs mass m_h [#4]. Note that this coupling can be substantially weaker than in the MSSM because in our model also h can be predominantly an $SU(2)\otimes U(1)$ singlet and therefore too weakly coupled to the Z. This happens especially for lower mass values, which are therefore certainly allowed within the spontaneously broken R parity model [#5]. Indeed, in fig. 2 we illustrate how the resulting branching ratio for the production of leptons in association with the lightest scalar Higgs h can be much weaker than expected in the SM. Limits have been placed on $BR(Z \to l\bar{l}h)$ from the nonobservation of acoplanar lepton pair events with the 1990 LEP data sample [22] at the level of few $\times 10^{-5}$ (note that the limit becomes weaker for low masses $m_h \lesssim 10$ GeV). Combining all of the experiments one can estimate that $BR(Z \to l\bar{l}h) \lesssim 5 \times 10^{-6}$ [23]. From fig. 2 we see that, even if the ZZh coupling attains its maximal allowed

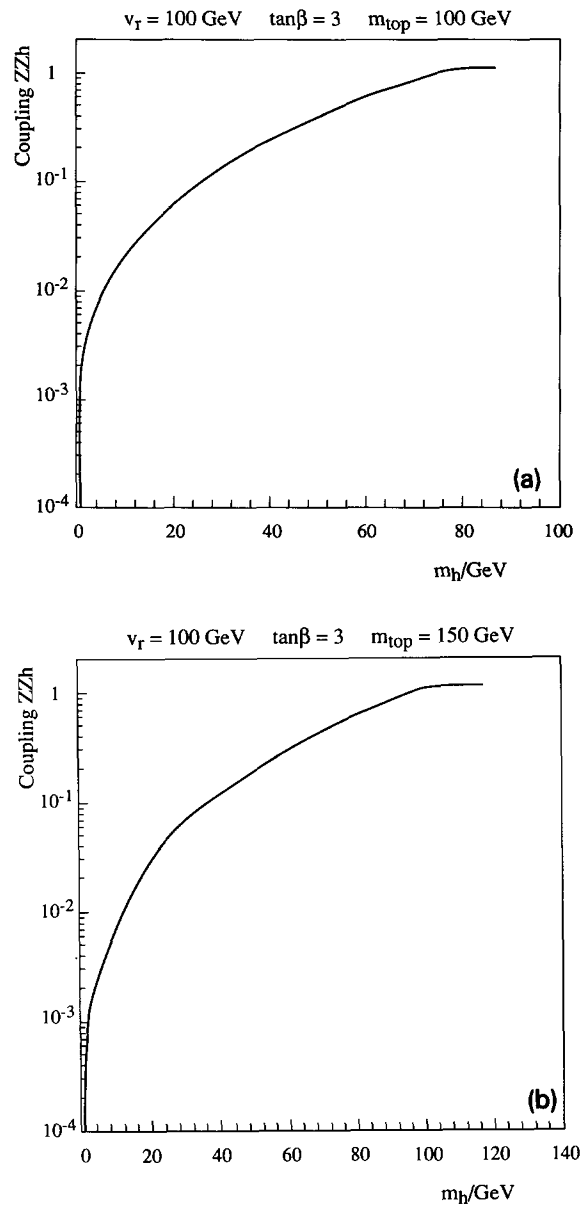

Fig. 1. (a) and (b) show the maximum ZZh coupling, normalized with respect to that of the SM, for $\tan\beta = 3$ and $m_{top} = 100$, 150 GeV respectively. All points underneath are possible in the model, and correspond to the unknown parameters being randomly varied as specified in the text, including only physically allowed regions and accepting only those points which are solutions that minimize the scalar potential of the theory.

values, LEP is barely starting to be sensitive to the Higgs, as predicted in this model. Similarly for the branching ratio for the associated production of quark

[#4] The production of the next-to-lightest scalar Higgs boson H is also possible but less likely.

[#5] Note that the Z lineshape, e.g., the contribution to Γ_Z^{inv} arising from $Z \to \nu\bar{\nu}h$ is negligible.

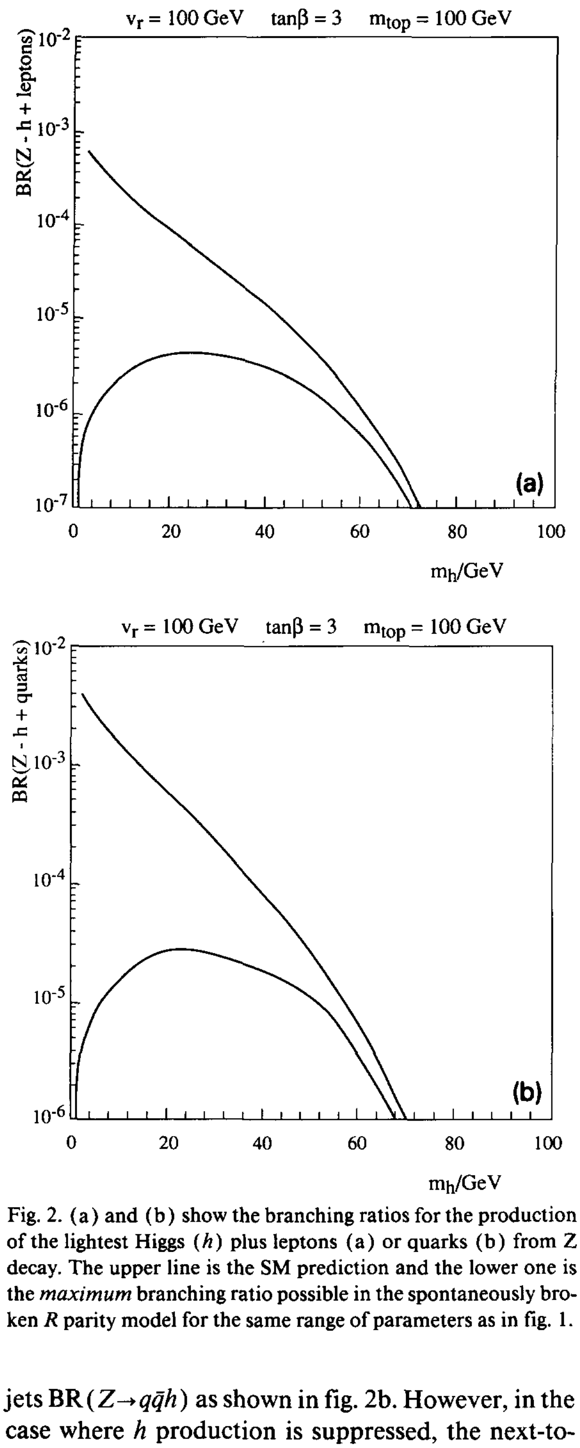

Fig. 2. (a) and (b) show the branching ratios for the production of the lightest Higgs (h) plus leptons (a) or quarks (b) from Z decay. The upper line is the SM prediction and the lower one is the *maximum* branching ratio possible in the spontaneously broken R parity model for the same range of parameters as in fig. 1.

jets BR$(Z\to q\bar{q}h)$ as shown in fig. 2b. However, in the case where h production is suppressed, the next-to-lightest of the scalar Higgses, H, couples to the Z with roughly the canonical strength expected for an iso-doublet Higgs. However, the H production rate will be suppressed with respect to that in the SM due to the smaller phase space available. For larger h masses, the maximum production rate of the lightest scalar Higgs boson attains roughly the canonical SM value.

4. Higgs decay

The most novel and characteristic aspect of the spontaneously broken R parity model is the existence of the massless pseudoscalar majoron, which follows from the spontaneous nature of R parity violation. This implies the presence of new Higgs decay channels with majoron emission [#6],

$$H_i \to JJ\,, \tag{25}$$

$$A_i \to H_i J\,. \tag{26}$$

In order to determine the rates for these decays we need to evaluate the corresponding trilinear couplings. By using the definition of z_i we rewrite everything in terms of mass eigenstates, using eq. (20). The resulting expressions are rather long but we have checked that in the MSSM limit defined by

$$h\to 0\,,\quad h_\nu \to 0\,,\quad h_0\to 0\,,\quad \hat{\mu}\to\mu\,, \tag{27}$$

one gets that the following expression for the cubic term of the interaction lagrangian

$$\mathscr{L} = -\frac{gM_Z}{4\cos\theta_w}\,[hAA\sin(\alpha+\beta)\cos 2\beta$$
$$-HAA\cos(\alpha+\beta)\cos 2\beta]\,, \tag{28}$$

in agreement with MSSM result.

In the present spontaneously broken R parity model the lightest CP-odd state is the massless majoron. However, because of its singlet nature, evident from eq. (9), all decays involving majoron emission in Z decays

$$Z\to H_i J \tag{29}$$

are automatically suppressed. The situation is quite

[#6] Note that the decays $A_i\to JJ$ are forbidden, since we assume CP conservation.

different for the majoron emitting Higgs decays. We now proceed to the determination of the effective H_iJJ coupling relevant for these decays. Substituting eq. (19) and eq. (20) in eq. (3) and keeping only the leading terms in eq. (9) we find

$$\begin{aligned} g_{H_iJJ} = -\sqrt{2}\Big[& P_{i6}\left(h^2 v_F + A h m_0 \frac{v_R v_S}{V^2}\right) \\ & + P_{i5} h^2 v_S \frac{v_R^2}{V^2} + P_{i4} h^2 v_R \frac{v_S^2}{V^2} + P_{i3} h_\nu^2 v_L \frac{v_R^2}{V^2} \\ & + P_{i2}\left(-h h_0 v_d \frac{v_R v_S}{V^2} + h_\nu^2 v_u \frac{v_R^2}{V^2}\right) \\ & - P_{i1} h h_0 v_u \frac{v_R v_S}{V^2}\Big]. \end{aligned} \tag{30}$$

The last two terms in eq. (30) indicate that in general both h and H have unsuppressed coupling to the majoron.

A simple calculation now enables us to obtain the Higgs decay rates to JJ and $b\bar{b}$. One finds

$$\Gamma(H_i \to b\bar{b}) = \frac{g_{H_ib\bar{b}}^2 m_{H_i}}{8\pi}, \tag{31}$$

$$\Gamma(H_i \to JJ) = \frac{g_{H_iJJ}^2}{32\pi m_{H_i}}, \tag{32}$$

so that one finds

$$\frac{\Gamma(H_i \to JJ)}{\Gamma(H_i \to b\bar{b})} = \frac{g_{H_iJJ}^2}{m_{H_i}^2} \frac{1}{g_{H_ib\bar{b}}^2}. \tag{33}$$

As a result we expect the "invisible" Higgs decay channel to be highly dominant, especially for low m_{H_i} values. We have studied the attainable values of the invisible Higgs decay branching ratio $\mathrm{BR_{inv}} = \mathrm{BR}(h \to JJ)$ for different m_h values in the range relevant for LEP studies. We found that the dominant Higgs decay mechanism is invisible, over most of the kinematical range available. Since the majoron is weakly coupled, it will escape undetected. Thus the $h \to JJ$ decay will lead to events with missing energy carried by majorons. For h mass values in a narrow range close to its maximum (~ 80 GeV for our choice of $\tan\beta$ and v_R) the h decay pattern approaches that of the SM and $\mathrm{BR_{inv}}$ becomes small. This maximum value of the h mass at which one changes from the novel regime of dominantly invisible h decay to the SM regime where $h \to b\bar{b}$ grows with the value of $\tan\beta$ and m_{top}.

5. Discussion

We have summarized the results of our study of the supersymmetric Higgs boson sector of the model with spontaneously broken R parity suggested previously [6]. We have shown that in this case Higgs production from Z decays proceeds mostly by the Bjorken process as in the standard model. The corresponding production rates can be weaker than in the standard model. As a result the LEP1 limit on the lightest scalar Higgs mass may be substantially weakened, especially in the low mass region. More strikingly, the main Higgs decay channel is likely to be "invisible", over most of the mass range accessible to LEP1, leading to events with large missing energy carried by majorons. This possibility should be considered in the planning of Higgs boson search strategies not only at LEP but also at high energy supercolliders.

Acknowledgement

This work was supported by CICYT (Spain) and by a CNPq fellowship (Brazil). We thank Michael Dittmar, Rohini Godbole and Roger Phillips for useful discussions.

References

[1] P.W. Higgs, Phys. Lett. 12 (1964) 132.
[2] J. Steinberger, in: Electroweak physics beyond the standard model, eds. J.W.F. Valle and J. Velasco (World Scientific, Singapore, 1992) p.3.
[3] H. Haber and G. Kane, Phys. Rev. 117 (1985) 75.
[4] J.W.F. Valle, Prog. Part. Nucl. Phys. 26 (1991) 91, and references therein.
[5] L. Hall and M. Suzuki, Nucl. Phys. B 231 (1984) 419.
[6] A. Masiero and J.W.F. Valle, Phys. Lett. B 251 (1990) 273.
[7] J.W.F. Valle, Phys. Lett. B 196 (1987) 157;
M.C. Gonzalez-Garcia and J.W.F. Valle, Nucl. Phys. B 355 (1991) 330.
[8] B.A. Campbell, S. Davison, J. Ellis and K. Olive, Phys. Lett. B 256 (1991) 457;
W. Fischler, G. Giudice, R. Leigh and S. Paban, Phys. Lett. B 258 (1991) 45.

[9] H. Dreiner and G.G. Ross, Oxford preprint OUTP-92-08P (1992).
[10] J.C. Romao, C.A. Santos and J.W.F. Valle, Phys. Lett. B 288 (1992) 311.
[11] P. Nogueira, J.C. Romao and J.W.F. Valle, Phys. Lett. B 251 (1990) 142.
[12] M.C. Gonzalez-Garcia, J.C. Romao and J.W.F. Valle, Nucl. Phys. B (1992), in press; Valencia preprint FTUV/91-42.
[13] J.C. Romao, N. Rius and J.W.F. Valle, Nucl. Phys. B 363 (1991) 369.
[14] J.C. Romao and J.W.F. Valle, Phys. Lett. B 272 (1991) 436; Nucl. Phys. B (1992), in press.
[15] R. Barbieri, S. Ferrara and C. Savoy, Phys. Lett. B 119 (1982) 343.
[16] J.W.F. Valle, Nucl. Phys. Proc. Suppl. 11 (1989) 118.
[17] R. Mohapatra and J.W.F. Valle, Phys. Rev. D 34 (1986) 1642.
[18] J.W.F. Valle, in: Weak and electromagnetic interactions in nuclei, ed. P. Depommier (Editions Frontières, Gif-sur-Yvette, 1989).
[19] J.E. Kim, Phys. Rev. 150 (1987) 1.
[20] H. Haber et al., Phys. Rev. Lett. 66 (1991) 1815;
R. Barbieri et al., Phys. Lett. B 258 (1991) 395.
[21] J. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B 262 (1991) 477.
[22] ALEPH Collab., D. Decamp et al., Phys. Lett. B 236 (1990) 233; B 237 (1990) 291; B 241 (1990) 141; B 245 (1990) 289; B 246 (1990) 306; preprints CERN-PPE/91-149;
DELPHI Collab., P. Abreu et al., Nucl. Phys. B 342 (1990) 1; Z. Phys. C 51 (1991) 25; preprint CERN-PPE/91-132;
L3 Collab., B. Adeva et al., preprint CERN-PPE/92-40, and references therein;
OPAL Collab., M.Z. Akrawy et al., Phys. Lett. B 236 (1990) 224; B 251 (1990) 211; B 253 (1991) 511;
OPAL Collab., P.D. Acton et al., Phys. Lett. B 268 (1991) 122.
[23] M. Dittmar, invited talk at Conf. on Search for new particles, status and prospects (ICTP, Trieste, May 1992).