Sneutrino–Higgs mixing in $WW$ and $ZZ$ production in supersymmetry with R-parity violation

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We consider new s-channel scalar exchanges in $e^+e^-	o ZZ$, $W^+W^-$ in supersymmetry with a small lepton number violation. We show that a small bilinear R-parity violating term which leads to sneutrino–Higgs mixing can give rise to a significant scalar resonance enhancement in $e^+e^-	o ZZ$, $W^+W^-$. We use the LEP2 measurements of the $WW$ and $ZZ$ cross-sections to place useful constraints on this scenario. We also find, under conservative assumptions on the relevant parameter space involved, that such an exchange of the sneutrino-like admixture in $e^+e^-	o ZZ$, $W^+W^-$ may be accessible to a 500 GeV $e^+e^-$ collider.

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In spite of the indisputable success of the Standard Model (SM) in confronting experimental data, there are strong theoretical motivations for new beyond the SM physics being just around the corner. One of the most attractive new physics scenarios is supersymmetry (SUSY), which offers a plethora of new phenomena that might be observed in upcoming future colliders. As opposed to the SM, lepton number does not have to be conserved in the SUSY new physics scenarios is supersymmetry (SUSY), which offers a plethora of new phenomena that might be observed in upcoming future colliders. As opposed to the SM, lepton number does not have to be conserved in the SUSY Lagrangian. In fact, there is no fundamental principle that enforces lepton number conservation.

The SUSY superpotential can violate lepton number (or more generally R-parity) via an R-parity violating (RPV) Yukawa-like trilinear term (RPVTT) in the purely leptonic sector, and via a mass-like RPV bilinear term (RPVBT) as follows

\[ W_{R PV} \supset \epsilon_{ab} \left[ \lambda_{ijk} \hat{L}^a_i \hat{E}^b_j / 2 - \mu_a \hat{L}^a_i \hat{H}^b_u \right] , \]  

where $i, j, k = 1, 2$ or 3 label the lepton generation and $a, b = 1, 2$ are SU(2) indices. In what follows we will assume that only $\mu_3 \neq 0$. The scalar potential contains the corresponding soft SUSY breaking RPV terms in addition to the usual R-parity conserving (RPC) ones. The relevant ones for our discussion are $b_3 \hat{L}_3 \hat{H}_u$ and $M_{\tilde{L}H}^2 \hat{L}_3 \hat{H}_d$ which lead to a non-vanishing VEV of the tau-sneutrino, $\nu_3$. However, since lepton number is not a conserved quantum number in this scenario, the $\tilde{H}_d$ and $\tilde{L}_3$ superfields loose their identity and can be rotated to a particular basis ($\tilde{H}'_d, \tilde{L}'_3$) in which either $\mu_3$ or $\nu_3$ is zero. In what follows, we find it convenient to choose the basis where $\nu_3 = 0$.

Throughout this paper we will assume a small lepton number violation in the SUSY Lagrangian. That is, $|\mu_3/\mu_0| \ll 1$, $|\lambda_{ijk}| \ll 1$ and $b_3/b_0 \ll 1$, where $\mu_0$ is the usual RPC Higgs mass term, $\mu_0 \tilde{H}_d \tilde{H}_u$, and $b_0 \tilde{H}_d \tilde{H}_u$ is the corresponding soft term. Note that in the $\nu_3 = 0$ basis the minimization of the scalar potential yields:

\[ b_3 \propto (M_{\tilde{L}H}^2 + \mu_3 \mu_0) \cot \beta, \]  

where $\tan \beta \equiv v_u/v_d$. Thus, in the general case, $b_3$ needs not vanish even if $\mu_3$ is vanishingly small, as may be suggested by low energy flavor changing processes (see e.g., $\nu_3$) and/or flavor changing $Z$-decays (see e.g., Bisset $et$ $al.$, in [3]). That is, if $M_{\tilde{L}H}^2 \gg \mu_3 \mu_0$ due to $\mu_3 \to 0$, then $b_3 \sim M_{\tilde{L}H}^2 \cot \beta$, in which case RPV in the scalar sector decouples from the RPV in the superpotential (i.e., $\mu_3$). Thus, small lepton number violation in the scalar potential can be realized by requiring only that $b_3 \ll b_0$.

Some of the interesting phenomenological implications of the RPVBT are tree-level neutrino masses and new scalar decay channels. In this letter we suggest yet a new signature that can serve as an exclusive probe of the RPVBT. In particular, one can have scalar resonances in massive gauge-boson pair production:

\[ e^+e^- \to \Phi_k \to VV \]  

where $\Phi_k$ are admixtures of the RPC CP-even neutral Higgs and tau-sneutrino fields as described below. Note that the CP-odd scalar states do not couple to $VV$ at tree-level. Such a resonance can arise with measurable consequences when the incoming $e^+e^-$ beam couples to the CP-odd component in $\Phi_k$ with a coupling $\propto \lambda \gg m_e/M_W$ in [4],

1The RPVTT $\lambda' \hat{L} \hat{Q} \hat{D}$ is not relevant for our discussion.
2The effects of $\mu_1 \neq 0$ and/or $\mu_2 \neq 0$ are not crucial for the main outcome of this paper.
3Note that the laboratory limit on the $\tau$-neutrino mass allows $b_3/b_0 \sim \mathcal{O}(1)$ [3].
while the VV final state couples to the Higgs components in $\Phi_k$. Therefore, this scalar exchange can be attributed only to the Higgs– sneutrino mixing phenomena via the RPVBT and is a viable mechanism for probing a RPVBT beyond a RPVTT, i.e., $\sigma(e^+e^\to \Phi_k \to VV)$ $\to 0$ as $b_3 \to 0$. It should be stressed that this resonance formation is different from previously suggested sneutrino resonances within RPV SUSY such as fermion pair production in leptonic colliders \[9,11\], since it is driven by RPV parameters in the soft breaking scalar sector and not purely by Yukawa-like RPV couplings in the superpotential.

Let us define the SU(2) components of the neutral Higgs and stau fields, respectively, as: $H_{d.u}^0 = (\phi_{d,u}^0 + \nu_{d,u}^0)/\sqrt{2}$ and $\tilde{\nu}_t \equiv (\nu_{d,u}^0 + v_\nu + \nu_3^0)/\sqrt{2}$, then setting $v_3 = 0$. The CP-even $3 \times 3$ symmetric scalar mass matrix is then obtained through the quadratic part of the scalar potential as: $\frac{1}{2} \Phi^0 M^2_{\Phi} (\Phi^0)^T$, where $\Phi^0 = (\xi^0_d, \xi^0_s, \nu^0_3)$.

In the RPC limit the Higgs and sneutrino sectors couple, i.e., $M^2_\Phi$ consists of the usual $2 \times 2$ upper left block corresponding to the two CP-even Higgs states (see e.g., \[9\]) and $(M^2_\Phi)_{33} = m^2_{\nu_3}$. $(M^2_\Phi)_{13} = b_3 \tan \beta$ and $(M^2_\Phi)_{23} = (M^2_\Phi)_{32} = -b_3$, which are responsible for the mixing of $\xi^0_3$ with $\nu_3^0$. As a result, the usual CP-even RPC Higgs states $H^0$ and $h^0$ $(m_H > m_W)$ acquire a small $\nu_3^0$ component and vice versa \[10\].

The new CP-even scalar mass-eigenstates (i.e., the physical states) will be denoted here after by $\Phi \equiv (H, h, \tilde{\nu}_t)$, where, for small RPV in the SUSY Lagrangian, $H, h$ and $\tilde{\nu}_t$ are the states dominated by $H^0$, $h^0$ and $\nu_3^0$, respectively. They are related to the weak eigenstates via $\Phi^0_j = S_{3k} \Phi_k$, where $S$ is the rotation matrix that diagonalizes $M^2_\Phi$ (i.e., $S^T M^2_{\Phi} S = \text{diag}(m_{H}^2, m_{h}^2, m_{\nu_3}^2)$) throughout the rest of the paper we use $m_{H}, m_{h}, m_{\nu_3}$ and $m^0_{\Phi}, m^0_{H}, m^0_{h}, m^0_{\nu_3}$ to denote the masses of the $H, h, \tilde{\nu}_t$ physical states and of the $H^0, h^0, \nu_3^0$ states, respectively). The interaction vertices of the physical states $\Phi$ are then obtained by rotating the Feynman rules of the RPC SUSY Lagrangian (see e.g., \[10\]) with the matrix $S$. Thus, if $\Lambda_{\Phi^0}$ is an interaction vertex involving a weak state, then $\Lambda_{\Phi_k}$, the vertex involving the physical state, is given by $\Lambda_{\Phi_k} = S_{3k} \Lambda_{\Phi^0}$.

Hence, the $\Phi_kVV$ coupling is given by:

$$\Lambda_{\Phi_kVV} = i(e/s_W)C_V m_V (c_\beta S_{1k} + s_\beta S_{2k}) g_{\mu\nu},$$

where $C_V = 1/(s_W^2)$ for $V = W(Z)$, $s_W, c_W = \sin \theta_W, \cos \theta_W$ and $c_\beta, s_\beta = \cos \beta, \sin \beta$. Note that for $b_3 \to 0$, $S_{11} = S_{22} = \cos \alpha, S_{12} = -S_{21} = \sin \alpha, S_{33} = 1$ and $S_{13,23,31,32} \to 0$, where $\alpha$ is the usual mixing angle of the RPC neutral CP-even Higgs sector \[9,10\].

The couplings $\Lambda_{\Phi_{e\nu}e\nu}$ are obtained from the RPVTT term in \[9\] and are $\Lambda_{\Phi_{e\nu}e\nu} = S_{3k} \Lambda_{\nu_3^0 e\nu} = i S_{3k} \lambda_{131}/\sqrt{2}$. In our numerical analysis we will set $\lambda_{131} = 0.1$ irrespective of $m_{\nu_3}$. We note, though, that the present $1\sigma$ limit, $\lambda_{131} \lesssim 0.06 \times m_{\tilde{e}_R}/[100 \text{ GeV}]$, does not rule out $\lambda_{131} \sim 0.3$ if the typical slepton mass is $m_{\tilde{e}_R} \sim m_{\nu_3} \sim 500 \text{ GeV}$; since $\lambda_{131} \to 0$ for $e^+e^- \to \Phi \to VV$ $\propto \lambda^2_{131}$ (see below), it can be easily rescaled for different values of $\lambda_{131}$.

$$\sigma^0_V \equiv \sigma(e^+e^- \to \Phi \to VV)$$

is thus given by:

$$\sigma^0_V = \frac{\beta_V (3 - 2\beta_V^2 + 3\beta_V^4)}{128 s^2_W} \frac{\alpha}{s(1 - \beta_V^2)} \lambda^2_{131} \times \sum_{i,j=1}^{3} S_{3i} S_{3j} A_i A_j \tilde{\Pi}_i \tilde{\Pi}_j,$$

where $\beta_V = 2(1)$ for $V = W(Z), \beta_V = \sqrt{1 - 4 M^2_V/s}$ (s is the square of the c.m. energy) and $A_k = (c_\beta S_{1k} + s_\beta S_{2k})$.

Also, $\tilde{\Pi}_k = (1 - x_k^2 + ix_k y_k)^{-1}$, with $x_k = m_{\Phi_k}/\sqrt{s}, y_k = \Gamma_{\Phi_k}/\sqrt{s}$ and $\Gamma_{\Phi_k}$ is the $\Phi_k$ width. The interferences between the SM diagrams and our $s$-channel scalar diagrams are $\propto m_e$ and therefore negligible. Thus, the total cross-section for $e^+e^- \to VV$ is simply the sum $\sigma^0_V = \sigma^0_{\Phi^0} + \sigma^0_V$.

Let us now establish our relevant low-energy SUSY parameter space. The usual RPC CP-even Higgs sector can be described at tree-level by only two parameters $\xi^0_3, \lambda_{131}$, conventionally chosen to be $m^0_\Phi$ - the pseudo-scalar Higgs mass in the RPC case - and $t_{3\beta} = \tan \beta$. Furthermore, with the assumption of small RPV, i.e., RPV/RPV $\ll 1$, $m^0_\Phi$ typically scales as $(m^0_{\Phi})^2 \sim b_3 t_\beta$ for $t_\beta^2 \gg 1$. Thus, without loss of generality we set $b_3 \equiv \varepsilon (m^0_A)^2 \cot \beta$, such that small lepton number violation in the scalar sector is parameterized by the dimensionless quantity $\varepsilon \sim b_3/b_0$. Then $\varepsilon \ll 1$ corresponds to $b_3 \ll b_0$. The parameter set $\{m^0_A, m^0_{\nu_3}, t_\beta, \varepsilon\}$ therefore completely fixes $M^2_\Phi$ at tree-level from which the rotation matrix $S$ and the tree-level masses $m_{\Phi_k}$ are derived.

We note that since $b_3 = \varepsilon (m^0_A)^2 \cot \beta$, when $\varepsilon \ll 1$ implying $b_3 \ll b_0$ and also $b_3 \ll (m^0_A)^2$, the masses of the physical CP-even states $m_H, m_h, m_{\nu_3}$ and of the CP-odd states (e.g., $m_A$) are only slightly shifted from the

\[4\] We use the superscript 0 to denote the particle states in the RPC limit.
corresponding states in the RPC limit \((m_H^0, m_{H^0}, m_{\nu^0}, \text{ and } m_A^0)\) as long as there is no accidental mass degeneracy among the scalar states \([12]\). In particular, the scalar masses will be shifted by terms proportional to \(b_k^2/([m_{\phi_0^k}]^2 - (m_{\psi_0^k})^2)\) with \(k \neq \ell\) (see e.g., \([2]\)). Thus, although we are using “bare” masses (i.e., the scalar masses in the RPC limit) as inputs, it should be kept in mind that the physical masses are only slightly shifted. For example, for \(\varepsilon = 0.1\) and \(|m_{\nu^0} - m_A^0|, |m_{H^0} - m_A^0| \sim 100\) GeV we find that the shift in \(m_{A_1}, |m_A - m_A^0|\), is at the level of a few percent at the most for both a low or a high tan \(\beta\) scenario (for more details see \([12]\)).

Moreover, the fact that the mass shifts due to \(\varepsilon \neq 0\) are proportional to the sign of \((m_{\phi_0^k} - m_{\psi_0^k})\) has important consequences on the light CP-even Higgs particle. In particular, we find that if \(m_A^0, m_{\nu^0} > m_h^0\) (as always chosen below), then \(m_h\) tends to decrease with \(\varepsilon\). We can thus use the present LEP2 limit on \(m_h\) to deduce the allowed range in e.g., the \(\varepsilon - m_{\nu^0}\) plane, for a given \(m_A^0\). In particular, the present LEP2 bound is roughly \(m_h \gtrsim 110\) GeV, for \(m_A \gtrsim 200\) GeV irrespective of \(t_\beta\) and in the maximal mixing scenario with a typical SUSY scale/squark mass of 1 TeV \([14,15]\).

Therefore, hereafter, we include the dominant higher order corrections (coming from the \(t - t\) sector) to the \((\tilde{\xi}_1^\nu, \tilde{\xi}_2^\nu)\) block in \(M_Z^2\), using the approximated formulae given in \([14]\) with the maximal mixing scenario (as defined in \([4]\) and setting the typical squark mass at \(m_{\tilde{q}} \sim 1\) TeV). For example, in Fig. \([1]\) we show the excluded region in the \(\varepsilon - m_{\nu^0}\) plane (the shaded area) from the recent LEP2 limit of \(m_h \gtrsim 110\) GeV which holds for the parameters set \(\tan \beta = 3\) and \(m_A^0 = 300\) or 600 GeV as used in Fig. \([1]\).

In what follows, we focus on the case of a heavy Higgs spectrum, in particular \((m_{\phi_0^k})^2 \gtrsim M_Z^2\), which leads to the near mass degeneracy \(m_{\phi_0^k} \sim m_{\phi_0^k}^0\) and equivalently \(m_H \sim m_A\). We find that with \((m_{\phi_0^k})^2 \gtrsim M_Z^2\), the \(\tilde{\nu}_+\) sneutrino-like state will potentially yield the dominant signal. This can be understood as follows: (i) A light Higgs \((h)\) resonance in on-shell \(VV\) pair production is theoretically excluded, since the c.m. energy required to produce an on-shell \(VV\) pair is at least \(\sim 25\) GeV above the highest possible \(m_h\) (the theoretical upper limit on \(m_h\) is \(\sim 135\) GeV). Therefore, the \(h\) contribution to \(\sigma_V^0\) is always negligible and in particular near a \(\tilde{\nu}_+\) resonance.\(^{\text{ii}}\) (ii) For \((m_{\phi_0^k})^2 \gtrsim M_Z^2\) and \(\varepsilon \ll 1\), a heavy Higgs \((H)\) resonance in \(\sigma_V^0\) will also be much smaller than a \(\tilde{\nu}_+\) resonance since

\[
S_{11} \, \varepsilon \to 0 \quad \text{cos} \, \frac{m_{\nu^0}^2 \sim M_Z^2}{\sin \beta}, \quad S_{21} \, \varepsilon \to 0 \quad \text{sin} \, \frac{m_{\nu^0}^2 \sim M_Z^2}{- \cos \beta},
\]

leading to \(A_1 \to 0\), where \(A_1\) is the reduced \(HVV\) coupling defined in \([1]\). In addition, for \(\varepsilon \to 0\) the element connecting \(\nu^0\) to \(H\) coupling to the incoming electron since it couples to \(e^+e^-\) through its dominant \(\nu^0\) component. In particular, \(S_{31} \to 1\) as \(\varepsilon \to 0\). Thus, even though \(\lambda_{\nu}, VV\) and \(\lambda_{HVV}\) are comparable (or equivalently \(A_1 \sim A_1\))\(^{\text{ii}}\) the \(\tilde{\nu}_+\) exchange contribution to \(\sigma_V^0\), being \(\propto S_{31} \times A_1\), will be more pronounced than the \(H\) exchange one since \(S_{31} \gg S_{31}\) for \(\varepsilon \ll 1\).

Therefore, the more favorite scenario for observing such a sneutrino-Higgs mixing resonance in \(VV\) pair production is when the \(\tilde{\nu}_+\) resonates. It should be noted, however, that \(\sigma_V^0\) may be further enhanced considerably if both \(m_{\nu^0}\) and \(m_H\) happen to lie close to the c.m. energy in the given experiment. As mentioned before, we do not consider in this paper such a possibility of an accidental mass degeneracy between the \(H\) and \(\tilde{\nu}_+\) states which may give rise to a “combined” \(H + \tilde{\nu}_+\) resonance. Hence, in what follows we will consider only the case of a sneutrino-like resonance in \(e^+e^-\) to \(VV\).

The \(\tilde{\nu}_+\) width, \(\Gamma_{\tilde{\nu}_+}\) in \([1]\), needs to be included, since it controls the behavior of \(\sigma_V^0\) in the vicinity of our \(\tilde{\nu}_+\) resonance. Assuming that the lightest neutralino \((\tilde{\chi}^0_1)\) is the Lightest SUSY Particle (LSP) and also that \(m_{\nu^0} > m_{\tilde{\chi}^+_1}\), where \(\tilde{\chi}^+_1\) is the lighter chargeino, then the RPC two-body decays \(\tilde{\nu}_+ \to \tilde{\chi}^0_1 \nu_\tau, \tilde{\chi}^+_1 \tau\) are open and dominate. Indeed, for \(m_{\nu^0}^2 \gg M_Z^2\) and following the traditional assumption of an underlying grand unification with a common gaugino mass parameter of \(m_{1/2} < m_{\nu^0}\), the mass hierarchy \(m_{\phi_0^0} < m_{\phi_0^0} \sim m_{\chi^+_1} < m_{\nu^0}\) and \(m_{\phi_0^0} \sim m_{\chi^+_1} > m_{\nu^0}\) is possible, e.g., when \(m_{\nu^0} < m_A\).\(^{\text{ii}}\) Thus, upon ignoring phase space factors, a viable conservative estimate is (see e.g., \([14-16]\)); \(T_{\tilde{\nu}_+} \sim \Gamma(\tilde{\nu}_+ \to \tilde{\chi}^0_{1,2} \nu_\tau) + \Gamma(\tilde{\nu}_+ \to \tilde{\chi}^+_1 \tau) \sim 10^{-2} m_{\nu^0}\), which we use below. Note that for the

\(^{\text{v}}\)Since \(b_3 \neq 0\) the \(hZZ\) coupling \([A_2\) in \([1]\)] is smaller than its value in the RPC case leading to a smaller \(e^+e^-\) to \(Zh\) production rate. The limits on \(m_h\) given in \([14,15]\) are therefore slightly weaker in the RPV case (see e.g., \([3]\)).

\(^{\text{vi}}\)Note, however, that in this scenario an \(s\)-channel \(h\) exchange may lead to similar resonant enhancement in \(e^+e^-\) to \(VV^*\) at lower c.m. energies, where \(V^*\) is an off-shell \(W\) or \(Z\).

\(^{\text{vii}}\)Since \((M_Z^2)_{1/2} = t_\beta, \tilde{\chi}^0_0\) acquires a larger \(\xi_0^0\) mixing (than a \(\xi_0^0\) mixing) which in turn implies a larger \(H\) mixing, since the \(H\) mass-eigenstate is mostly \(\xi_0^0\) weak-state when \((m_\nu^0)^2 \gg m_Z^2\) and \(t_\beta^2 \gg 1\).
FIG. 1. The shaded area in the $\varepsilon - m_{0}^{s\nu}$ plane $\varepsilon \equiv b_{3}t_{\beta}/(m_{A}^{0})^{2}$ is excluded by the recent LEP2 limit on the light Higgs mass $m_{h} \gtrsim 110$ GeV. This excluded region is independent of $\lambda_{131}$.

ranges of $\varepsilon$, $m_{A}^{0}$ and $m_{s\nu}^{0}$ considered the possible RPV decays are sufficiently smaller and $\Gamma_{\tilde{\nu}_{+}} \sim \Gamma_{\tilde{\nu}_{2}}$ since $S_{33} \rightarrow 1$. Also, for reasons explained above, $\Gamma_{H}$ and $\Gamma_{h}$ have a negligible effect on the $\tilde{\nu}_{+}$ resonance and are therefore neglected.

Before presenting our numerical results we note the following: (i) Sufficiently away from threshold ($\beta_{V} \rightarrow 1$), $\sigma_{W}^{0}/\sigma_{Z}^{0} \sim (\delta_{W}c_{W}^{2}/\delta_{Z}) \cdot (M_{Z}/M_{W})^{2} \sim 2$ and, since typically $\sigma_{W}^{SM}/\sigma_{Z}^{SM} > 10$, the relative effect of the scalar exchange cross-section is more pronounced in the ZZ channel. (ii) As mentioned above, for $\varepsilon \ll 1$ and in the decoupling limit [i.e., $(m_{A}^{0})^{2} \gg m_{Z}^{2}$], $\Lambda_{3\nu V} \rightarrow 1$ and $\Lambda_{HVV} \rightarrow 0$. At the same time, when $t_{\beta}^{2} \gg 1$, $\xi_{u}^{0} \rightarrow H$ and $\xi_{d}^{0} \rightarrow h$ so that, accordingly, for $t_{\beta}^{2} \gg 1$, $\left(\Lambda_{3\nu V}^{0}/\Lambda_{3\nu V}^{0}\right) \gg 1$. Thus, since $(M_{Z}^{2})_{23} = b_{3} = \varepsilon (m_{A}^{0})^{2}/t_{\beta}$, the $\tilde{\nu}_{u}^{0} - \xi_{u}^{0}$ mixing decreases with $\tan \beta$ and so as $t_{\beta}$ increases the sneutrino “prefers” to mix more with $\xi_{d}^{0}$ which has a suppressed coupling to $VV$ in this limit. As a consequence, the sneutrino-like resonance effect in $\sigma_{V}^{0}$ drops with $\tan \beta$ in the limit of small RPV and $m_{A}^{2} \gg M_{Z}^{2}$.

In Fig. 2 we show $\sigma_{Z}^{0}$ as a function of $m_{s\nu}^{0}$ for c.m. energies of $\sqrt{s} = 200$ and 500 GeV. This is shown for $t_{\beta} = 3$ and for $m_{A}^{0} = 300$ GeV (left side) or $m_{A}^{0} = 600$ GeV (right side) (a more detailed investigation of the parameter
space involved will be given in [12]). For definiteness we take \( \varepsilon = 0.05 \), \( 0.1 \) and \( \lambda_{131} = 0.1 \). The SM cross-sections \( \sigma_Z^{SM}(\sqrt{s} = 200 \text{ GeV}) \sim 1.29 \text{ [pb]} \) and \( \sigma_Z^{SM}(\sqrt{s} = 500 \text{ GeV}) \sim 0.41 \text{ [pb]} \) are also shown by the horizontal solid lines.

The SM cross-sections \( \sigma^{SM}_{Z} \) are also shown by the horizontal solid lines.

![Graph](image_url)

**Fig. 2.** \( \sigma^{0}_{Z} \) as a function of \( m^{0}_{\tilde{\nu}} \), for \( m^{0}_{A} = 300 \text{ GeV} \) (left plot) and \( m^{0}_{A} = 600 \text{ GeV} \) (right plot). For both values of \( m^{0}_{A} \), \( \sigma^{0}_{Z} \) is shown for the c.m. energies \( \sqrt{s} = 200 \text{ GeV} \) with \( \varepsilon = 0.1 \), \( 0.05 \) (left curves) and \( \sqrt{s} = 500 \text{ GeV} \) with \( \varepsilon = 0.1 \), \( 0.05 \) (right curves). \( \lambda_{131} = 0.1 \) is used (note that \( \sigma^{0}_{Z} \) scales as \( \lambda_{131}^{2} \)). The SM ZZ cross-sections for \( \sqrt{s} = 200 \) and 500 GeV are also shown by the horizontal solid lines.

We see that, as expected, \( \sigma^{0}_{Z} \) is larger for a smaller \( |m^{0}_{A} - m^{0}_{\tilde{\nu}}| \) mass splitting, since the sneutrino–Higgs mixing phenomena is proportional to factors of \( [(m^{0}_{A})^{2} - (m^{0}_{\tilde{\nu}})^{2}]^{-1} \) (see discussion above). Clearly, the scalar exchange cross-section can be statistically significant even if the mass of the sneutrino-like scalar is away from the resonance, i.e., within a range of \( m_{\tilde{\nu}} - \sqrt{s} \leq \Delta \), where, as we shall see below, \( \Delta \) may range from a few GeV to a few tens of GeV depending on \( \varepsilon \) and the rest of the SUSY parameter space involved.

Thus, for the case of \( \sqrt{s} \) around 200 GeV, we can use the measured values of the \( WW \) and \( ZZ \) cross-sections at LEP2 to place further bounds on the \( \varepsilon - m^{0}_{\tilde{\nu}} \) plane for a given \( m^{0}_{A} \) and \( t_{\beta} \). This is shown in Fig. 3 where we have

\( \sigma^{0}_{V} \) is insensitive to the signs of \( \varepsilon \) and \( \lambda_{131} \).
again set $m_A^0 = 300$ GeV or $m_A^0 = 600$ GeV, $t_\beta = 3$, $\lambda_{131} = 0.1$ and used the measured $\sigma_Z$ and $\sigma_W$, combined by the 4 LEP experiments, from the 183, 189, 192, 196, 200, 202, 205 and 207 GeV LEP2 runs as given in [16]. In particular, for each run we take the measured and the SM cross-sections (also given in [16]), $\sigma_V^{\text{exp}}; \sigma_V^{\text{SM}}$, and require that $\sigma_V^0 < (\sigma_V^{\text{exp}} - \sigma_V^{\text{SM}}) + \sqrt{\Delta\sigma_V^{\text{exp}}^2 + \Delta\sigma_V^{\text{SM}}^2}$. 

\[ \text{FIG. 3. 1\sigma excluded regions in the } \varepsilon - m_{s\nu}^0 \text{ plane from the LEP2 measurements of the } \text{WW and ZZ cross-sections (see text).} \]

Evidently, the limits coming from the ZZ and WW cross-sections measurements give further restrictions at low $\varepsilon$ values below $\sim 0.2$ (in a sneutrino mass range of several tens of GeV [11], for which there are no bounds coming from the LEP2 limits on $m_h$ (see Fig. 1). Note that the fingers like shape of the shaded area in Fig. 3 is an artifact of the fact that we are using a discrete set of c.m. energies in accordance with the LEP2 runs.

9For the ZZ and WW SM cross-sections we use the results of the ZZTO and YFSWW3 Monte-Carlos, respectively, where we take a 2% theoretical error for the ZZTO prediction and no error for the YFSWW3 one, see [14].

10We do not include the cases in which $(\sigma_V^{\text{exp}} - \sigma_V^{\text{SM}}) + \sqrt{(\Delta\sigma_V^{\text{exp}})^2 + (\Delta\sigma_V^{\text{SM}})^2} < 0$.

11Note that, since $b_3 = \varepsilon (m_A^0)^2 / t_\beta$, these 1\sigma limits can be directly translated into limits on the $b_3 - m_{\nu_s}^0$ plane.
Alternatively, for the case of a 500 GeV $e^+e^-$ collider we can find the mass range of the sneutrino-like scalar for which its contribution to the WW and ZZ cross-sections may be observable with a statistical significance of at least $3\sigma$ by requiring $\left(\sigma_{V}^{WW}/\sqrt{\sigma_{V}^{ZZ}+\sigma_{SM}^{0}}\right) > 3$. For example, we find that with an integrated luminosity of $L = 100$ fb$^{-1}$, $m_{A}^0 = 600$ GeV and $\beta_t = 3$, a more than $3\sigma$ signal can arise in the ZZ case within the sneutrino mass ranges $490 \text{ GeV} \lesssim m_{s\nu} \lesssim 509 \text{ GeV}$ and $495 \text{ GeV} \lesssim m_{s\nu} \lesssim 505 \text{ GeV}$ for $\varepsilon \sim 0.1$ and 0.05, respectively. The corresponding $3\sigma$ mass intervals in the WW case are typically a factor $\sim 1.5$ smaller for the same values of $\beta_t$, $m_{A}^0$ and $\varepsilon$. These $3\sigma$ mass ranges are further enlarged if an angular cut on the c.m. scattering angle, $\theta$, is imposed. For example, with $0 \lesssim \cos \theta \lesssim 1$, we find that, for $m_{A}^0 = 600$ GeV, $\beta_t = 3$ and $\varepsilon \sim 0.1$ or 0.05, the sneutrino resonance will be observable at $3\sigma$ in the WW channel within the mass ranges $489 \text{ GeV} \lesssim m_{s\nu} \lesssim 511 \text{ GeV}$ or $495 \text{ GeV} \lesssim m_{s\nu} \lesssim 505 \text{ GeV}$, respectively. These mass ranges are comparable to the ones obtained in the ZZ case with no angular cut.

To summarize, a small lepton number violation scenario, which incorporates small trilinear and bilinear RPV terms into the SUSY Lagrangian, can lead to a significant scalar resonance enhancement in $e^+e^- \rightarrow ZZ, W^+W^-$ due to mixings between the sneutrino and the Higgs particles, which may be accessible to a 500 GeV $e^+e^-$ collider. We also find that useful limits can be placed on this scenario from the LEP2 measurements of the WW and ZZ cross-sections and from the LEP2 limits on the light Higgs mass. Finally, we note that a similar scalar resonance may arise in top-quark pair production due to the sneutrino--Higgs mixing phenomena (see [12]). Such a resonance enhancement in $e^+e^- \rightarrow t\bar{t}$ should give further evidence in favor of the bilinear RPV SUSY scenario since the absence of a tree-level sneutrino–top–anti-top trilinear RPV coupling and the fact that the Higgs-electron-positron coupling is $\propto m_t$, makes the sneutrino–Higgs mixing the only viable mechanism for generating an observable resonance signal in $e^+e^- \rightarrow t\bar{t}$ within the SUSY framework.

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