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Four-Jet Signal at LEP2 and Supersymmetry

M. Carena^{a,b}, G.F. Giudice^{a1}, S. Lola^a and C.E.M. Wagner^a

^a Theory Division, CERN
Geneva, Switzerland

^b Deutsches Elektronen-Synchrotron, DESY
Hamburg, Germany

Abstract

ALEPH has reported a significant excess of four-jet events in the LEP runs above the Z^0 resonance, which however has not been confirmed by the other LEP collaborations. We assume here that this excess corresponds to a physics signal and try to interpret it in the context of supersymmetric models with R -parity violation. Associated production of a left and right selectron can explain all the distinctive features of the ALEPH data: the value of the cross section, the dijet mass difference, the absence of bottom quarks in the final state, and the dijet charge content. Our proposed scenario makes definite predictions, which can be tested at future LEP runs at higher energies.

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¹On leave of absence from INFN, Sezione di Padova, Padua, Italy.

1 Introduction

One of the most intriguing and controversial results of the LEP run above the Z^0 resonance has been the excess of four-jet events reported by ALEPH [1, 2]. As the other three experimental collaborations working at LEP do not observe any anomaly in four-jet topologies, the resolution of the experimental controversy is a most urgent issue. All experimental collaborations are actively working on the question, which hopefully will be settled by further study and, most importantly, by the new runs at higher energies. From the theoretical side, we believe that it is important to investigate if the reported ALEPH data can be interpreted as a consistent physics signal. At least such a study can be used as a benchmark to compare the present results with future data at higher \sqrt{s} .

The ALEPH four-jet events have been selected from the data recorded at centre-of-mass energies between 130 and 172 GeV. An excess is observed in the distribution of the sum of the two dijet invariant masses constructed by pairing jets with the smallest dijet mass difference. This distribution shows a peak at 106.1 ± 0.8 GeV, corresponding to 18 events observed with 3.1 expected from QCD background [2]. If interpreted as particle pair production, this corresponds to a cross section of 2.5 ± 0.7 pb when only data with \sqrt{s} in the range between 130 and 161 GeV are considered, and of 1.5 ± 0.8 pb when all data with \sqrt{s} in the range between 130 and 172 GeV are considered [2]. This cross section is too large for Higgs bosons or for electroweakly-interacting scalar particles, whose productions are suppressed by a factor β^3 . Here β is the final-state particle velocity in the centre-of-mass, which is, for the relevant kinematical configuration, about 0.6 at $\sqrt{s} = 130$ GeV and 0.8 at $\sqrt{s} = 172$ GeV. The inferred value of the cross section could be accommodated by production of fermions with electroweak couplings or of scalar particles with a substantial colour or multiplicity factor.

The dijet mass difference distribution of the selected 18 events is consistent with a value around 10 GeV [1, 2]. Combining this with the information on the dijet mass sum, it can be concluded that the pair-produced particles should have masses of about 58 and 48 GeV, respectively. Pair-production of equal-mass particles is disfavoured.

At the moment little information can be extracted from angular distributions. From measurements of the “rapidity-weighted” jet charge, a variable that statistically retains information on the electric charge of the primary parton [3], one concludes [1] that the pair-produced particles have a sizeable charge. Electrically neutral particles are therefore disfavoured. Finally there is little or no presence of b quarks in the final states [1, 2]. This is another reason to reject the hypothesis of Higgs-boson production.

In this paper we want to study whether the ALEPH data, assumed here to correspond to a real physics signal, can be explained by pair production of a left-handed and a right-handed selectron, each particle decaying into two quarks, as an effect of R -parity violating interactions. Other interpretations of the four-jet events have already been presented in the literature [4], but to our knowledge this is the first example of a consistent picture of pair production and decay of particles with different masses, in agreement with the results of the ALEPH analysis.

2 Particle Production

Let us start by considering the particle-production cross sections. Left-handed or right-handed selectrons can be pair-produced at LEP in any of the channels $\tilde{e}_L\tilde{e}_L$, $\tilde{e}_R\tilde{e}_R$, and $\tilde{e}_L\tilde{e}_R$. The interactions involved in the production of the pairs $\tilde{e}_L\tilde{e}_L$ or $\tilde{e}_R\tilde{e}_R$ always require that the incoming electron and positron have opposite helicities (*i.e.* collinear spin vectors). This means that the cross section for scalar particle production has a β^3 suppression, corresponding to a p -wave suppression near threshold. On the other hand, the production of $\tilde{e}_L\tilde{e}_R$ pairs involves the interaction of an electron and a positron with the same helicity (*i.e.* opposite spin vectors), and the cross section near threshold is proportional to β , corresponding to an s -wave.

The differential cross section for $\tilde{e}_L\tilde{e}_R$ production is

$$\frac{d\sigma}{dt}(e^+e^- \rightarrow \tilde{e}_L^+\tilde{e}_R^-) = \frac{d\sigma}{dt}(e^+e^- \rightarrow \tilde{e}_R^+\tilde{e}_L^-) = \frac{g^4 \tan^4 \theta_W}{64\pi s} \sum_{a,b=1}^4 \frac{A_a A_b M_{\chi_a^0} M_{\chi_b^0}}{(t - M_{\chi_a^0}^2)(t - M_{\chi_b^0}^2)}, \quad (1)$$

$$A_a \equiv N_{a1} \left(N_{a1} + \frac{N_{a2}}{\tan \theta_W} \right). \quad (2)$$

Here N_{a1} and N_{a2} are the B -ino and W_3 -ino components of the a -th neutralino with mass $M_{\chi_a^0}$. In the limit of a purely B -ino state, only one neutralino contributes to the sum in eq. (1). Therefore the $U(1)$ gaugino mass M_1 is the most important parameter entering eq. (1). For simplicity, we will concentrate on the case in which the lightest neutralino is a pure B -ino. Because of the necessary helicity flip proportional to the gaugino mass, the cross section decreases only as M_1^{-2} , for large M_1 .

To reproduce the kinematical configuration suggested by the ALEPH data, we choose $m_{\tilde{e}_L} = 58$ GeV and $m_{\tilde{e}_R} = 48$ GeV. The electron sneutrino mass is then also determined by the weak $SU(2)$ relation

$$m_{\tilde{\nu}_e}^2 = m_{\tilde{e}_L}^2 + (1 - \sin^2 \theta_W) \cos 2\beta M_Z^2, \quad (3)$$

where $\tan \beta$ is the usual ratio of Higgs vacuum expectation values. We assume here that the sleptons of the second and third generations are heavier than those of the first one, and cannot have been produced at LEP. We will comment in sect. 5 on the case in which slepton masses are universal in flavour.

In fig. 1 we show the cross sections at $\sqrt{s} = 130$ and 172 GeV for the different production channels $\tilde{e}_L\tilde{e}_R$, $\tilde{e}_L\tilde{e}_L$, $\tilde{e}_R\tilde{e}_R$, $\tilde{\nu}_e\tilde{\nu}_e$, and $\chi_1^0\chi_1^0$ as a function of the $U(1)$ gaugino mass M_1 , in the limit of large μ and M_2 . The cross sections are corrected for initial-state radiation and have been generated by the numerical code SUSYXS [5]. The important result is that, for M_1 less than about 100 GeV, the $\tilde{e}_L\tilde{e}_R$ production cross section is large, and consistent with the value suggested by the ALEPH data. Also, in the range $M_1 = 80$ –100 GeV, all other particle–antiparticle production channels are quite suppressed. For the $\tilde{e}_L\tilde{e}_L$ and $\tilde{e}_R\tilde{e}_R$ channels, this is the result of an efficient destructive interference between the s -channel γ/Z exchange and the t -channel neutralino/chargino exchange. To sufficiently suppress the $\tilde{\nu}_e$ cross section, we have to choose $m_{\tilde{\nu}_e}$ close to the upper bound determined by eq. (3). This implies that $\tan \beta$ is close

to 1, and the top Yukawa coupling is quite large. This can be made consistent with the absence of a Landau pole below the grand unification scale only if some new physics threshold exists. Indeed, the need of an effective supersymmetry-breaking scale Λ_{SUSY} much lower than the Planck scale is also suggested, in our scenario, by the presence of small slepton masses together with larger gaugino masses. In fact, for very large values of Λ_{SUSY} , such hierarchy of masses would require a large amount of fine-tuning between the value of the slepton and the gaugino mass parameters at the high energy scale (for a recent discussion of the dependence of the renormalization group evolution of the scalar mass parameters on the effective supersymmetry breaking scale, see ref. [6]).

The results presented in fig. 1 correspond to the case in which the lightest neutralino is a pure B -ino. Had we assumed unification of gaugino masses, and values of the higgsino mass μ not too large, then the cross section for $\tilde{e}_L \tilde{e}_R$ production could be larger than what is shown in fig. 1, as a consequence of the mixing between B -ino and W_3 -ino, see eq. (1). However the $\tilde{\nu}_e$ production cross section would also sizeably increase, because of the constructive interference between chargino and Z exchange contributions. For instance, for large values of $|\mu|$, and $M_2 = 500$ (300) GeV, the charged slepton production cross sections at $\sqrt{s} = 172$ GeV are not significantly modified, but the sneutrino cross section is enhanced from 0.4 pb to 0.51 (0.72) pb for $m_{\tilde{\nu}_e} = 53$ GeV, while for $m_{\tilde{\nu}_e} = 58$ GeV the cross section is enhanced from 0.31 pb to 0.40 (0.57) pb. Hence, values of $M_2 \gtrsim 500$ GeV will efficiently suppress the $\tilde{\nu}_e$ production cross section.

The differential cross section for $\tilde{e}_L \tilde{e}_R$ production, eq. (1), leads to an angular distribution that is different from the usual scalar-particle pair production via gauge bosons in the s -channel with $d\sigma/dt \propto (ut - m^4)$. In fig. 2 we compare the two distributions as a function of the angle θ between the beam direction and one of the two dijet momenta ($0 < \theta < \pi/2$). At the moment, the experimental uncertainties are too large for us to distinguish between the two distributions. If the charge of the primary parton is identified, one can measure the forward-backward asymmetry \mathcal{A}_{FB} of the dijet system with a definite charge. In the case of ordinary scalar particle pair production, the distribution is symmetric in the forward and backward regions, and $\mathcal{A}_{FB} = 0$. This is however not true for the distribution in eq. (1), which produces the following integrated forward-backward asymmetry

$$\mathcal{A}_{FB} = \frac{\sqrt{[s - (m_{\tilde{e}_L} + m_{\tilde{e}_R})^2][s - (m_{\tilde{e}_L} - m_{\tilde{e}_R})^2]}}{s + 2M_1^2 - m_{\tilde{e}_L}^2 - m_{\tilde{e}_R}^2}. \quad (4)$$

Here we have assumed that the B -ino is an approximate mass eigenstate, and defined the forward and backward regions with respect to initial- and final-state particles with the same electric charge. For $M_1 = 80$ –100 GeV, \mathcal{A}_{FB} is large, about 40–60% at $\sqrt{s} = 130$ GeV and 30–50% at $\sqrt{s} = 172$ GeV.

3 Particle Decay

In order to generate the four-jet final state from the slepton pair, we have to introduce some R -parity violating interaction. The only renormalizable operator that couples quarks to leptons has the following expression in the superpotential:

$$\lambda_{ijk} L_L^i Q_L^j \bar{D}_R^k . \quad (5)$$

We have used here a standard notation for lepton and quark chiral superfields, and denoted the generation indices as i, j, k . We assume that one of the couplings λ_{ijk} is much larger than all the others; this coupling determines the decay mode of the lightest supersymmetric particle. As we do not want to consider top or bottom quarks in the final state, we are interested only in the couplings λ_{1jk} with $j, k = 1, 2$. If in the future more experimental information on the flavour content of the jets becomes available, we will be able to further restrict the choice of the operators.

Our interpretation of the four-jet events as slepton pairs requires that the R -parity violating decay mode has a branching ratio close to 1. Thus it is important to compare the rate for \tilde{e}_L decay into two quarks,

$$\Gamma(\tilde{e}_L^- \rightarrow \bar{u}_j d_k) = \frac{3\lambda_{1jk}^2}{16\pi} m_{\tilde{e}_L} , \quad (6)$$

with the R -parity conserving decay rates. Indeed, \tilde{e}_L can decay into the lightest supersymmetric particle, \tilde{e}_R , through neutralino exchange. In the approximation $M_{\chi^0} \gg m_{\tilde{e}_L}, m_{\tilde{e}_R}$, the decay widths are

$$\Gamma(\tilde{e}_L^- \rightarrow e^- e^+ \tilde{e}_R^-) = \frac{g^4 \tan^4 \theta_W}{3(8\pi)^3} m_{\tilde{e}_L}^3 F_1 \left(\frac{m_{\tilde{e}_R}^2}{m_{\tilde{e}_L}^2} \right) \sum_{a,b=1}^4 \frac{A_a A_b}{M_{\chi_a^0} M_{\chi_b^0}} , \quad (7)$$

$$F_1(x) = (1-x)(1+10x+x^2) + 6x(1+x) \ln x , \quad (8)$$

$$\Gamma(\tilde{e}_L^- \rightarrow e^- e^- \tilde{e}_R^+) = \frac{g^4 \tan^4 \theta_W}{3(16\pi)^3} m_{\tilde{e}_L}^5 F_2 \left(\frac{m_{\tilde{e}_R}^2}{m_{\tilde{e}_L}^2} \right) \sum_{a,b=1}^4 \frac{A_a A_b}{M_{\chi_a^0}^2 M_{\chi_b^0}^2} , \quad (9)$$

$$F_2(x) = (1-x)(1-7x-7x^2+x^3) - 12x^2 \ln x . \quad (10)$$

Also, \tilde{e}_L can decay into $\tilde{\nu}_e$ via W exchange

$$\Gamma(\tilde{e}_L^- \rightarrow f \bar{f}' \tilde{\nu}_e) = N_{ff'} \frac{G_F^2 m_{\tilde{e}_L}^5}{3(4\pi)^3} F_2 \left(\frac{m_{\tilde{\nu}_e}^2}{m_{\tilde{e}_L}^2} \right) , \quad (11)$$

Here $N_{ff'}$ is a colour factor, equal to 9, when summed over the light quarks and leptons in the final state. The decay rate for $\tilde{e}_L^- \rightarrow e^- \nu_e \tilde{\nu}_e$ can be neglected, as it is suppressed by the chargino mass.

Figure 3 shows the value of $BR(\tilde{e}_L^- \rightarrow \bar{u}_j d_k)$, as a result of a phase-space integration in the limit of a purely B -ino neutralino, but with no approximations on the ratio $M_{\chi_1^0}/m_{\tilde{e}_L}$. The R -parity violating mode dominates the \tilde{e}_L decays for λ_{1jk} larger than few times 10^{-4} . These values for the R -parity violating coupling constants are consistent with present bounds, as we discuss in the following.

Experimental bounds on λ_{1jk} depend on the values of the squark masses, which mediate the effective four-fermion interactions between the leptons and quarks. We give here the bounds for a typical squark mass of 300 GeV, although the value of the squark mass does not enter into our analysis. The heavier the squarks are, the weaker the bounds on λ_{1jk} become. From charged current universality, one finds $\lambda_{11k} < 0.1$ [7]. From limits on $BR(K^+ \rightarrow \pi^+ \nu \bar{\nu})$, one finds $\lambda_{1jk} < 0.03$ [8], although this limit depends on assumptions about the flavour structure. From radiative contributions to the electron neutrino mass, one can get significant limits only for R -parity violating operators that involve a third generation index [9]. From negative searches of neutrinoless double- β decay, one obtains an interesting limit on $\lambda_{111} < 8 \times 10^{-3}$ [10], and a bound on the product $\lambda_{121} \lambda_{112} < 3 \times 10^{-5}$ [11]. Experiments at HERA have set bounds on λ_{1jk} of about 10^{-1} for squark masses of 200 GeV; these bounds disappear for values of the squark masses above 300 GeV [12]. The only problematic constraint comes from cosmological considerations about the survival of a baryon asymmetry created at the very early stages of the Universe, which requires $\lambda_{1jk} < 10^{-7}$ [13]. However, this limit does not apply to cosmological models with low-temperature baryogenesis, and can also be evaded under certain conditions [14].

We therefore conclude that there is a large range of λ_{1jk} values, consistent with present bounds on R -parity violation, in which \tilde{e}_L decays almost entirely into a quark pair. This is true, although \tilde{e}_L is not the lightest supersymmetric particle, because the R -parity violating two-body decay is more important than kinematically suppressed three-body decay modes.

In our scenario, \tilde{e}_R , the lightest supersymmetric particle, does not participate in the R -parity violating interaction in eq. (5), which involves only quarks and left-handed leptons. Therefore the \tilde{e}_R decay will occur either via the small mixing ϕ between \tilde{e}_R and \tilde{e}_L ,

$$\Gamma(\tilde{e}_R^- \rightarrow \bar{u}_j d_k) = \frac{3}{16\pi} \lambda_{1jk}^2 \sin^2 \phi m_{\tilde{e}_R}, \quad (12)$$

or via virtual neutralino and \tilde{e}_L exchange,

$$\Gamma(\tilde{e}_R^- \rightarrow e^+ e^- \bar{u}_j d_k) = \frac{3g^4 \tan^4 \theta_W \lambda_{1jk}^2 m_{\tilde{e}_R}^3}{4(4\pi)^5} G_1 \left(\frac{m_{\tilde{e}_L}^2}{m_{\tilde{e}_R}^2} \right) \sum_{a,b=1}^4 \frac{A_a A_b}{M_{\chi_a^0}^2 M_{\chi_b^0}^2}, \quad (13)$$

$$G_1(x) = \frac{(x-1)}{6} (4x^2 + 25x + 1) \ln \left(\frac{x}{x-1} \right) - \frac{(12x^2 + 123x + 13)}{18} + x(3x+2) \text{Li} \left(\frac{1}{x} \right), \quad (14)$$

$$\Gamma(\tilde{e}_R^- \rightarrow e^- e^- u_j \bar{d}_k) = \frac{3g^4 \tan^4 \theta_W \lambda_{1jk}^2 m_{\tilde{e}_R}^5}{(8\pi)^5} G_2 \left(\frac{m_{\tilde{e}_L}^2}{m_{\tilde{e}_R}^2} \right) \sum_{a,b=1}^4 \frac{A_a A_b}{M_{\chi_a^0}^2 M_{\chi_b^0}^2}, \quad (15)$$

$$G_2(x) = \frac{(x-1)}{6}(5x^3 - 27x^2 - 15x + 1) \ln\left(\frac{x}{x-1}\right) - \frac{(60x^3 - 354x^2 - 460x + 43)}{72} - 6x^2 \text{Li}\left(\frac{1}{x}\right). \quad (16)$$

In fig. 4 we show the $BR(\tilde{e}_R \rightarrow \bar{u}_j d_k)$ as a function of the mixing angle ϕ , in the limit of a purely B -ino neutralino. Again, although eqs. (13)–(15) have been derived in the approximation $M_{\chi_a^0} \gg m_{\tilde{e}_L}, m_{\tilde{e}_R}$, the results plotted in fig. 4 follow from a numerical integration of phase space with no restrictive assumptions.

The mixing angle ϕ is related to the higgsino mass μ and to the trilinear coupling A by the relation

$$\sin \phi \simeq \frac{m_e(A - \mu \tan \beta)}{m_{\tilde{e}_L}^2 - m_{\tilde{e}_R}^2} = \left(\frac{A - \mu \tan \beta}{200 \text{ GeV}}\right) 10^{-4}. \quad (17)$$

Therefore the most plausible values for $\sin \phi$ lie in the range around 10^{-4} . From fig. 4 we then infer that, in this range, \tilde{e}_R dominantly decays into two jets. In conclusion, although \tilde{e}_R does not participate in the R -parity violating interaction, the small left–right mixing ensures that the preferred \tilde{e}_R decay mode is into a quark pair, rather than into a phase-space suppressed four-body final state.

The R -parity violating coupling λ_{1jk} does not influence the \tilde{e}_R decay branching ratio, as long as it is non-vanishing. It determines however the \tilde{e}_R lifetime, which is

$$\tau_{\tilde{e}_R} = \left(\frac{10^{-4}}{\sin \phi}\right)^2 \left(\frac{10^{-2}}{\lambda_{1jk}}\right)^2 2 \times 10^{-13} \text{ s}. \quad (18)$$

For the relevant kinematical configuration, this correspond to a decay vertex displacement of about

$$d_{\tilde{e}_R} = \sqrt{\frac{[s - (m_{\tilde{e}_L} + m_{\tilde{e}_R})^2][s - (m_{\tilde{e}_L} - m_{\tilde{e}_R})^2]}{4sm_{\tilde{e}_R}^2}} \tau_{\tilde{e}_R} \simeq \left(\frac{10^{-4}}{\sin \phi}\right)^2 \left(\frac{10^{-2}}{\lambda_{1jk}}\right)^2 50 \text{ to } 80 \text{ } \mu\text{m}. \quad (19)$$

Depending on the values of $\sin \phi$ and λ_{1jk} , this could be measured at LEP.

4 Prospects for LEP Searches at Higher Energies

The best testing ground for the plausibility of the ALEPH data will come with the new LEP runs at higher \sqrt{s} . If the ALEPH signal is real and our interpretation correct, we should expect pair production of $\tilde{e}_L \tilde{e}_L$, $\tilde{e}_R \tilde{e}_R$, $\tilde{e}_L \tilde{e}_R$, and $\tilde{\nu}_e \tilde{\nu}_e$ with the rates shown in table 1. The preferred sneutrino decay mode is $\tilde{\nu}_e \rightarrow \bar{d}_j d_k$, as the R -parity conserving decay modes $\tilde{\nu}_e \rightarrow \nu_e e^\pm \tilde{e}_R^\mp$ are suppressed by phase space and by the small mixing between B -ino and W_3 -ino. Therefore all slepton production processes correspond to four-jet events, although the peaks in the distributions of the sum and difference of dijet masses depend on the process.

Charginos are expected to be too heavy to be produced at LEP, even if gaugino mass unification holds. There is however a chance to observe the lightest neutralino χ_1^0 , if the parameter M_1 is in the lower part of the allowed range. The B -ino state χ_1^0 decays into two jets and an electron with more than 80% probability, or else into two jets and a neutrino. The relevant production cross sections are also shown in table 1.

Finally there are very good prospects for the discovery of the Higgs boson. Within the supersymmetric model with minimal Higgs structure, the low values of $\tan\beta$ assumed here imply that the lightest Higgs boson has Standard Model-like couplings and a mass, coming almost entirely from radiative corrections, roughly below 80 GeV [15].

5 The Case of Flavour Universality

As we have mentioned before, in our analysis we have assumed that sleptons of the second and third generations are heavier than those of the first. This assumption is not inconsistent with the strong bounds on individual lepton number conservation, derived from $\mu \rightarrow e\gamma$ and similar processes. An approximate lepton flavour conservation can be the result of an alignment between leptons and sleptons, as a consequence of additional global symmetries [16] or of a dynamical principle [17].

Let us suppose now that the slepton supersymmetry-breaking masses are universal in flavour. Because of the mixing effect, we find that the lightest smuon and stau are lighter than \tilde{e}_R by an amount

$$\Delta_{\tilde{\mu}\tilde{e}} \simeq \frac{(A - \mu \tan\beta)^2 m_\mu^2}{2m_{\tilde{e}_R}(m_{\tilde{e}_L}^2 - m_{\tilde{e}_R}^2)} = \left(\frac{A - \mu \tan\beta}{200 \text{ GeV}}\right)^2 4 \text{ MeV} , \quad (20)$$

$$\Delta_{\tilde{\tau}\tilde{e}} \simeq \frac{(A - \mu \tan\beta)^2 m_\tau^2}{2m_{\tilde{e}_R}(m_{\tilde{e}_L}^2 - m_{\tilde{e}_R}^2)} = \left(\frac{A - \mu \tan\beta}{200 \text{ GeV}}\right)^2 1 \text{ GeV} . \quad (21)$$

Thus the mainly right-handed stau is the lightest supersymmetric particle.

Because of the absence of t -channel contributions, pair productions of smuons, staus, and their corresponding sneutrinos do not suffer from destructive interference and have relatively large cross sections. Some indicative numbers are the following: For $m_{\tilde{\mu}_R} = 48$ GeV, the $\tilde{\mu}_R\tilde{\mu}_R$ production cross section is 0.55 pb for $\sqrt{s} = 130$ GeV and 0.54 pb for $\sqrt{s} = 172$ GeV. For $m_{\tilde{\tau}_R} = 47$ GeV, the $\tilde{\tau}_R\tilde{\tau}_R$ production cross section is 0.59 pb for $\sqrt{s} = 130$ GeV and 0.57 pb for $\sqrt{s} = 172$ GeV. For $m_{\tilde{\mu}_L} = 58$ GeV, the $\tilde{\mu}_L\tilde{\mu}_L$ production cross section is 0.18 pb for $\sqrt{s} = 130$ GeV and 0.43 pb for $\sqrt{s} = 172$ GeV. For $m_{\tilde{\tau}_L} = 59$ GeV, the $\tilde{\tau}_L\tilde{\tau}_L$ production cross section is 0.14 pb for $\sqrt{s} = 130$ GeV and 0.41 pb for $\sqrt{s} = 172$ GeV. Finally, for $m_{\tilde{\nu}} = 58$ (53) GeV, the $\tilde{\nu}\tilde{\nu}$ production cross section is 0.21 (0.49) pb for $\sqrt{s} = 130$ GeV and 0.31 (0.4) pb for $\sqrt{s} = 172$ GeV.

The simultaneous presence of R -parity violating interactions with $\lambda_{ijk} \neq 0$, for different values of the index i , is severely constrained by lepton flavour-transition processes like $\mu \rightarrow e\gamma$. We are then led to assume that the second- and third-generation sleptons do not participate

in the R -parity violating interaction, and consequently their decays have to involve real or virtual \tilde{e}_R , \tilde{e}_L , or $\tilde{\nu}_e$. Their signatures are therefore two jets accompanied by soft leptons or small amounts of missing energy. The presence of the leptons and/or neutrinos is a necessary feature of the transition between different generations of sleptons. Since such events have not been reported by any of the LEP experimental collaborations, we believe that the case of universality has to be rejected.

6 Conclusions

In this paper we have assumed that the controversial ALEPH excess of four-jet events corresponds indeed to a physics signal and we have interpreted it in the context of a supersymmetric model with R -parity violation. If we consider non-universal mass terms for gauginos and for sleptons with different flavours, we find that $\tilde{e}_L\tilde{e}_R$ production can reproduce the four-jet events, while the production of other associated supersymmetric particles occurs at a much lower rate. Within an acceptable range of R -parity violating couplings, both \tilde{e}_L and \tilde{e}_R naturally have a decay branching ratio into two quarks very close to 1.

Our model is compatible with all discernible features emerging from the ALEPH data: the value of the cross section, the dijet mass difference, the absence of bottom quarks in the final state, and the jet charge content. It also predicts a specific angular distribution and a large forward-backward asymmetry in the jet charge. LEP runs at higher energies will be able to confirm or rule out this scenario.

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Table 1: Predictions for new particle production at future LEP runs. We have taken $m_{\tilde{e}_L} = 58$ GeV and $m_{\tilde{e}_R} = 48$ GeV.

Process	Particle mass [GeV]	σ at $\sqrt{s} = 186$ GeV [pb]	σ at $\sqrt{s} = 195$ GeV [pb]
$\tilde{e}_L\tilde{e}_R$	$M_1 = 80$	1.46	1.37
$\tilde{e}_L\tilde{e}_R$	$M_1 = 100$	1.33	1.27
$\tilde{e}_L\tilde{e}_L$	$M_1 = 80$	0.20	0.19
$\tilde{e}_L\tilde{e}_L$	$M_1 = 100$	0.23	0.22
$\tilde{e}_R\tilde{e}_R$	$M_1 = 80$	0.42	0.47
$\tilde{e}_R\tilde{e}_R$	$M_1 = 100$	0.20	0.22
$\tilde{\nu}_e\tilde{\nu}_e$	$m_{\tilde{\nu}_e} = 53$	0.36	0.33
$\tilde{\nu}_e\tilde{\nu}_e$	$m_{\tilde{\nu}_e} = 58$	0.29	0.28
$\chi_1^0\chi_1^0$	$M_1 = 80$	0.62	0.80

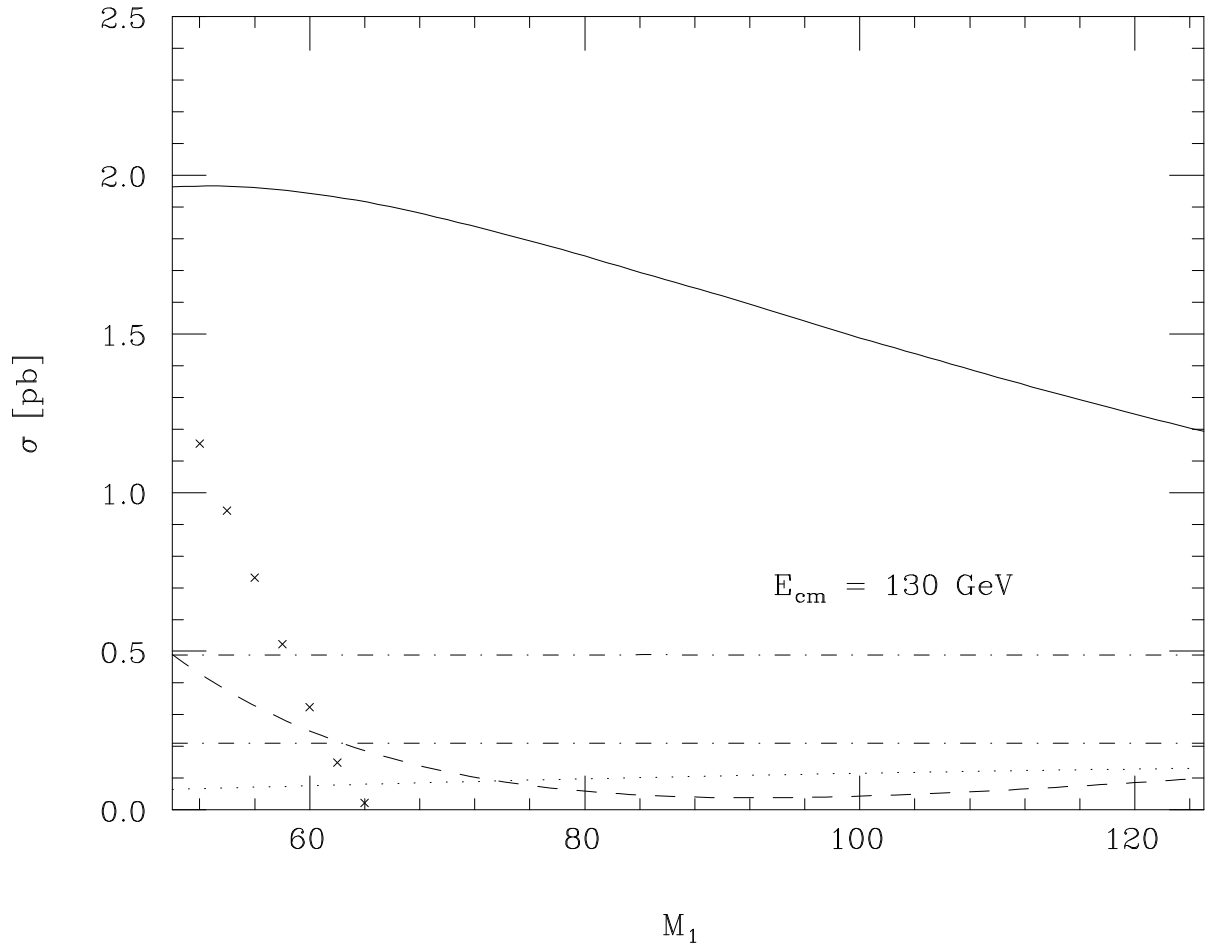


Figure 1: a) Cross sections at $\sqrt{s} = 130$ for the production channels: $\tilde{e}_L\tilde{e}_R$ (solid lines), $\tilde{e}_L\tilde{e}_L$ (dotted lines), $\tilde{e}_R\tilde{e}_R$ (dashed lines), $\tilde{\nu}_e\tilde{\nu}_e$ (dot-dashed lines), and $\chi_1^0\chi_1^0$ (crosses) as a function of the $U(1)$ gaugino mass M_1 . We have taken $m_{\tilde{e}_L} = 58$ GeV and $m_{\tilde{e}_R} = 48$ GeV. The upper dot-dashed line corresponds to $m_{\tilde{\nu}_e} = 53$ GeV and the lower one to $m_{\tilde{\nu}_e} = 58$ GeV.

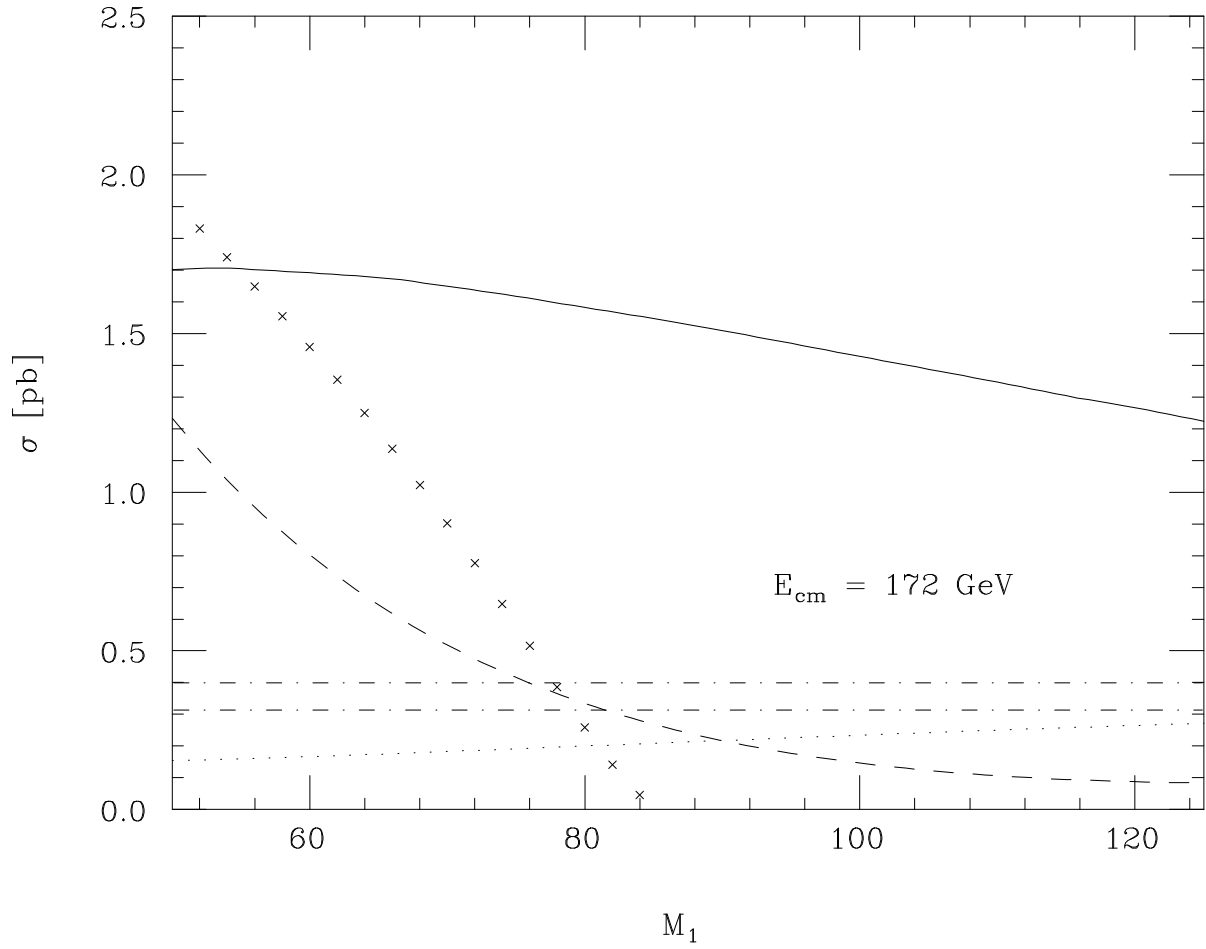


Figure 1: b) The same as Fig.1.a but at $\sqrt{s} = 172$ GeV.

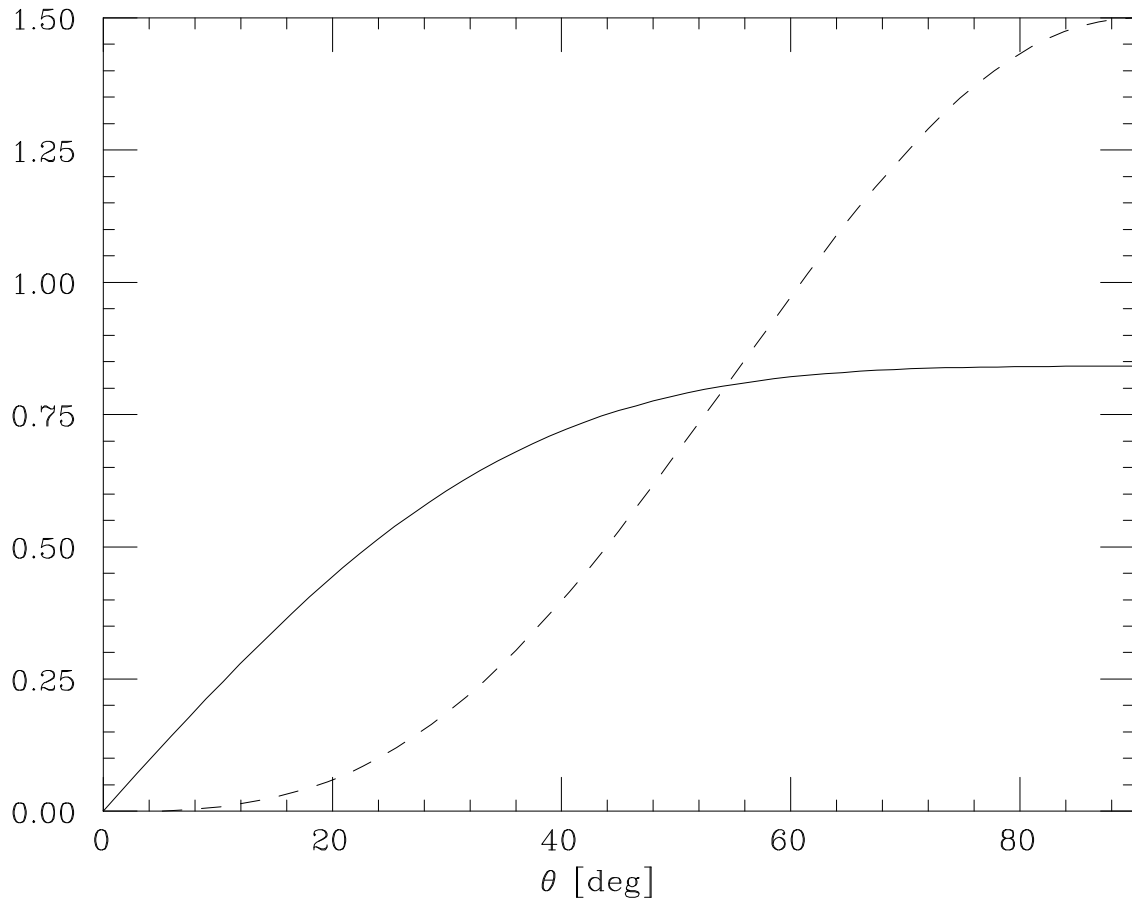


Figure 2: Distributions in the angle θ between the dijet momentum and the beam direction for scalar-particle pair production via gauge bosons in the s -channel (dashed line) and for $\tilde{e}_L \tilde{e}_R$ production (solid line). The curves are normalized so that their integrals over $0 < \theta < \pi/2$ are equal to 1. We have taken $m_{\tilde{e}_L} = 58$ GeV, $m_{\tilde{e}_R} = 48$ GeV, $M_1 = 90$ GeV, and $\sqrt{s} = 136$ GeV.

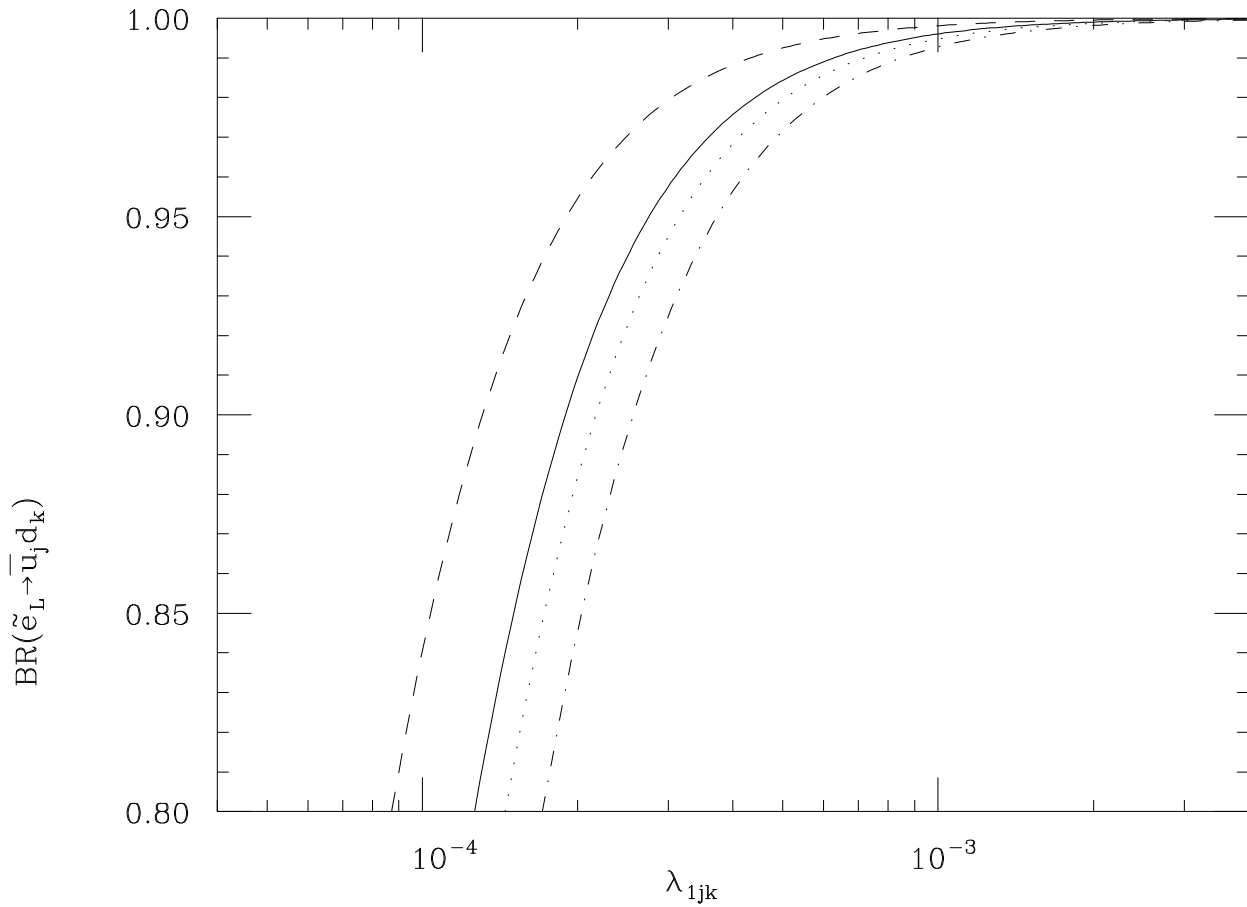


Figure 3: Branching ratio for $\tilde{e}_L \rightarrow \bar{u}_j d_k$, as a function of the R -parity violating coupling λ_{1jk} . We have taken $m_{\tilde{e}_L} = 58$ GeV, $m_{\tilde{e}_R} = 48$ GeV, and $M_1 = 80$ GeV, $m_{\tilde{\nu}_e} = 53$ (dot-dashed line); $M_1 = 80$ GeV, $m_{\tilde{\nu}_e} = 58$ (dotted line); $M_1 = 100$ GeV, $m_{\tilde{\nu}_e} = 53$ (solid line); $M_1 = 100$ GeV, $m_{\tilde{\nu}_e} = 58$ (dashed line).

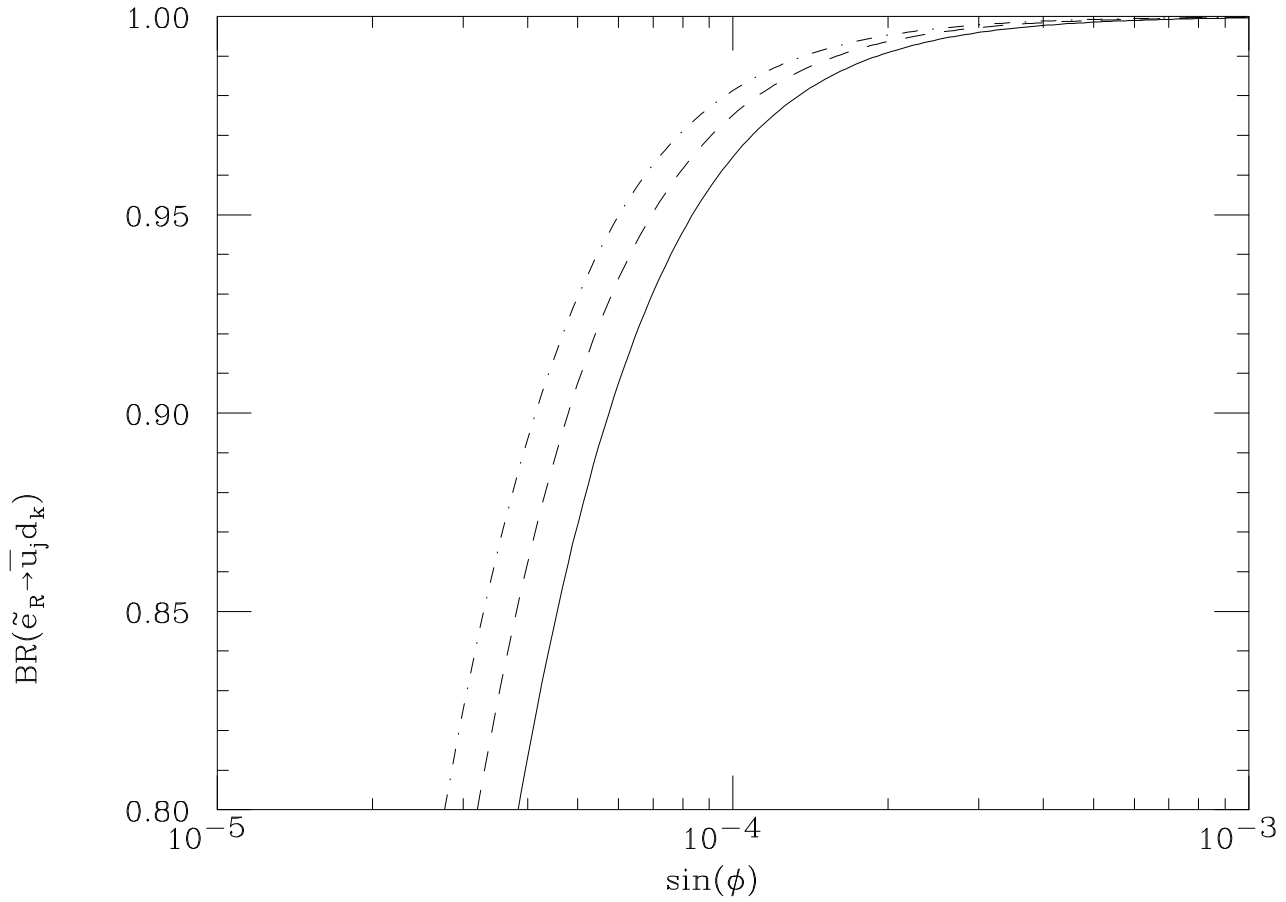


Figure 4: Branching ratio for $\tilde{e}_R \rightarrow \bar{u}_j d_k$, as a function of the left–right mixing angle ϕ . We have taken $m_{\tilde{e}_L} = 58$ GeV, $m_{\tilde{e}_R} = 48$ GeV, $M_1 = 80$ GeV (solid line), 90 GeV (dashed line), or 100 GeV (dot-dashed line).