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Manifestations of R -Parity Violation in Ultrahigh-Energy Neutrino Interactions

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Abstract

Supersymmetric couplings that do not respect R -parity can induce significant changes in the interaction rates of ultrahigh-energy neutrinos through the direct-channel production of superpartner resonances, and can provide new sources of extremely energetic τ -leptons. We analyze the possible observable consequences of \mathcal{R} transitions in large-volume neutrino telescopes.

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I. INTRODUCTION

By observing long-range muons produced in the charged-current reactions $(\nu_\mu, \bar{\nu}_\mu)N \rightarrow \mu^\mp + \text{anything}$, we can hope to detect ultrahigh-energy (UHE) neutrinos from extraterrestrial sources. The diffuse spectrum of neutrinos produced in the interactions of energetic protons in active galactic nuclei (AGNs), neutrinos emitted by gamma-ray bursters, and neutrinos produced as a result of pion photoproduction on the cosmic microwave background are among the targets of UHE neutrino astronomy [1]. Extensive predictions exist for the charged-current and neutral-current cross sections at energies up to 10^{21} eV [2–4], and neutrino telescopes with instrumented volumes approaching 1 km^3 are under active study [1,5].

The primary motivation for neutrino observatories is to search for sources and to probe the processes that produce UHE γ rays in AGNs. However, unconventional processes can modify the expected interaction rates of UHE neutrinos and provide evidence for new particles or new interactions of direct interest to particle physics.

In this paper, we explore supersymmetric processes mediated by R -parity-violating (\mathcal{R}) interactions and investigate their consequences for UHE neutrino interaction rates. We find that direct-channel production of squarks through \mathcal{R} couplings with ordinary particles can significantly enhance the UHE neutrino–nucleon cross sections, and can alter the neutral-current to charged-current ratio for neutrino–nucleon interactions. Similarly, direct-channel production of sleptons can enhance the neutrino–electron cross section at resonance and may provide a new source of UHE τ leptons. These effects may, in time, be observable in neutrino telescopes, and our analysis raises some issues to be borne in mind as km^3 -class neutrino observatories come into being.

II. R -PARITY IN SUPERSYMMETRIC THEORIES

In the standard electroweak theory, gauge invariance suffices to forbid terms in the Lagrangian that change either baryon number or lepton number. The most general supersymmetric (SUSY) extension of the standard model [6], in contrast, allows couplings that change lepton number or baryon number. In general, such terms may lead to an unacceptably short proton lifetime. A simple, though *ad hoc*, solution to this problem is to impose a discrete symmetry called R -parity, which implies a conserved multiplicative quantum number, $R \equiv (-1)^{3B+L+S}$, where B is baryon number, L is lepton number, and S is spin [7]. All ordinary particles are R -parity-even, while all superpartners are R -parity-odd. If R -parity is conserved, superpartners must be produced in pairs and the lightest superpartner, or LSP, is absolutely stable.

However, we have no principle that requires us to impose R -parity on the SUSY Lagrangian. What is more, \mathcal{R} interactions can improve the agreement between theory and precision electroweak measurements (*e.g.*, $Z^0 \rightarrow b\bar{b}, c\bar{c}$), and also offer ready explanations [8] for experimental anomalies such as the high- Q^2 excess reported at HERA [9,10]. Accordingly, it is of interest to consider an R -parity-violating extension of the minimal supersymmetric standard model (MSSM). The most general \mathcal{R} terms in the superpotential consistent with

Lorentz invariance, gauge symmetry, and supersymmetry are¹

$$W_{\mathcal{R}} = \lambda_{ijk} L^i L^j \bar{E}^k + \lambda'_{ijk} L^i Q^j \bar{D}^k + \lambda''_{ijk} \bar{U}^i \bar{D}^j \bar{D}^k \quad (2.1)$$

where i, j, k are generation indices, $L^i \ni (\nu^i, e^i)_L$ and $Q^i \ni (u^i, d^i)_L$ are the left-chiral superfields, and $E^i \ni e^i_R$, $D^i \ni d^i_R$, and $U^i \ni u^i_R$ are the right-chiral superfields, respectively. The Yukawa couplings λ_{ijk} , λ'_{ijk} , and λ''_{ijk} are *a priori* arbitrary, so the \mathcal{R} superpotential (2.1) introduces 45 free parameters.

The LLE and LQD terms change lepton number, whereas the UDD term changes baryon number. Since our interest here is in UHE neutrino interactions, we shall explicitly forbid the UDD interactions [13] as the most economical way to avoid unacceptably rapid proton decay. Expanding the superfield components in (2.1), we obtain the interaction Lagrangians

$$\mathcal{L}_{LLE} = \lambda_{ijk} \left\{ \tilde{\nu}_L^i e_L^j \bar{e}_R^k + \tilde{e}_L^i \nu_L^j \bar{e}_R^k + \tilde{e}_R^{kc} \nu_L^i e_L^j \right\} + \text{h.c.} \quad (2.2)$$

and

$$\begin{aligned} \mathcal{L}_{LQD} = \lambda'_{ijk} \left\{ \tilde{\nu}_L^i d_L^j \bar{d}_R^k - \tilde{e}_L^i u_L^j \bar{d}_R^k + \tilde{d}_L^j \nu_L^i \bar{d}_R^k \right. \\ \left. - \tilde{u}_L^j e_L^i \bar{d}_R^k + \tilde{d}_R^{kc} \nu_L^i d_L^j - \tilde{d}_R^{kc} e_L^i u_L^j \right\} + \text{h.c.} \end{aligned} \quad (2.3)$$

The \mathcal{R} couplings in (2.2) and (2.3) modify supersymmetric phenomenology in several important ways: processes that change lepton number are allowed, superpartners can be produced singly, and the LSP—now unstable against decay into ordinary particles—is no longer constrained to be a neutral color singlet to avoid cosmological embarrassments.

The remarkable agreement between present data and standard-model expectations implies very restrictive bounds on the strength of \mathcal{R} operators [14–17]. LQD -type interactions of electron neutrinos or antineutrinos with the first-generation quarks found in nucleon targets are highly constrained. We consider for definiteness the case of a 200-GeV/ c^2 sfermion. The absence of a signal for neutrinoless double-beta decay implies that $\lambda'_{111} \leq 0.002$, for squark and gluino masses of 200 GeV/ c^2 . Universality of the charged-current interactions implies that $\lambda'_{112,113} \leq 0.04$ (at the 2σ level). The couplings $\lambda'_{121,131}$ are also constrained: if the Cabibbo mixing in the light-quark sector arises solely from the *down* quarks, one may argue that the upper bound on the branching ratio for the decay $K \rightarrow \pi \nu \bar{\nu}$ [18] limits $\lambda'_{121} < 0.02$ [19]. If on the other hand this mixing arises from the *up* sector, an effect would have been seen in neutrinoless double beta decay. In either case, λ'_{121} cannot be very large. Finally, atomic parity violation gives the constraint $\lambda'_{121,131} \leq 0.07$. In view of these constraints, we shall not consider the effects of R -parity violation on $\nu_e N$ and $\bar{\nu}_e N$ cross sections.

Experimental limits on LQD couplings that involve ν_μ are less restrictive. So too are the limits on LLE couplings, both for ν_μ and ν_e . To illustrate some possible consequences of

¹We suppress here the $SU(2)_L$ and $SU(3)_c$ indices. Symmetry under $SU(2)_L$ implies that the first term is antisymmetric under $i \leftrightarrow j$, while $SU(3)_c$ symmetry dictates that the third term is antisymmetric under $j \leftrightarrow k$. We neglect bilinear terms that mix lepton and Higgs superfields [11]. Discussions of the phenomenological implications of such terms can be found in the literature [12].

\mathcal{R} couplings for UHE neutrino interactions, we shall analyze separately the effects of LQD terms on $(\nu_\mu, \bar{\nu}_\mu)N$ scattering and of LLE terms on $(\nu_{e,\mu}, \bar{\nu}_{e,\mu})e$ scattering [20], considering in turn the influence of massive squarks and sleptons in light of existing constraints on λ'_{ijk} and λ_{ijk} . In Table I we summarize the constraints on the \mathcal{R} Yukawa couplings that are relevant for our studies. In each example we consider, we shall assume that only one \mathcal{R} coupling can be sizeable [21].

III. NEUTRINO-NUCLEON SCATTERING

The dominant mechanisms for producing UHE photons and neutrinos are expected to be

$$p(p/\gamma) \rightarrow \pi^0 + \text{anything} \quad (3.1)$$

$$\quad \quad \quad \downarrow$$

$$\quad \quad \quad \gamma\gamma$$

and

$$p(p/\gamma) \rightarrow \pi^\pm + \text{anything} \quad (3.2)$$

$$\quad \quad \quad \downarrow$$

$$\quad \quad \quad \mu\nu_\mu$$

$$\quad \quad \quad \downarrow$$

$$\quad \quad \quad e\nu_e\nu_\mu.$$

If π^+ , π^- , and π^0 are produced in equal numbers, the relative populations of the neutral particles will be $2\gamma : 2\nu_\mu : 2\bar{\nu}_\mu : 1\nu_e : 1\bar{\nu}_e$. There are no significant conventional sources of ν_τ and $\bar{\nu}_\tau$ [22].

At low neutrino energies ($E_\nu \ll M_W^2/2M$, where M_W is the intermediate-boson mass and M is the nucleon mass), differential and total cross sections for the reactions $\nu N \rightarrow \mu + \text{anything}$ are approximately proportional to the neutrino energy. For neutrino energies in the range $10 \text{ GeV} \leq E_\nu \leq 1 \text{ TeV}$, the cross sections computed [2] using the CTEQ3 parton distributions [24] are reproduced by simple power-law forms:

$$\begin{aligned} \sigma_{\text{CC}}(\nu N) &= 8.66 \times 10^{-39} \text{ cm}^2 \left(\frac{E_\nu}{1 \text{ GeV}} \right)^{0.953} \\ \sigma_{\text{NC}}(\nu N) &= 2.66 \times 10^{-39} \text{ cm}^2 \left(\frac{E_\nu}{1 \text{ GeV}} \right)^{0.958} \\ \sigma_{\text{CC}}(\bar{\nu} N) &= 3.83 \times 10^{-39} \text{ cm}^2 \left(\frac{E_\nu}{1 \text{ GeV}} \right)^{0.982} \\ \sigma_{\text{NC}}(\bar{\nu} N) &= 1.34 \times 10^{-39} \text{ cm}^2 \left(\frac{E_\nu}{1 \text{ GeV}} \right)^{0.983} . \end{aligned} \quad (3.3)$$

Above $E_\nu \approx 10^{12} \text{ eV}$, the gauge-boson propagator restricts the momentum transfer Q^2 to values near M_W^2 and damps the cross section. (Parallel remarks apply to the neutral-current reactions $\nu N \rightarrow \nu + \text{anything}$ mediated by Z exchange.) For neutrino energies in the range $10^{15} \text{ eV} \leq E_\nu \leq 10^{21} \text{ eV}$, good representations of the cross sections are given by [2]:

$$\begin{aligned}
\sigma_{\text{CC}}(\nu N) &= 2.69 \times 10^{-36} \text{ cm}^2 \left(\frac{E_\nu}{1 \text{ GeV}} \right)^{0.402} \\
\sigma_{\text{NC}}(\nu N) &= 1.06 \times 10^{-36} \text{ cm}^2 \left(\frac{E_\nu}{1 \text{ GeV}} \right)^{0.408} \\
\sigma_{\text{CC}}(\bar{\nu} N) &= 2.53 \times 10^{-36} \text{ cm}^2 \left(\frac{E_\nu}{1 \text{ GeV}} \right)^{0.404} \\
\sigma_{\text{NC}}(\bar{\nu} N) &= 0.98 \times 10^{-36} \text{ cm}^2 \left(\frac{E_\nu}{1 \text{ GeV}} \right)^{0.410} .
\end{aligned} \tag{3.4}$$

The constraints we have reviewed on \mathcal{R} couplings involving muon neutrinos or antineutrinos and up or down quarks permit significant modifications to the standard-model cross sections for $\nu_\mu N$ and $\bar{\nu}_\mu N$ interactions.² We choose to limit our considerations to interactions involving first-generation quarks, which are the predominant constituents of the nucleon. A similar analysis could be carried out for the minority constituents s, c , *etc.* The relevant R -parity-violating couplings then are λ'_{21j} , which drives reactions mediated by right-handed squarks, and λ'_{2j1} , which drives reactions mediated by left-handed squarks.

$\nu_\mu N$ interactions. The charged-current reaction $\nu_\mu N \rightarrow \mu^- + \text{anything}$ can receive contributions from the s -channel formation process $\nu_\mu d_L \rightarrow \tilde{d}_R^k \rightarrow \mu^- u_L$, which involves valence quarks, and from u -channel exchange of \tilde{d}_R^k in the reaction $\nu_\mu \bar{u} \rightarrow \bar{d} \mu^-$, which involves only sea quarks. The cross section formulas are collected in the Appendix. We consider the influence of a right-handed squark \tilde{d}_R^k with mass $\tilde{m}_R^k = 200(400) \text{ GeV}/c^2$ and \mathcal{R} Yukawa couplings $\lambda'_{21k} = 0.2(0.4)$, which are about the maximum values allowed. With these choices for masses and Yukawa couplings, the total width of the right-handed squark is $\Gamma(\tilde{d}_R^k) = 0.32 \text{ GeV}$ for $\lambda'_{21k} = 0.2$, and $\Gamma(\tilde{d}_R^k) = 2.55 \text{ GeV}$ for $\lambda'_{21k} = 0.4$, with equal branching fractions into $\mu^- u$ and $\nu_\mu d$. We have made here the implicit assumption that the squark decays to gauginos are suppressed compared to the R -parity-violating decays.³ We will revisit this assumption when we discuss the contribution of $LL\bar{E}$ operators to neutrino-electron scattering. Left-handed squarks couple either to neutrinos or to charged leptons, but not both, so cannot contribute to the charged-current reaction.

The resonance condition for forming \tilde{d}_R^k in νq interactions is $x = \tilde{m}_R^k{}^2/2ME_\nu$, where x is the momentum fraction carried by the quark and M is the nucleon mass. As a consequence of the spread in quark momenta, the resonance peaks are not narrow, but are broadened and shifted above the threshold energies $E_\nu^{\text{th}} = \tilde{m}_R^k{}^2/2M \approx 2.13$ and $8.52 \times 10^4 \text{ GeV}$, respectively. We have calculated the cross sections with and without the \mathcal{R} contribution using the CTEQ3 parton distributions [24]. Figure 1(a) shows that the modifications to the standard-model

²Since the cases we investigate do not involve first-generation leptons, they have no observable consequences for the $e^\pm p$ reactions studied at HERA.

³This is the case for relatively large μ and M_2 , which result in large gaugino masses. While the Tevatron bounds on second generation scalar leptoquarks are impressively tight for light sfermions, they are greatly relaxed for smaller branching ratios to leptoquark-type decays and for higher sfermion masses. See Ref. [23].

cross section are appreciable: in the two examples we consider, the charged-current cross section is enhanced by up to 60% and 30%, respectively. Far above the resonance bump, the cross section is enhanced by about 20% over the standard-model level.

The squark contributions modify the angular distributions as well. In Figure 2(a), we plot the mean inelasticity $y_{\text{av}} \equiv (1/\sigma) \int_0^1 dy y (d\sigma/dy)$ as a function of E_ν . The isotropic distribution of the squark decay products enhances y_{av} at the resonance bump and beyond.

The right-handed squark \tilde{d}_R^k has a similar influence on the neutral-current reaction $\nu_\mu N \rightarrow \nu_\mu + \text{anything}$. There is an s -channel contribution from the formation process $\nu_\mu d_L \rightarrow \tilde{d}_R^k \rightarrow \nu_\mu d_L$, which involves valence quarks, and a u -channel contribution to the reaction $\nu_\mu \bar{d} \rightarrow d\nu_\mu$, which involves sea quarks. Again, the resonance bumps shown in Figure 3(a) significantly enhance the standard-model cross section: the neutral-current cross section is increased by as much as 120% and 55% in our two examples. Far above the resonance bumps, the increase is about 40%. The \tilde{d}_R^k contribution is more visible than in the charged-current case of Figure 1(a) because $\sigma_{\text{NC}}/\sigma_{\text{CC}} \approx 0.4$ in the standard model, whereas the \mathcal{R} cross sections are the same for neutral-current and charged-current interactions.

Left-handed squarks can contribute to the neutral-current reaction. An s -channel contribution arises from \tilde{d}_L^{kc} formation in the reaction $\nu_\mu \bar{d}_R \rightarrow \tilde{d}_L^{kc} \rightarrow \nu_\mu \bar{d}_R$ involving the light-quark sea of the nucleon, while u -channel \tilde{d}_L^k -exchange drives the reaction $\nu_\mu d_R \rightarrow d_R \nu_\mu$ involving valence quarks as well. As we did for \tilde{d}_R^k , we consider squark masses $\tilde{m}_L^k = 200$ and 400 GeV/ c^2 , and \mathcal{R} Yukawa couplings $\lambda'_{2k1} = 0.2$ and 0.4 . In distinction to the \tilde{d}_R^k , the \tilde{d}_L^{kc} decays only into $\nu_\mu \bar{d}$ through \mathcal{R} interactions; its total width $\Gamma(\tilde{d}_L^k) = \frac{1}{2}\Gamma(\tilde{d}_R^k)$. The influence of \tilde{d}_L^k on the νN neutral-current cross section is shown in Figure 3(a). The unit branching fraction into the neutral-current mode means that the effect is pronounced, even though the s -channel reaction involves only sea quarks. Because the sea is softer than the valence, the influence of \tilde{d}_L^k is deferred to energies well above the nominal threshold. The maximum enhancement over the standard model is about 80% (50%) for $\tilde{m}_L^k = 200$ (400) GeV/ c^2 . This enhancement persists to the highest energies we consider.

The neutral-current to charged-current ratio is an interesting diagnostic for the new physics represented by \mathcal{R} couplings. In Figure 4(a) we compare the ratio $\sigma_{\text{NC}}/\sigma_{\text{CC}}$ expected in the standard model with our four \mathcal{R} examples. The most prominent feature is the rapid rise of $\sigma_{\text{NC}}/\sigma_{\text{CC}}$ at the onset of \tilde{d}_L^k formation; the neutral-current to charged-current ratio is roughly doubled for squarks of 200 or 400 GeV/ c^2 .

The ability to measure the neutral-current to charged-current ratio has not, until now, been conceived as a strength of neutrino telescopes. Our study of R -parity-violating interactions shows that it would add markedly to the sensitivity that neutrino observatories will have to new physics.

$\bar{\nu}_\mu N$ interactions. The charged-current reaction $\bar{\nu}_\mu N \rightarrow \mu^+ + \text{anything}$ can receive contributions from the s -channel formation reaction $\bar{\nu}_\mu d_L \rightarrow \tilde{d}_R^{kc} \rightarrow \mu^+ \bar{u}_L$, which involves sea quarks, and from u -channel exchange of \tilde{d}_R^k in the reaction $\bar{\nu}_\mu u_L \rightarrow d_L \mu^+$, which involves valence quarks. For the cases we consider, illustrated in Figure 1(b), the enhancement over the standard-model cross section is no more than 20% (15%) for $\tilde{m}_R^k = 200$ (400) GeV/ c^2 .

The right-handed squark \tilde{d}_R^{kc} has a similar influence on the neutral-current reaction $\bar{\nu}_\mu N \rightarrow \bar{\nu}_\mu + \text{anything}$. There is an s -channel contribution from the formation process $\bar{\nu}_\mu \bar{d}_L \rightarrow \tilde{d}_R^{kc} \rightarrow \bar{\nu}_\mu \bar{d}_L$, which involves only sea quarks, and a u -channel contribution to the

reaction $\bar{\nu}_\mu d_L \rightarrow d_L \bar{\nu}_\mu$, which involves valence quarks as well. The changes to the standard-model cross sections shown in Figure 3(b) are relatively more important than in the charged-current case, because the standard-model neutral-current cross section is smaller. The maximum enhancements are 50% (30%) for $\tilde{m}_R^k = 200$ (400) GeV/ c^2 . We see in Figure 4(b) that the right-handed squark does not lead to a dramatic effect in the neutral-current to charged-current ratio.

Left-handed squarks can alter the neutral-current cross section significantly. An s -channel contribution arises from the formation reaction $\bar{\nu}_\mu d_R \rightarrow \tilde{d}_L^k \rightarrow \bar{\nu}_\mu d_R$ involving valence quarks, and u -channel \tilde{d}_L^k exchange drives the reaction $\bar{\nu}_\mu \bar{d}_R \rightarrow \bar{d}_R \bar{\nu}_\mu$, involving sea quarks only. The influence of \tilde{d}_L^k on the $\bar{\nu}_\mu N$ neutral-current reaction is shown in Figure 3(b). For the cases we consider, \mathcal{R} contributions can be 3 times (2 times) the standard-model cross section, for $\tilde{m}_L^k = 200$ (400) GeV/ c^2 . As a consequence, the effect of a left-handed squark on the neutral-current to charged-current ratio is extremely pronounced: we see in Figure 4(b) that $\sigma_{\text{NC}}/\sigma_{\text{CC}}$ is quadrupled (doubled) for our two examples. Far above threshold, σ_{NC} is increased by 95% (75%), and the neutral-current to charged-current ratio remains twice ($1\frac{3}{4}$ times) the standard-model value.

Consequences for neutrino observatories. Until now, we have considered $\nu_\mu N$ and $\bar{\nu}_\mu N$ interactions separately. Neutrino telescopes will not distinguish between events induced by neutrinos and antineutrinos. Thus we should rather consider the sum of the $\nu_\mu N$ and $\bar{\nu}_\mu N$ cross sections. The four panels in Figure 5 depict σ_{CC} , y_{av} for charged-current interactions, σ_{NC} , and $\sigma_{\text{NC}}/\sigma_{\text{CC}}$ for summed $\nu_\mu N + \bar{\nu}_\mu N$ interactions. The significance of the resonance threshold is reduced, but the deviations from the standard-model expectations remain significant. The increases in the mean inelasticity and the neutral-current to charged-current ratio seem especially promising signals for unconventional interactions, because they do not require a knowledge of the incident neutrino flux.

IV. NEUTRINO INTERACTIONS ON ELECTRON TARGETS

Because of the electron's small mass, neutrino-electron interactions in matter can generally be neglected with respect to neutrino-nucleon interactions [25]. There is one exceptional case in the standard model: resonant formation of the intermediate boson in $\bar{\nu}_e e \rightarrow W^-$ interactions at 6.3 PeV. The resonant cross section is larger than the νN cross section at any energy up to 10^{21} eV. This singular case makes it important to bear the neutrino-electron cross sections in mind when assessing the capabilities of neutrino telescopes. We now consider similar singular cases that can arise through R -parity-violating interactions. Once we allow LLE couplings, a wide variety of new production and decay channels open up. Because the \mathcal{R} couplings are constrained to be small, only channels that involve resonant sparticle production can display sizeable effects.

νe interactions. It follows from the interaction Lagrangian (2.2) that a left-handed slepton can only give rise to additional u -channel diagrams generated by the operator $L_i L_j \bar{E}_k$, with $i \neq j$ from SU(2) invariance. R -parity-violating contributions to the reactions $(\nu_\mu, \nu_e)e \rightarrow (\nu_\mu, \nu_e)e$ are mediated by $\tilde{\tau}_L$ or $(\tilde{e}_L, \tilde{\mu}_L)$. These lead to unobservably small modifications to the standard-model cross section.

In contrast, a right-handed slepton can appear in an s -channel diagram, and may thus

be produced as a resonance. Because of the SU(2) invariance of the $L_i L_j \bar{E}_k$ operator, the options are $(\nu_\mu, \nu_\tau)e \rightarrow \tilde{e}_R^k$. For the astrophysically interesting case of incident ν_μ , and considering only one \mathcal{R} operator at a time, the two final states $\nu_\mu e$ and $\nu_e \mu$ occur with equal probability in the decay of \tilde{e}_R^k . Within the standard model, $\nu_\mu e \rightarrow \nu_\mu e$ elastic scattering proceeds through a t -channel Z -exchange graph. At low energies, the cross section grows with neutrino energy, but reaches a plateau at 10 pb around $E_\nu = 10^{17}$ eV. The charge-exchange reaction $\nu_\mu e \rightarrow \mu \nu_e$ proceeds by W -exchange, which leads to a cross section about an order of magnitude larger. In either channel, the magnitude of the resonant cross section is governed by the strength of the \mathcal{R} coupling λ_{12k} and the slepton mass. Before discussing the observability of the $\nu_\mu e_L \rightarrow \tilde{e}_R^k \rightarrow (\nu_\mu e_L, \nu_e \mu_L)$ signals, let us explore the signals that arise in $\bar{\nu}_e e$ and $\bar{\nu}_\mu e$ interactions.

$\bar{\nu}_e e$ and $\bar{\nu}_\mu e$ interactions. Resonant slepton production can occur in the reactions $\bar{\nu}_e e_R \rightarrow \tilde{\tau}_L$ or $\tilde{\mu}_L$ and $\bar{\nu}_\mu e_R \rightarrow \tilde{\tau}_L$ or \tilde{e}_L . If we suppose that only one \mathcal{R} coupling is nonzero, then the left-handed slepton cannot mediate flavor-changing processes such as $\bar{\nu}_e e \rightarrow \bar{\nu}_\mu \mu$.

In Figure 6 we present the cross sections for (a) $\bar{\nu}_e e$ and (b) $\bar{\nu}_\mu e$ elastic scattering. Within the standard model, $\bar{\nu}_e e$ elastic scattering is driven by s -channel W^- formation and t -channel Z^0 exchange; the behavior of $\bar{\nu}_\mu e$ elastic scattering is very similar to that of the reaction $\nu_\mu e \rightarrow \nu_\mu e$ described above. A left-handed slepton with a single \mathcal{R} coupling may occur as an s -channel resonance in either $\bar{\nu}_e e$ elastic scattering or $\bar{\nu}_\mu e$ elastic scattering, but not both. If R -parity–*conserving* decays are absent, as assumed in Figure 6, the slepton resonance is very narrow and impressively tall. The peak cross section is

$$\sigma_{\text{peak}}(\bar{\nu}_i e \rightarrow \bar{\nu}_i e) \approx \frac{8\pi}{m_{\tilde{e}_L^k}^2} \cdot \left(\frac{\Gamma(\tilde{e}_L^k \rightarrow \bar{\nu}_i e)}{\Gamma_{\text{tot}}(\tilde{e}_L^k)} \right)^2, \quad (4.1)$$

where the final factor is unity, by assumption. The total width of the slepton resonance is $\Gamma_{\text{tot}}(\tilde{e}_L^k) = (\lambda^2/16\pi)m_{\tilde{e}_L^k}$. For the case of $\nu_\mu e \rightarrow \tilde{e}_R^k$ scattering discussed above, the peak cross section is similar in form. In that case, however, \tilde{e}_R^k decays with equal probability to $\nu_\mu e$ or $\nu_e \mu$. The peak cross section for $\nu_\mu e \rightarrow \nu_\mu e$ or $\nu_\mu e \rightarrow \nu_e \mu$ is accordingly rescaled by a factor 1/4, and the total width is doubled.

Detecting the slepton resonance. The modest flux of UHE neutrinos and the factor-of-two energy resolution anticipated for neutrino observatories mean that it will be difficult to resolve such a narrow structure from the standard-model background. One distinguishing characteristic is that the slepton resonance will only be produced in downward-going interactions. In water-equivalent units, the interaction length is

$$L_{\text{int}}^{(e)} = \frac{1}{\sigma(E_\nu)(10/18)N_A}, \quad (4.2)$$

where $N_A = 6.022 \times 10^{23} \text{ mol}^{-1} = 6.022 \times 10^{23} \text{ cm}^{-3}$ (water equivalent) is Avogadro's number and $(10/18)N_A$ is the number of electrons in a mole of water. At the peak of a 200- (400-) GeV/ c^2 slepton resonance produced in $\bar{\nu}e$ interactions, the interaction length is 122 (488) kmwe. The resonance is effectively extinguished for neutrinos that traverse the Earth, whose diameter is 1.1×10^5 kmwe. Similar remarks apply to the $\nu_\mu e_L \rightarrow \tilde{e}_R^k \rightarrow (\nu_\mu e_L, \nu_e \mu_L)$ cases.

A more promising strategy to observe a slepton resonance is to consider final states that would clearly stand out above the background. One such possibility arises if both LLE and LQD operators exist. Then the slepton could decay into two jets, for which the background from $\bar{\nu}_e e \rightarrow W^- \rightarrow q\bar{q}$ is negligible at high energies. Neutral-current $\nu N \rightarrow \nu + \text{anything}$ interactions might constitute an additional—and large—background.

A much more interesting possibility arises if neutralinos are relatively light. Now the slepton may also decay into the corresponding lepton and a light neutralino, which decays in turn into leptons and neutrinos. Reactions that produce final-state τ -leptons are of special interest. The decay length of a 1-PeV τ is about 50 meters, so the production and subsequent decay of a τ at ultrahigh energies will give rise to a characteristic “double-bang” signature in a water or ice Čerenkov detector. We consider

$$\nu_\mu e_L^- \rightarrow \tilde{\tau}_R^- \rightarrow \tau^- \tilde{\chi}^0 \quad \begin{array}{l} \searrow \\ \downarrow \\ \searrow \end{array} \tau_R^+ \nu_e \mu_L^- \text{ or } \tau_R^+ \nu_\mu e_L^- \quad (4.3)$$

and

$$\bar{\nu}_e e_R^- \rightarrow \tilde{\tau}_L^- \rightarrow \tau^- \tilde{\chi}^0 \quad \begin{array}{l} \searrow \\ \downarrow \\ \searrow \end{array} \tau_L^+ \bar{\nu}_e e_R^-, \quad (4.4)$$

where $\bar{\nu}$ refers to $\bar{\nu}_e$ or $\bar{\nu}_\mu$. Close to the resonance (*i.e.*, where the t - and u -channel exchanges can be neglected in comparison to the s -channel pole), the cross section for $\tau \tilde{\chi}^0$ production can be approximated by a Breit-Wigner formula,

$$\begin{aligned} \sigma &= \frac{8\pi s}{m_{\tilde{\tau}}^2} \frac{\Gamma(\tilde{\tau} \rightarrow e^- \nu) \Gamma(\tilde{\tau} \rightarrow \tau^- \tilde{\chi}^0)}{(s - m_{\tilde{\tau}}^2)^2 + m_{\tilde{\tau}}^2 \Gamma_{\text{total}}^2} \left[\frac{s - m_{\tilde{\chi}^0}^2}{m_{\tilde{\tau}}^2 - m_{\tilde{\chi}^0}^2} \right]^2 \\ &\rightarrow \frac{8\pi}{m_{\tilde{\tau}}^2} B(\tilde{\tau} \rightarrow \nu e) B(\tilde{\tau} \rightarrow \tau \tilde{\chi}^0), \text{ as } s \rightarrow m_{\tilde{\tau}}^2. \end{aligned} \quad (4.5)$$

The signal cross section can thus be deduced from Figure 6(b), scaled by the appropriate branching fractions. The standard-model background for τ production arises only from the reaction $\bar{\nu}_e e \rightarrow W^- \rightarrow \bar{\nu}_\tau \tau$; it drops below a few pb at about 10^{17} eV. To good approximation, the production of taus through a slepton resonance with a mass of 200 GeV/ c^2 or greater is background-free.

In Figure 7 we plot branching-fraction contours for the decays $\tilde{\tau}_{L,R} \rightarrow \tilde{\chi}^0 \tau_{L,R}$ in the $M_2 - \mu$ plane for small and large values of $\tan \beta$, the ratio of the vacuum expectation value of the Higgs boson that couples to up quarks to that of the Higgs boson that couples to down quarks. Here M_2 is the SU(2) gaugino mass parameter and μ is the Higgsino mass parameter. The R -parity-conserving decay mode competes with, and even dominates, the \cancel{R} decay mode(s) for a wide range of supersymmetric parameters. For $\tilde{\tau}_R$, the \cancel{R} decays are divided equally between the $\nu_e \mu^-$ and $\nu_\mu e^-$ modes, but the R -conserving decay into charginos is very suppressed. At small and moderate values of M_2 and μ , the tau-neutralino decay dominates and the total width of $\tilde{\tau}_R$ can be as large as 0.9 GeV. For $\tilde{\tau}_L$, the only \cancel{R} decay is into $\nu_e \mu^-$, but the R -conserving decays can include both the $\tau \tilde{\chi}^0$ mode and the $\nu_\tau \tilde{\chi}^-$ mode. The best bet for unconventional tau production thus seems to be through the production of $\tilde{\tau}_R$ in $\nu_\mu e$ collisions.

Because we know of no conventional astrophysical sources of tau neutrinos, the observation of an excess of energetic taus in a neutrino telescope has been suggested as a signature for neutrino oscillations [26]. We learn from the above discussion that if a massive slepton is produced through \mathcal{R} interactions, it may well be a source of energetic taus that represents a different sort of physics beyond the standard model.

V. CONCLUSION

Supersymmetric theories with R -parity-violation are an interesting generalization of the usual physics beyond the standard model. Relaxing the assumption that R -parity is conserved alters the phenomenology of supersymmetry in a crucial way. In accelerator searches for supersymmetric particles, the missing energy signature is removed because the lightest supersymmetric particle can decay into ordinary particles, decay chains are modified, and superpartners can be produced singly. We have shown that R -parity-violating processes can also have significant effects on ultrahigh-energy neutrino cross sections, if the \mathcal{R} couplings are as large as allowed by current bounds.

We have considered the influence of squarks with \mathcal{R} couplings on the cross sections for UHE $\nu_\mu N$ and $\bar{\nu}_\mu N$ interactions, and the effect of sleptons with \mathcal{R} couplings on $\nu_\mu e$, $\bar{\nu}_e e$ and $\bar{\nu}_\mu e$ interactions. For the case of neutrino-nucleon collisions, we have found that the neutral-current to charged-current ratio can be a very useful diagnostic for new physics. The ability to measure neutral-current interactions and characterize their energy should receive increased attention in design studies for neutrino telescopes.

A remarkable sign for R -parity-violating effects in neutrino-electron scattering would be the formation of slepton resonances in narrow energy bands. The large cross sections at resonance might be observed through the extinction of neutrinos passing through the Earth near the resonance energy or by the anomalous production of tau leptons in downward-going events only. We have explored scenarios that lead to the production of energetic taus through R -parity-violating neutrino-electron interactions. If a $\tilde{\tau}$ has both R -parity-conserving and R -parity-violating decays, we would expect resonant production of $\tau\tilde{\chi}^0$, with a subsequent decay of the LSP into $\tau\ell\nu$. Alternatively, if more than one \mathcal{R} coupling is appreciable, we could expect resonant $\tilde{e}^k \rightarrow \tau\nu$ production. Either could provide an explanation for τ events that does not invoke neutrino oscillations. The key to separating the two \mathcal{R} effects lies in observing the decay products of the $\tilde{\chi}^0$ by searching for associated leptons in τ -bearing events.

It is plain that the observation of the unconventional effects we have examined here is some time away, even if sizeable \mathcal{R} couplings do occur in Nature. Our intent is to widen the scientific horizons of neutrino observatories and to encourage imaginative thinking about the range of measurements they might make.

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APPENDIX: CROSS SECTION FORMULAE

It is straightforward to calculate the inclusive cross section for the reactions $\nu_\mu N \rightarrow \mu + \text{anything}$ and $\nu N \rightarrow \nu + \text{anything}$, where $N = \frac{1}{2}(p + n)$ is an isoscalar nucleon, in the renormalization-group-improved parton model. The differential cross section is written in terms of the Bjorken scaling variables $x = Q^2/2ME_\nu$ and $y = \nu/E_\nu$ as

$$\frac{d^2\sigma}{dx dy} = \frac{ME_\nu}{16\pi} \sum_f \left[|a_f|^2 + |b_f|^2 (1-y)^2 \right] x f(x, Q^2), \quad (\text{A1})$$

where $-Q^2$ is the squared invariant momentum transfer between the incident and outgoing lepton, $\nu = E_\nu - E'$ (with E' the energy of the outgoing lepton in the target frame) is the lepton energy loss, M_W is the W -boson mass, and $G_F = 1.16632 \times 10^{-5} \text{ GeV}^{-2}$ is the Fermi constant. Here $f(x, Q^2)$ represents the parton densities within the nucleon; the coefficients a_f and b_f for interactions on parton species f depend on the particular process under consideration.

For charged-current interactions, the only nonzero coefficients are

$$\begin{aligned} a_{d_j}^{\text{CC}} &= \frac{g^2}{Q^2 + M_W^2} - \sum_k \frac{|\lambda'_{ijk}|^2}{xs - \tilde{m}_R^{k2} + i\tilde{m}_R^k \Gamma(\tilde{d}_R^k)}, \\ b_{u_j}^{\text{CC}} &= \frac{g^2}{Q^2 + M_W^2} - \sum_k \frac{|\lambda'_{ijk}|^2}{Q^2 - xs - \tilde{m}_R^{k2}}, \end{aligned} \quad (\text{A2})$$

where $s = 2ME_\nu$ is the square of the c.m. energy, the square of the $SU(2)_L$ gauge coupling is $g^2 = 8G_F M_W^2 / \sqrt{2}$, and the λ'_{ijk} are the \cancel{R} Yukawa couplings.

For neutral-current interactions, we have

$$\begin{aligned} a_{d_j}^{\text{NC}} &= \frac{g^2}{2(1-x_W)} \frac{L_d}{Q^2 + M_Z^2} - \sum_k \frac{|\lambda'_{ijk}|^2}{xs - \tilde{m}_R^{k2} + i\tilde{m}_R^k \Gamma(\tilde{d}_R^k)} \\ b_{d_j}^{\text{NC}} &= \frac{g^2}{2(1-x_W)} \frac{R_d}{Q^2 + M_Z^2} + \sum_k \frac{|\lambda'_{ijk}|^2}{Q^2 - xs - \tilde{m}_L^{k2}} \\ a_{u_j}^{\text{NC}} &= \frac{g^2}{2(1-x_W)} \frac{L_u}{Q^2 + M_Z^2} \\ b_{u_j}^{\text{NC}} &= \frac{g^2}{2(1-x_W)} \frac{R_u}{Q^2 + M_Z^2} \end{aligned} \quad (\text{A3})$$

where M_Z is the Z -boson mass, $x_W \equiv \sin^2 \theta_W$ is the weak mixing parameter, and the chiral couplings are $L_{u,d} = \pm 1 - 2x_W Q_{u,d}$ and $R_q = -2x_W Q_q$, with Q_q the electric charge of quark q . For neutrino-antiquark interactions, the coefficients can be obtained by crossing symmetry, *viz.* ($a \leftrightarrow b, xs \leftrightarrow Q^2 - xs$).

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FIGURES

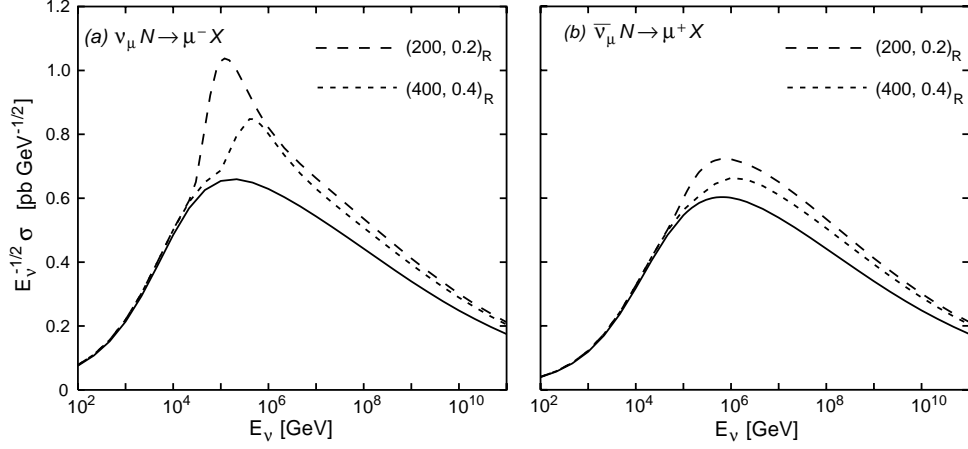


FIG. 1. Charged-current cross sections for (a) $\nu_\mu N$ and (b) $\bar{\nu}_\mu N$ interactions. The solid lines show the predictions of the standard model. The dashed (short-dashed) curves include the contributions of a right-handed squark, \tilde{d}_R^k , with mass $\tilde{m} = 200(400)$ GeV/ c^2 and \mathcal{R} coupling $\lambda'_{21k} = 0.2(0.4)$.

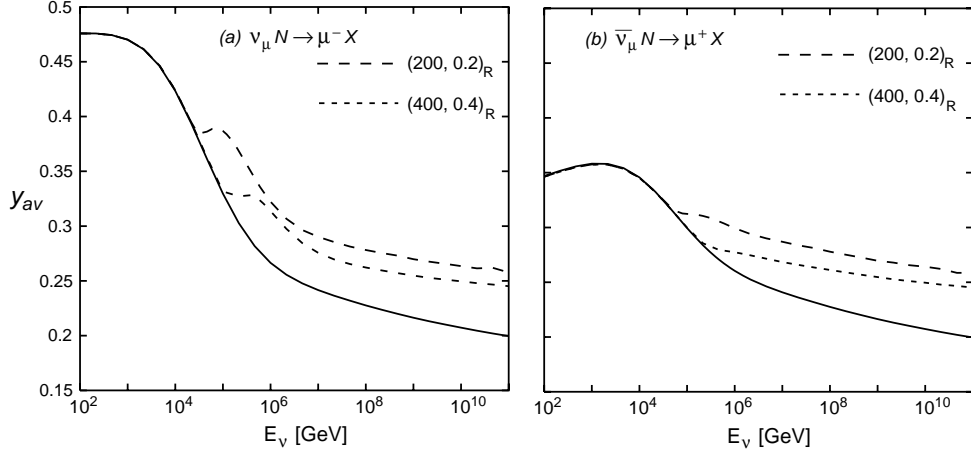


FIG. 2. Mean inelasticity parameter y_{av} for (a) $\nu_\mu N$ and (b) $\bar{\nu}_\mu N$ charged-current interactions. The solid lines show the predictions of the standard model. The dashed (short-dashed) curves include the contributions of a right-handed squark, \tilde{d}_R^k , with mass $\tilde{m} = 200(400)$ GeV/ c^2 and \mathcal{R} coupling $\lambda'_{21k} = 0.2(0.4)$.

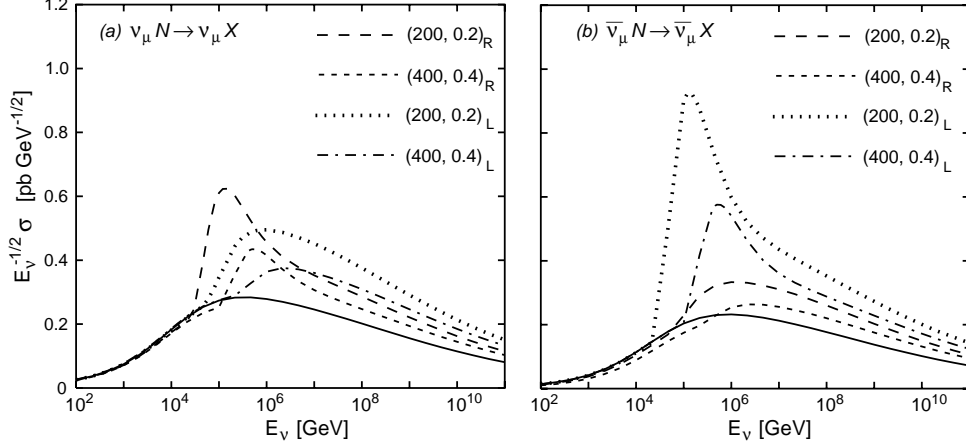


FIG. 3. Neutral-current cross sections for (a) $\nu_\mu N$ and (b) $\bar{\nu}_\mu N$ interactions. The solid lines show the predictions of the standard model. The dashed (short-dashed) curves include the contributions of a right-handed squark, \tilde{d}_R^k , with mass $\tilde{m} = 200(400)$ GeV/ c^2 and \mathcal{R} coupling $\lambda'_{21k} = 0.2(0.4)$. The dotted (dot-dashed) curves include the contributions of a left-handed squark, \tilde{d}_L^k , with mass $\tilde{m} = 200(400)$ GeV/ c^2 and \mathcal{R} coupling $\lambda'_{2k1} = 0.2(0.4)$.

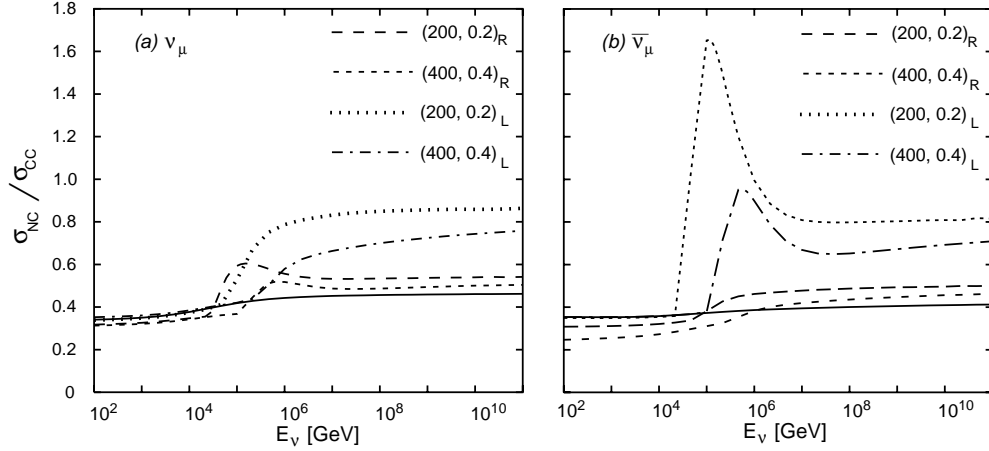


FIG. 4. Neutral-current to charged-current ratios for (a) $\nu_\mu N$ and (b) $\bar{\nu}_\mu N$ interactions. The solid lines show the predictions of the standard model. The dashed (short-dashed) curves include the contributions of a right-handed squark, \tilde{d}_R^k , with mass $\tilde{m} = 200(400)$ GeV/ c^2 and \mathcal{R} coupling $\lambda'_{21k} = 0.2(0.4)$. The dotted (dot-dashed) curves include the contributions of a left-handed squark, \tilde{d}_L^k , with mass $\tilde{m} = 200(400)$ GeV/ c^2 and \mathcal{R} coupling $\lambda'_{2k1} = 0.2(0.4)$.

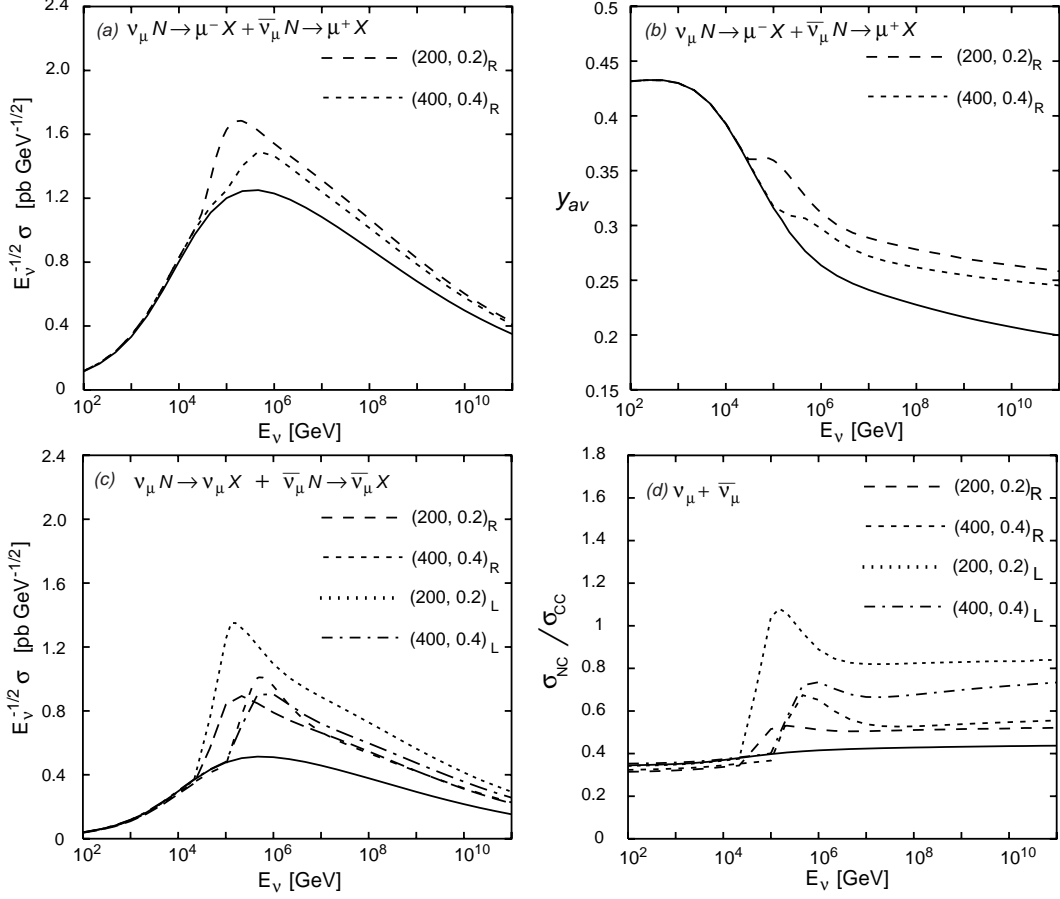


FIG. 5. Observables for $\nu N + \bar{\nu} N$ interactions. (a) charged-current cross section; (b) mean inelasticity y_{av} for charged-current interactions; (c) neutral-current cross section; and (d) neutral-current to charged-current ratio. The solid lines show the predictions of the standard model. The dashed (short-dashed) curves include the contributions of a right-handed squark, \tilde{d}_R^k , with mass $\tilde{m}_R^k = 200$ (400) GeV/c^2 and \mathcal{R} coupling $\lambda'_{21k} = 0.2$ (0.4). The dotted (dot-dashed) curves include the contributions of a left-handed squark, \tilde{d}_L^k , with mass $\tilde{m}_L^k = 200$ (400) GeV/c^2 and \mathcal{R} coupling $\lambda'_{2k1} = 0.2$ (0.4).

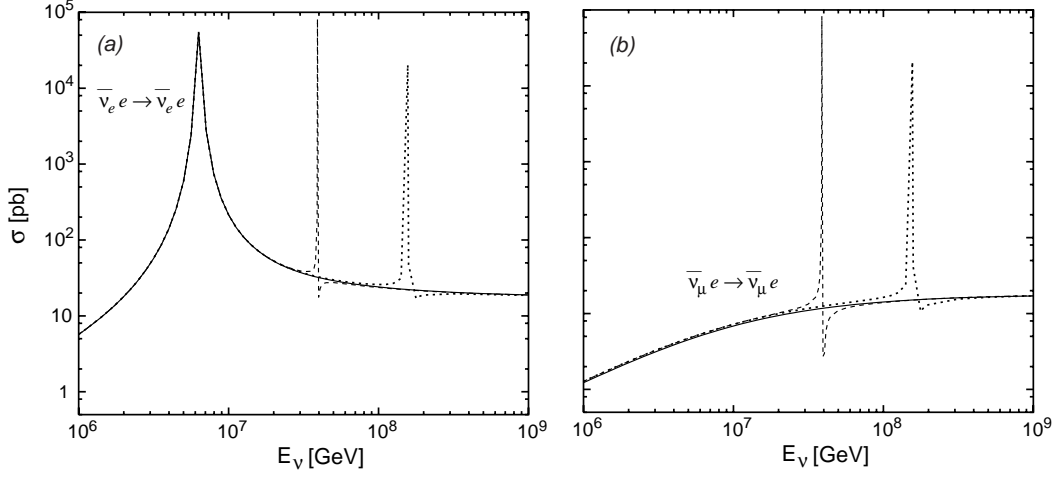


FIG. 6. (a) Cross section for $\bar{\nu}_e e$ elastic scattering in the standard model (solid line), showing the W^- resonance at 6.3 PeV, and for resonant production of a 200-GeV/ c^2 (short dashes) or 400-GeV/ c^2 (dotted line) slepton through an \mathcal{R} coupling. (b) $\bar{\nu}_\mu e \rightarrow \bar{\nu}_\mu e$ cross sections for resonant production of a 200-GeV/ c^2 (short dashes) or 400-GeV/ c^2 (dotted line) slepton through an \mathcal{R} coupling, compared with the featureless standard-model prediction (solid line). The \mathcal{R} couplings are $\lambda = (0.1, 0.2)$ for $m_{\tilde{e}_L^k} = (200, 400)$ GeV/ c^2 .

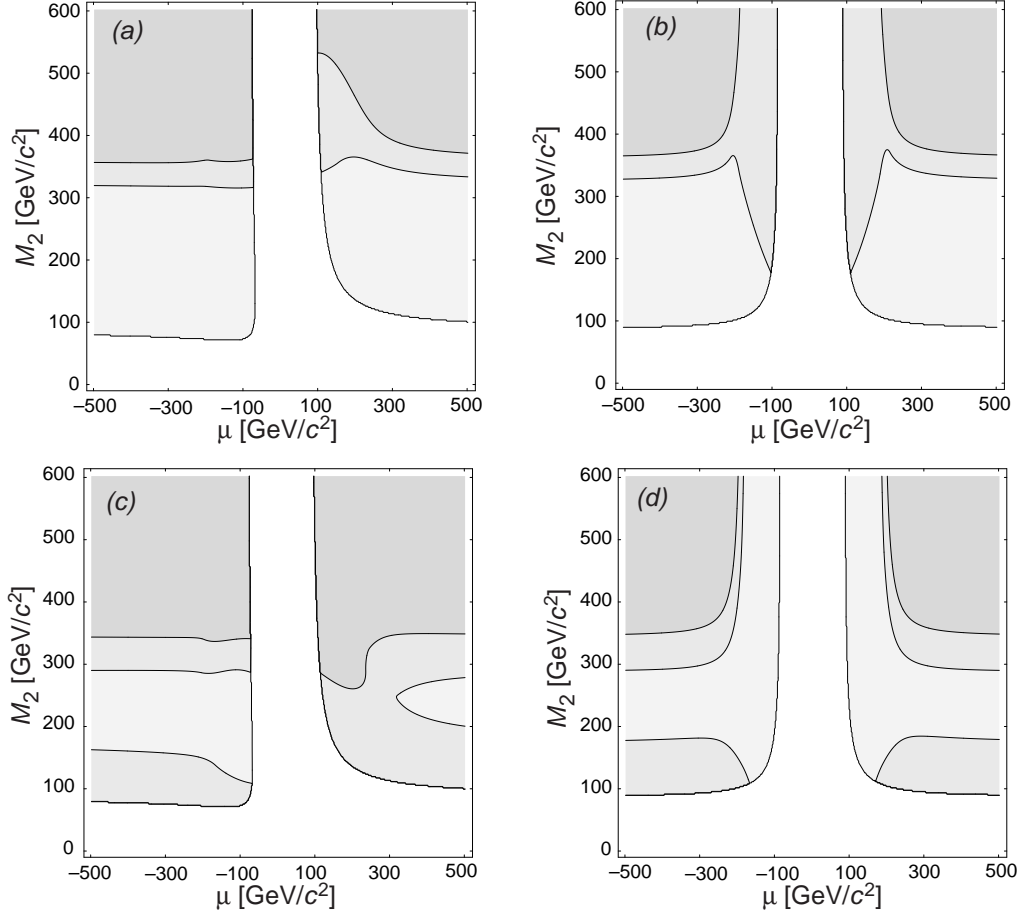


FIG. 7. Branching ratios for $\tilde{\tau}$ decays to $\tau +$ neutralinos for $m_{\tilde{\tau}_{R,L}} = 200 \text{ GeV}/c^2$ and $\lambda = 0.1$, as functions of parameters of the minimal supergravity model. The four cases correspond to: (a) $\tilde{\tau}_R$ production with $\tan \beta = 2$, (b) $\tilde{\tau}_R$ production with $\tan \beta = 60$, (c) $\tilde{\tau}_L$ production with $\tan \beta = 2$, and (d) $\tilde{\tau}_L$ production with $\tan \beta = 60$. The shaded regions denote branching fractions 0 – 0.3 (dark), 0.3 – 0.6 (medium), 0.6 – 1.0 (light). The U(1) gaugino mass M_1 is determined in terms of the SU(2) gaugino mass M_2 by the unification relation $M_1 = (5/3)M_2 \tan^2 \theta_W$. We exclude regions that correspond to chargino masses below the LEP2 bound of $85 \text{ GeV}/c^2$.

TABLES

TABLE I. Experimental constraints (at one or two standard deviations) on the R -parity-violating Yukawa couplings of interest, for the case of 200-GeV/ c^2 sfermions. For arbitrary sfermion mass, multiply the limits by $(m_{\tilde{f}}/200 \text{ GeV}/c^2)$, except for λ'_{221} .

\mathcal{R} Coupling	Limited by
$\lambda_{12k} < 0.1$ (2σ)	charged-current universality
$\lambda_{131,132,231} < 0.12$ (1σ)	$\Gamma(\tau \rightarrow e\nu\bar{\nu})/\Gamma(\tau \rightarrow \mu\nu\bar{\nu})$
$\lambda_{133} < 0.006$ (1σ)	ν_e Majorana mass
$\lambda'_{21k} < 0.18$ (1σ)	π decay
$\lambda'_{221} < 0.36$ (1σ)	D decay
$\lambda'_{231} < 0.44$ (2σ)	ν_μ deep inelastic scattering