PRAMANA
—journal of
physics

Vol. 51, Nos 1 & 2,
July/August 1998
pp. 123–133

An introduction to explicit R-parity violation

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Abstract. I discuss the theoretical motivations for R-parity violation, review the experimental bounds and outline the main changes in collider phenomenology compared to conserved R-parity.

Keywords. R-parity violation; collider phenomenology

PACS Nos 11.30; 12.50

1. Introduction

Until recently, R-parity violation ($\not{R}_p$) has been considered an unlikely component of the supersymmetric extension of the standard model (SM). In the past two years, it has motivated potentially favoured solutions to experimentally observed discrepancies (e.g. R_b, R_c, ALEPH four-jet events, HERA high Q^2 excess). It is the purpose of this chapter to present $\not{R}_p$ as an equally well motivated supersymmetric extension of the SM and provide an introductory guide. I start out with the definition of R_p and the most serious problem of proton decay. Then I discuss the various motivations for $\not{R}_p$, contrasting them with the R_p-conserving MSSM. Afterwards, I give an overview of the phenomenology of $\not{R}_p$.

2. What is R-parity?

R-parity (R_p) is a discrete multiplicative symmetry. It can be written as [1]

$$R_p = (-1)^{3B+L+2S}. \tag{1}$$

Here B denotes the baryon number, L the lepton number and S the spin of a particle. The electron has $R_p = +1$ and the selectron has $R_p = -1$. In fact, for all superfields of the supersymmetric SM, the SM field has $R_p = +1$ and its superpartner has[1] $R_p = -1$. R_p is conserved in the MSSM, superpartners can only be produced in pairs (all initial states at colliders are R_p even) and the LSP is stable. When extending the SM with supersymmetry one doubles the particle content to accomodate the superpartners and adds an additional Higgs doublet superfield. The minimal symmetries required to construct the Lagrangian are the gauge symmetry of the SM: $G_{\rm SM} = SU(3)_c \times SU(2)_L \times U(1)_Y$ and supersymmetry

[1] In general symmetries for which the anticommuting parameters, θ, transform non-trivially (and thus superpartners differently) are denoted R-symmetries. They can be discrete (R_p), global continuous, or even gauged [2, 3]. R-symmetries can be broken without supersymmetry being broken.

(including Lorentz invariance). The most general superpotential with these symmetries and this particle content (cf. Ch. 1) is [4]

$$W = W_{\rm MSSM} + W_{\not R_p}, \tag{2}$$

$$W_{\rm MSSM} = h^e_{ij} L_i H_1 \bar{E}_j + h^d_{ij} Q_i H_1 \bar{D}_j + h^u_{ij} Q_i H_2 \bar{U}_j + \mu H_1 H_2, \tag{3}$$

$$W_{\not R_p} = \tfrac{1}{2}\lambda_{ijk} L_i L_j \bar{E}_k + \lambda'_{ijk} L_i Q_j \bar{D}_k + \tfrac{1}{2}\lambda''_{ijk} \bar{U}_i \bar{D}_j \bar{D}_k + \kappa_i L_i H_2. \tag{4}$$

$i, j = 1,2,3$ are generation indices and a summation is implied. L_i (Q_i) are the lepton (quark) $SU(2)_L$ doublet superfields. $\bar{E}_j$ ($\bar{D}_j$, $\bar{U}_j$) are the electron (down- and up-quark) $SU(2)_L$ singlet superfields. λ, λ', and λ'' are Yukawa couplings. The κ_i are dimensionful mass parameters. The $SU(2)_L$ and $SU(3)_C$ indices have been suppressed. When including them we see that the first term in $W_{\not R_p}$ is anti-symmetric in $\{i,j\}$ and the third term is anti-symmetric in $\{j,k\}$. Therefore $i \neq j$ in $L_i L_j \bar{E}_k$ and $j \neq k$ in $\bar{U}_i \bar{D}_j \bar{D}_k$. Equation (4) thus contains $9 + 27 + 9 + 3 = 48$ new terms beyond those of the MSSM.

The last term in eq. (4), $L_i H_2$, mixes the lepton and the Higgs superfields. In supersymmetry L_i and H_1 have the same gauge and Lorentz quantum numbers and we can redefine them by a rotation in (H_1, L_i). The terms $\kappa_i L_i H_2$ can then be rotated to zero in the superpotential [5]. If the corresponding soft supersymmetry breaking parameters B_i are aligned with the κ_i they are simultaneously rotated away [5, 6]. However, the alignment of the superpotential terms with the soft breaking terms is not stable under the renormalization group equations [7]. Assuming an alignment at the unification scale, the resulting effects are small [7] except for neutrino masses [7, 8]. The effects can be further suppressed by a horizontal symmetry. Throughout the rest of this chapter, I will assume the $L_i H_2$ terms have been rotated away[2]

$$W_{\not R_p} = \lambda_{ijk} L_i L_j \bar{E}_k + \lambda'_{ijk} L_i Q_j \bar{D}_k + \lambda''_{ijk} \bar{U}_i \bar{D}_j \bar{D}_k. \tag{5}$$

Expanding for example the $LL\bar{E}$ term into the Yukawa couplings yields

$$\mathcal{L}_{LL\bar{E}} = \lambda_{ijk}[\tilde{\nu}^i_L \bar{e}^k_R e^j_L + \tilde{e}^j_L \bar{e}^k_R \nu^i_L + (\tilde{e}^k_R)^* (\bar{\nu}^i_L)^c e^j_L - (i \leftrightarrow j)] + \text{h.c.} \tag{6}$$

The tilde denotes the scalar fermion superpartners. These terms thus violate lepton-number. The $LQ\bar{D}$ terms also violate lepton number and the $\bar{U}\bar{D}\bar{D}$ terms violate baryon number. The entire superpotential (5) violates R_p.

3. Proton decay and discrete symmetries

The combination of lepton- and baryon-number violating operators in the Lagrangian can possibly lead to rapid proton decay. For example the two operators $L_1 Q_1 \bar{D}_k$ and $\bar{U}_1 \bar{D}_1 \bar{D}_k$ ($k \neq 1$) can contribute to proton decay at tree-level via s-channel squark exchange. On dimensional grounds we estimate

$$\Gamma(P \rightarrow e^+ \pi^0) \approx \frac{\alpha(\lambda'_{11k})\alpha(\lambda''_{11k})}{\tilde{m}^4_{dk}} M^5_{\rm proton}. \tag{7}$$

[2] The ambiguity on bounds due to rotations in (L_i, H_1) space has been discussed in [9].

Here $\alpha(\lambda) = \lambda^2/(4\pi)$. Given that [10] $\tau(P \to e\pi) > 10^{32}$ yr, we obtain

$$\lambda'_{11k} \cdot \lambda''_{11k} \lesssim 2 \cdot 10^{-27} \left(\frac{\tilde{m}_{dk}}{100\,\mathrm{GeV}} \right)^2 . \tag{8}$$

For a more detailed calculation see [11]. This bound is so strict that the only natural explanation is for at least one of the couplings to be zero. Thus the simplest supersymmetric extension of the SM is excluded: an extra symmetry is required to protect the proton.

In the MSSM, R_p is imposed by hand. This forbids all the terms in $W_{\not{R}_p}$ and thus protects the proton. An alternative discrete symmetry with the same physical result is matter parity

$$(L_i, \bar{E}_i, Q_i, \bar{U}_i, \bar{D}_i) \to -(L_i, \bar{E}_i, Q_i, \bar{U}_i, \bar{D}_i), \quad (H_1, H_2) \to (H_1, H_2). \tag{9}$$

This forbids all terms with an odd power of matter fields and thus forbids all the terms in $W_{\not{R}_p}$. However, there are other solutions, which protect the proton equally well. If baryon number is conserved the proton can not decay. Thus forbidding just the interactions $\bar{U}_i\bar{D}_j\bar{D}_k$ is sufficient. This can be achieved by baryon-parity

$$(Q_i, \bar{U}_i, \bar{D}_i) \to -(Q_i, \bar{U}_i, \bar{D}_i), \quad (L_i, \bar{E}_i, H_1, H_2) \to (L_i, \bar{E}_i, H_1, H_2). \tag{10}$$

This symmetry thus protects the proton but allows for $\not{R}_p$ via the $L_iL_j\bar{E}_k$ and $L_iQ_j\bar{D}_k$ operators. If only the interactions $\bar{U}_i\bar{D}_j\bar{D}_k$ are allowed *and* the proton is lighter than the LSP the proton is stable as well. This can be achieved by lepton parity

$$(L_i, \bar{E}_i) \to -(L_i, \bar{E}_i), \quad (Q_i, \bar{U}_i, \bar{D}_i, H_1, H_2) \to (Q_i, \bar{U}_i, \bar{D}_i, H_1, H_2). \tag{11}$$

Baryon-parity and lepton parity are two possible solutions to maintain a stable proton and allow for $\not{R}_p$. There is a large number of discrete symmetries which can achieve this [12].

4. Motivation

The symmetries discussed in the previous section were all imposed *ad hoc* with no deeper motivation than to ensure the stability of the proton. On this purely phenomenological level there is no reason to prefer the models with conserved R_p versus those with $\not{R}_p$. However, this is not a satisfactory view of the weak-scale picture. Hopefully, the correct structure will emerge from a simpler theory at a higher energy predicting either R_p-conservation or $\not{R}_p$.

Grand unified theories: In GUTs quarks and leptons are typically in common multiplets and thus have the same quantum numbers. The discrete symmetries protecting the proton and resulting in $\not{R}_p$ typically treat quarks and leptons differently and thus seem incompatible with a GUT. All the same, several GUT models have been constructed [5, 13–15] which have low-energy $\not{R}_p$. This is typically achieved by non-renormalizable GUT scale operators involving Higgs fields. These operators become renormalizable $\not{R}_p$-operators after the GUT symmetry has been broken. Such models have been constructed for the GUT gauge groups $SU(5)$ [5, 14], $SO(10)$ [14], and $SU(5) \times U(1)$ [13, 14]. They have been constructed such that the only set of low-energy operators is $LL\bar{E}$ or $LQ\bar{D}$, or $\bar{U}\bar{D}\bar{D}$, respectively. There is thus no problem with proton decay. In order to ensure that only the required set of non-renormalizable operators are allowed, additional symmetries are required beyond the GUT gauge group. This is true for both R_p-conservation and $\not{R}_p$. Thus from a grand unified point of view there is no preference for either R_p-conservation or $\not{R}_p$.

String theory: In string theories unification can be achieved without a simple gauge group. There is thus no difficulty in having distinct quantum numbers for quarks and lepton superfields. Indeed R_p-conserving and $\not{R}_p$ string theories have been constructed [16]. At present, there does not seem to be a preference at the string level for either of the two.

In both string theory and in GUTs, there is no generic prediction for the size of the $\not{R}_p$-Yukawa couplings. This is analogous to the fermion mass problem.

Discrete gauge symmetries: There has been a further attack on this problem from a slightly different angle. If a discrete symmetry is a remnant of a broken gauge symmetry it is called a discrete gauge symmetry. It has been argued that quantum gravity effects maximally violate all discrete symmetries unless they are discrete gauge symmetries [17]. The condition that the underlying gauge symmetry be anomaly-free can be translated into conditions on the discrete symmetry. A systematic analysis of all $\mathcal{Z}_N$ symmetries [18] has been performed. The result was that only two symmetries were discrete gauge anomaly-free: R_p and baryon-parity (10). Baryon-parity was slightly favoured since in addition it prohibited dimension-5 proton-decay operators. It has since been shown [19] that the non-linear constraints in [18] are model dependent thus possibly allowing an even larger set of discrete symmetries.

Given the quantum gravity argument it is more appealing to determine the low-energy structure directly from gauge symmetries instead of discrete symmetries. This can possibly even be connected with the fermion-mass or flavour problem. This is an on-going field of research and it is too early to draw any conclusions. I just point out that gauged models with $\not{R}_p$ have been constructed [3, 6].

In conclusion, from the theoretical understanding of unification, there is no clear preference between R_p and $\not{R}_p$. In light of the very distinct phenomenology which we discuss below, it is thus mandatory to experimentally search for both possibilities. R_p-conservation and $\not{R}_p$ have the same minimal particle content. They also in principle have the same kind of symmetries, as we have just argued: $G_{\rm SM}$ plus an additional symmetry to protect the proton. They should thus both be considered as different versions of the MSSM. We shall denote the R_p conserving version of the MSSM as R_p-MSSM and the R_p-violating version as $\not{R}_p$-MSSM.

5. Indirect bounds

The $\not{R}_p$ interactions can contribute to various (low-energy) processes through the virtual exchange of supersymmetric particles [20]. To date, all data are in good agreement with the SM. This leads directly to bounds on the $\not{R}_p$ operators. When determining such limits one must make some simplifying assumptions due to the large number of operators. In the following, we shall assume that one $\not{R}_p$ operator at a time is dominant while the others are negligible. We thus do not include the sometimes very strict bounds on products of operators, for example from $\mu \to e\gamma$ [21]. This is an important assumption but not unreasonable. It holds for the SM for example, where the top quark Yukawa coupling is almost a factor 40 larger than the bottom Yukawa coupling. Since we do not know the origins of Yukawa couplings, we do not know whether this is a generic feature.

Before presenting the complete bounds, I shall discuss one example [22] to show how such bounds can be obtained. The operator $L_1 L_3 \bar{E}_k$ can contribute to the decay $\tau \to e\nu\bar{\nu}$ at tree-level via slepton exchange. For large slepton masses, $\tilde{m}(\tilde{e}^k_R)$, this interaction is described by an effective 4-fermion Lagrangian (after Fierz re-ordering) [22]

$$\mathcal{L}_{\rm eff} = \frac{|\lambda_{13k}|^2}{2\tilde{m}^2}(\bar{e}_L \gamma^\mu \nu_{eL})(\bar{\nu}_{\tau L}\gamma_\mu \tau_L). \tag{12}$$

This has the same structure as the term in the effective SM Lagrangian and thus leads to an apparent shift in the Fermi constant for tau decays. Considering the ratio $R_\tau \equiv \Gamma(\tau \to e\nu\bar{\nu})/\Gamma(\tau \to \mu\nu\bar{\nu})$, the contribution from $R\!\!\!/_p$ relative to the SM contribution is [22]

$$R_\tau = R_\tau(\mathrm{SM})\left[1 + 2\frac{M_W^2}{g^2}\left(\frac{|\lambda_{13k}|^2}{\tilde{m}^2(\tilde{e}^k_R)}\right)\right]. \tag{13}$$

Using the experimental value [10] $R_\tau / R_\tau(\mathrm{SM}) = 0.99889 \pm 0.0071$ we obtain the bounds

$$|\lambda_{13k}| < 0.04\left(\frac{\tilde{m}(\tilde{e}^k_R)}{100\,\mathrm{GeV}}\right), \quad (1\sigma) \quad k = 1,2,3, \tag{14}$$

which are given in table 1. The strictest bounds on the remaining operators are also summarized in table 1.

Table 1. Strictest bounds on $R\!\!\!/_p$ Yukawa couplings for $\tilde{m} = 100\,\mathrm{GeV}$. The physical processes from which they are obtained are summarized in the main text.

ijk	λ_{ijk}	ijk	λ'_{ijk}	ijk	λ'_{ijk}	ijk	λ'_{ijk}	ijk	λ''_{ijk}
121	$0.05^{a\dagger}$	111	0.001^{d}	211	0.09^{h}	311	0.16^{k}	112	$10^{-6,\ell}$
122	$0.05^{a\dagger}$	112	$0.02^{a\dagger}$	212	0.09^{h}	312	0.16^{k}	113	$10^{-5,m}$
123	$0.05^{a\dagger}$	113	$0.02^{a\dagger}$	213	0.09^{h}	313	0.16^{k}	123	1.25^{**}
131	0.04^{b}	121	$0.035^{e\dagger}$	221	0.18^{i}	321	0.20^{f*}	212	1.25^{**}
132	0.04^{b}	122	0.06^{c}	222	0.18^{i}	322	0.20^{f*}	213	1.25^{**}
133	0.004^{c}	123	0.20^{f*}	223	0.18^{i}	323	0.20^{f*}	223	1.25^{**}
231	0.05^{b}	131	$0.035^{e\dagger}$	231	$0.22^{j\dagger}$	331	0.26^{g}	312	0.43^{g}
232	0.05^{b}	132	0.33^{g}	232	0.39^{g}	332	0.26^{g}	313	0.43^{g}
233	0.05^{b}	133	0.002^{c}	233	0.39^{g}	333	0.26^{g}	323	0.43^{g}

The bounds in table 1 are obtained from the following physical processes: [a]charged current universality [22, 10], [b]$\Gamma(\tau \to e\nu\bar{\nu})/\Gamma(\tau \to \mu\nu\bar{\nu})$ [22, 10], [c]bound on the mass of ν_e [5, 23, 24], [d]neutrinoless double-beta decay [25, 26], [e]atomic parity violation [27–29], [f]$D^0 - \bar{D}^0$ mixing [30–32], [g]$R_\ell = \Gamma_{\rm had}(Z^0)/\Gamma_\ell(Z^0)$ [33, 20], [h]$\Gamma(\pi \to e\bar{\nu})/\Gamma(\pi \to \mu\bar{\nu})$ [26], [i]$BR(D^+ \to \bar{K}^{0*}\mu^+\nu_\mu)/BR(D^+ \to \bar{K}^{0*}e^+\nu_\mu)$ [35, 20], [j]ν_μ deep-inelastic scattering [22], [k]$BR(\tau \to \pi\nu_\tau)$ [35, 20], [ℓ]heavy nucleon decay [11], and [m]$n - \bar{n}$ oscillations [34, 11].

The bounds denoted by † are 2σ bounds, the other bounds are at the 1 sigma level. The bounds denoted by (**) are not direct experimental bounds.

The bounds denoted by * are based on a further *assumption* about the absolute mixing in the (SM) quark sector.

The bounds are all given for $\tilde{m} = 100\,\mathrm{GeV}$ and they become weaker with increasing $\tilde{m}$. They each depend on a specific scalar mass and have various functional dependences on this mass.

λ_{ijk}	$m_{\tilde{e}_{Rk}}/100\,\mathrm{GeV}$	λ'_{111}	$(m_{\tilde{q}}/100\,\mathrm{GeV})^2(m_{\tilde{g}}/1\,\mathrm{TeV})^{1/2}$
$\lambda'_{11k}, \lambda'_{21k}$	$m_{\tilde{d}_{Rk}}/100\,\mathrm{GeV}$	λ'_{1j1}	$m_{\tilde{q}_{Lj}}/100\,\mathrm{GeV}$
$\lambda_{133}, \lambda'_{1jj}$	$\sqrt{m_{\tilde{\tau},\tilde{d}_j}/100\,\mathrm{GeV}}$	λ'_{123}	$\sqrt{m_{\tilde{b}_R}/100\,\mathrm{GeV}}$
λ'_{231}	$m_{\tilde{\nu}_{\tau L}}/100\,\mathrm{GeV}$	λ'_{32k}	$\sqrt{m_{\tilde{d}_{Rk}}/100\,\mathrm{GeV}}$

(15)

For λ'_{111} the dependence can be on either $m_{\tilde{d}_{Rk}}$, or $m_{\tilde{u}_{Lk}}$. The bound in table 1 is given for $\tilde{m}(\tilde{g}) = 1\,\mathrm{TeV}$. For λ'_{132}, λ'_{22k}, λ'_{23k}, λ'_{31k}, and λ'_{32k} one must consult the appropriate references since the dependence is only given numerically. The bounds on $\lambda''_{112,113}$ from heavy nucleon decay and $n-\bar{n}$ oscillations have very strong mass dependences [11].

I have updated the previous bounds from charged current universality [22], from lepton-universality [22] and from[3] R_ℓ using more recent data [10]. For the bound from the electron neutrino mass I have used the upper bound [24] $m_{\nu_e} < 5\,\mathrm{eV}$. The PDG number is 10–15 eV [10] and is very conservative [36]. The bound on λ from m_{ν_e} scales with the square root of the upper bound on m_{ν_e}. For the bound from atomic parity violation I have used the theory value [29]: $Q_W^{\mathrm{th}} = -73.17 \pm 0.13$. The error includes the variations due to the unknown Higgs mass. I have also used the recent new experimental number [28]. For the bound from $D^0-\bar{D}^0$ mixing, I have updated the bound from [32] to include a lattice calculation of [31] B_D and a more updated value of f_D [30]. I have also included a 10% error to account for the quenched approximation. The ** bounds are obtained [37, 11] from the requirement that the $R\!\!\!/_p$-coupling remains within the unitarity bound up to the grand unified scale of 10^{16} GeV. This need not be the case.

As stated before, we do not know the physical origin of Yukawa couplings or superpotential terms. It is a reasonable (but not necessary) assumption that their structure is determined by some symmetry at an energy scale well above the electroweak scale, e.g. the GUT or the Planck scale. We would expect this symmetry to be in terms of the weak *current* eigenstates. Such a symmetry could then give us a single dominant operator, for example

$$L_1 Q_1 \bar{D}_1 = \lambda'_{111}(-\tilde{e}_L u_L \bar{d}_R + \tilde{\nu}_{eL} d_L \bar{d}_R \ldots). \tag{16}$$

Below the electroweak scale the quarks become massive and we must rotate them to their *mass* eigenstate basis. (The squarks must separately also be rotated by a different rotation but that is not relevant to these bounds.) In (16) there are then separate rotations: $d_L \rightarrow \mathcal{D}_{1j} d'_{jL}$ and $u_L \rightarrow \mathcal{U}_{1j} u'_{jL}$ which generate extra $R\!\!\!/_p$ terms suppressed by mixing angles [38].

For the quarks we do not know the *absolute* mixing of the down-quark sector, $\mathcal{D}_{ij}$, or of the up-quark sector $\mathcal{U}_{ij}$ and thus do not know by how much to rotate the up- and down-quark current eigenstates. The *relative* mixing of these two sectors is given by the

[3] I thank Gautam Bhattacharyya for providing me with updates on the bounds resulting from R_ℓ.

CKM-matrix [10] of the SM. If we assume the relative rotation is solely due to an absolute mixing in the up-quark sector ($\mathcal{D}_{ij}=1$) the best bounds are those given in table 1. Those denoted by * are specifically based on this mixing assumption. If however the relative mixing is solely due to absolute mixing in the down-quark sector ($\mathcal{U}_{ij}=1$) the $D^0-\bar{D}^0$ mixing bounds no longer apply. There are then significantly stricter bounds on many couplings from measurements of $K^+\to\pi^+\nu\nu$ decays [32]

$$\lambda'_{ijk}<0.012,\ (90\%CL),\quad j\neq 3. \tag{17}$$

For table 1 we have adopted the conservative estimate that the mixing is solely due to the up-quark sector since we do not know the absolute mixing. We therefore did not include the bounds (17).[4]

6. Changes to R_p-MSSM

On the Lagrangian level the only change to the R_p-MSSM is the inclusion of the operators in $W_{\not R_p}$ which give new lepton- and baryon number violating Yukawa couplings. There are several changes in the phenomenology of supersymmetry due to these couplings [38].

1. Lepton- or baryon-number is violated as discussed in § 5.
2. The LSP is not stable and can decay in the detector. It is no longer a dark matter candidate.
3. The neutralino is not necessarily the LSP.
4. The single production of supersymmetric particles is possible.

2. If for example the neutralino is the LSP and the dominant $\not R_p$ operator is $L_1Q_2\bar{D}_1$ it can decay to an electron and two quarks at tree-level. For LSP$=\tilde\gamma$ the decay rate is [39, 40]

$$\Gamma_{\tilde\gamma}=\frac{3\alpha\lambda'^2_{121}}{128\pi^2}\frac{M^5_{\chi^0_1}}{\tilde m^4}. \tag{18}$$

The decay occurs in the detector if $c\gamma_L\tau(\tilde\gamma)\lesssim 1\,m$, or

$$\lambda'_{121}>1.4\cdot 10^{-6}\sqrt{\gamma_L}\left(\frac{\tilde m}{200\,\mathrm{GeV}}\right)^2\left(\frac{100\,\mathrm{GeV}}{M_{\tilde\gamma}}\right)^{5/2}. \tag{19}$$

where γ_L is the Lorentz boost factor. This is well below the bound of table 1. Recall also for comparison, that in the SM Yukawa couplings can be very small: for the electron $h^e=3\cdot 10^{-6}$. We have presented these numerical results for a photino for simplicity and clarity. The full analysis with a neutralino LSP has been performed in [41, 42]. It involves several subtleties due to the R_p-MSSM parameter space which can have significant effects on the lifetime. Due to the LSP decay, supersymmetry with broken R_p has no natural dark matter candidate.

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[4] There is a possible loop-hole. The symmetry at the high energy scale could just produce such a combination of couplings that is rotated to one single dominant coupling at low energy. After all, it is possibly the same symmetry which produces the single dominant quark Yukawa coupling in the SM. However, I do not adopt this philosophy here.

3. In the R_p-MSSM the stable LSP must be charge and colour neutral for cosmological reasons (cf. Chapter 15). In the $\not{R}_p$-MSSM there is no preference for the nature of the unstable LSP. It can be any of the following[5]

$$\mathrm{LSP} \in \{\chi_1^0, \chi_1^\pm, \tilde{g}, \tilde{q}, \tilde{t}, \tilde{\ell}, \tilde{\nu}\}. \tag{20}$$

In each case the collider phenomenology can be quite distinct.

4. In the $\not{R}_p$-MSSM there are resonant and non-resonant single particle production mechanisms. The resonant production mechanisms are

$$e^+ + e^- \to \tilde{\nu}_{Lj}, \quad L_1 L_j \bar{E}_1, \tag{21}$$

$$e^- + u_j \to \tilde{d}_{Rk}, \quad L_1 Q_j \bar{D}_k, \tag{22}$$

$$e^- + \bar{d}_k \to \bar{\tilde{u}}_{Lj}, \quad L_1 Q_j \bar{D}_k, \tag{23}$$

$$\bar{u}_j + d_k \to \tilde{e}^-_{Li}, \quad L_i Q_j \bar{D}_k, \tag{24}$$

$$d_j + \bar{d}_k \to \tilde{\bar{\nu}}_{Li}, \quad L_1 Q_j \bar{D}_k, \tag{25}$$

$$\bar{u}_i + \bar{d}_j \to \tilde{d}_{Rk}, \quad \bar{U}_i \bar{D}_j \bar{D}_k, \tag{26}$$

$$d_j + d_k \to \bar{\tilde{u}}_{Ri}, \quad \bar{U}_i \bar{D}_j \bar{D}_k. \tag{27}$$

These processes can be realized at e^+e^--colliders, at HERA, and at hadron colliders, respectively. There are many further t-channel single sparticle production processes. For example at an e^+e^--collider, we can have $e^+ + e^- \to \tilde{\chi}_1^0 + \nu_j$ via t-channel selectron exchange. The t-channel exchange of squarks (sleptons) can also contribute to $q\bar{q}$ ($\ell\bar{\ell}$) pair production, leading to indirect bounds [44].

7. Collider phenomenology

The supersymmetric signals for $\not{R}_p$ will be a combination of supersymmetric production and decay to R_p even final states. Supersymmetric particles can be produced in pairs via MSSM gauge couplings or singly as in (21)–(27). The former benefits from large couplings while being kinematically restricted to masses $< \sqrt{s}/2$. The latter case has double the kinematic reach but suffers from typically small Yukawa couplings. Combining the various production modes with the decays and the different dominant operators leads to a wide range of potential signals to search for. Instead of systematically listing them I shall focus on two examples. Throughout we shall assume a neutralino LSP.

7.1 *Squark pair production at the Tevatron*

Squark pair production at the Tevatron proceeds via the known gauge couplings of the R_p-MSSM

$$q\bar{q}, gg \to \tilde{q} + \bar{\tilde{q}}. \tag{28}$$

In $\not{R}_p$, once produced the squarks decay to an R_p even final state. Let us consider a dominant $L_i L_j \bar{E}_k$ operator. The couplings λ_{ijk} are bounded to be smaller than gauge

[5] The stop is listed separately since it has a special theoretical motivation [43] and leads to quite distinct phenomenology given that the top quark is so heavy.

Table 2. Squark mass bounds from the Tevatron for various dominant $\not{R}_p$-operators [45].

	$L_{1,2}Q_{1,2}\bar{D}_k$,	$L_{1,2}L_3\bar{E}_3$	$L_1L_2\bar{E}_3$	$L_{1,2}L_3\bar{E}_{1,2}$	$L_1L_2\bar{E}_{1,2}$
$m_{\tilde{q}}$	100 GeV	100 GeV	140 GeV	160 GeV	175 GeV

couplings. Thus we expect the squarks to cascade decay to LSPs as in the MSSM. The LSPs in turn will then decay via the operator $L_iL_j\bar{E}_k$ to two charged leptons and a neutrino each. If each squark decays directly to the LSP (assuming it is the second lightest)

$$q'\bar{q}', gg \to \tilde{q} + \bar{\tilde{q}} \to q\bar{q} + \tilde{\chi}_1^0\tilde{\chi}_1^0 \to q\bar{q} + l^+l^-l^+l^-\nu\nu. \tag{29}$$

We therefore have a multi-lepton signal which is detectable [45]. To date it has not been searched for with $\not{R}_p$ in mind. However, before the top quark discovery there was a bound from CDF on a di-lepton production cross section. Making corresponding cuts and with some simple assumptions this can be translated into a bound on the rate of the process (29) and thus a lower bound on the squark mass [45]. The assumptions are: (i) $BR(\tilde{q} \to \tilde{\gamma}q) = 100\%$, ($m_{\tilde{q}} < m_{\tilde{g}}$), (ii) LSP$=\tilde{\gamma}$ with $M_{\tilde{\gamma}} = 30\,\mathrm{GeV}$, (iii) λ, λ' satisfy the bound (19). For various dominant operators the bounds are given in table 2. No attempt was made to consider final state τ's due to lack of data. These bounds are comparable to the R_p-MSSM squark mass bounds. Since, the theoretical analysis has been improved to allow for neutralino LSPs, more involved cascade decays and the operator $\bar{U}\bar{D}\bar{D}$ [38, 46]. However, to date no experimental analysis has been performed.

7.2 *Resonant squark production at HERA*

HERA offers the possibility to test the operators $L_1Q_j\bar{D}_k$ via resonant squark production [47][6]

$$e^+ + d_k \to \tilde{u}_j \to (e^+ + d_k, \tilde{\chi}_1^0 + u_j, \tilde{\chi}_1^+ + d_j), \tag{30}$$

$$e^+ + \bar{u}_j \to \bar{\tilde{d}}_k \to (e^+ + \bar{u}_j, \bar{\nu}_e + \bar{d}_j, \tilde{\chi}_1^0 + \bar{d}_k). \tag{31}$$

We have included what are most likely the dominant decay modes. The neutralino and chargino will decay as

$$\tilde{\chi}_1^0 \to (e^\pm, \nu) + 2\,\mathrm{jets}, \quad \tilde{\chi}_1^+ \to (e^+, \nu) + 2\,\mathrm{jets}. \tag{32}$$

The neutralino can decay to the electron or positron since it is a Majorana fermion. We are thus left with several distinct decay topologies. (i) If the squark is the LSP it will decay to $e^+ + q$ or $\bar{\nu}_e + q$ ($\bar{\tilde{d}}_k$). The first looks just like neutral current DIS, except that for $x_{Bj} \approx \tilde{m}^2(\tilde{q})/s$ it results in a flat distribution in y_e whereas NC-DIS gives a $1/y_e^2$ distribution. The latter looks just like CC-DIS. (ii) If the gauginos are lighter than the squark the gaugino decay will dominate [42][7]. The clearest signal is a high p_T electron which is

[6] HERA has accumulated most of its data as a positron proton collider.
[7] The gaugino decays could be suppressed by phase space or by partial cancellations of the neutralino couplings [41, 48].

Table 3. Exclusion upper limits at 95% CL on λ'_{1jk} for $\tilde{m}(\tilde{q}) = 150\,\mathrm{GeV}$ and $\tilde{m}(\tilde{\chi}_1^0) = 40\,\mathrm{GeV}$ for two different dominant admixtures of the neutralino.

	λ'_{111}	λ'_{112}	λ'_{113}	λ'_{121}	λ'_{122}	λ'_{123}	λ'_{131}	λ'_{132}	λ'_{133}
$\tilde{\gamma}$-like	0.056	0.14	0.18	0.058	0.19	0.30	0.06	0.22	0.55
$\tilde{Z}^0$-like	0.048	0.12	0.15	0.048	0.16	0.26	0.05	0.19	0.48

essentially background free. The high p_T positron or the missing p_T of the neutrino can also be searched for.

All five signals have been searched for by the H1 collaboration [49] in the 1994 e^+ data ($\mathcal{L} = 2.83\,\mathrm{pb}^{-1}$). The observations were in excellent agreement with the SM. The resulting bounds on the couplings are summarized in table 3. After rescaling the bounds of table 1 we see that the direct search is an improvement for λ'_{121}, λ'_{131}, and λ'_{132}.

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