

R-Parity Violation and Unification

G.F. Giudice¹ and R. Rattazzi

Theory Division, CERN, CH-1211, Genève 23, Switzerland

Abstract

The reported anomaly in deep-inelastic scattering at HERA has revived interest in the phenomenology of *R*-parity violation. From the theoretical point of view, the existence of *R*-violating interactions poses two considerable problems. The first one concerns the flavour structure of the interactions and the origin of an appropriate suppression of flavour-changing neutral-current processes and lepton-family transitions. The second one concerns the way of embedding *R*-violating interactions in a grand unified theory (GUT) without introducing unacceptable nucleon decay rates. We show that the second problem can be solved by a mechanism which is purely group theoretical and does not rely on details of the flavour theory. We construct explicit GUT models in which our mechanism can be realized.

CERN-TH/97-76
April 1997

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

1. The anomaly in deep-inelastic e^+p scattering events reported by H1 [1] and ZEUS [2] and the excess of four-jet events observed by ALEPH [3], but not confirmed by the other LEP experiments [4], have revived interest [5, 6] in the phenomenology of R -parity violating interactions. Certainly more statistics is required to understand if experiments are really observing some signal of new physics. Nevertheless we believe it is timely and important to investigate what are the consequences of R -parity violation in our understanding of the theoretical framework of supersymmetric models.

The two main problems arising from R -parity violation are connected with *flavour* and with *unification*. As we emphasize below, we believe that these two problems are quite different both in their quantitative aspect (the unification problem being numerically more acute) and in their conceptual aspect. In this paper we will present a solution to the unification problem which is independent of the flavour structure of the theory, and relies only on GUT symmetry properties. At present, the question of flavour remains unresolved, since we do not know any fully convincing theory which generates the peculiar hierarchy of the Yukawa couplings – let alone the R -parity violating couplings. We believe that separating the two puzzles, and solving the unification problem in terms of GUTs, leads to important theoretical progress since, as we argue below, it is probably unrealistic to hope that an ultimate flavour theory can provide the right cure.

2. We concentrate on the R -violating interaction suggested to explain the large- Q^2 data at HERA, which is given by the term in the superpotential

$$W = \lambda_{ijk} Q_L^i \bar{D}_R^j L_L^k \quad (1)$$

with $\lambda_{i11} \gtrsim 4 \times 10^{-2}$ [5] and i equal to 2 or 3. Here i, j, k refer to generation indices and we employ a standard notation for quark and lepton superfields. The *flavour* problem arises because the generation structure of the operator in eq. (1) is in general not aligned with the generation structure of the Yukawa interactions

$$W = h_{ij}^d Q_L^i \bar{D}_R^j \bar{H} + h_{ij}^u Q_L^i \bar{U}_R^j H + h_{ij}^e L_L^i \bar{E}_R^j \bar{H}. \quad (2)$$

We work in a basis where h^d and h^e are diagonal, and h^u is a diagonal matrix times the Kobayashi-Maskawa matrix. Because of the mismatch in flavour space, squarks and sleptons mediate effective four-fermion interactions which lead to flavour-changing neutral-current processes and lepton-family transitions. For instance, measurements of the $K^0 - \bar{K}^0$ mixing parameters Δm_K and ϵ imply [7]

$$|\lambda_{121}| < \left(\frac{4 \times 10^{-2}}{|\lambda_{211}|} \right) \left(\frac{m_{\tilde{\nu}}}{200 \text{ GeV}} \right)^2 10^{-7}, \quad (3)$$

$$|\lambda_{121}| < \left(\frac{4 \times 10^{-2}}{|\lambda_{211}|} \right) \left(\frac{m_{\tilde{\nu}}}{200 \text{ GeV}} \right)^2 \frac{7 \times 10^{-10}}{\sin \delta_\lambda}, \quad (4)$$

where δ_λ is the relative phase between the two λ couplings. Bounds on μ - e conversion imply [7]

$$|\lambda_{i12}| < \left(\frac{4 \times 10^{-2}}{|\lambda_{i11}|} \right) \left(\frac{m_{\tilde{u}_i}}{200 \text{ GeV}} \right)^2 5 \times 10^{-6}, \quad (5)$$

for any $i = 1, 2, 3$. While limits on a single λ_{ijk} coupling are weak enough [8] to allow for an important phenomenological rôle of R -parity breaking, the product of two λ couplings with different generation indices is severely constrained. A successful theory of flavour and R -parity violation should explain the origin of this strong hierarchy.

Let us now turn to the *unification* problem. If the interaction of eq. (1) has to be embedded in a trilinear term arising from a GUT, then the superpotential in general also contains the interactions

$$W = \lambda'_{ijk} \bar{E}_R^i L_L^j L_L^k + \lambda''_{ijk} \bar{U}_R^i \bar{D}_R^j \bar{D}_R^k . \quad (6)$$

While the λ and λ' couplings violate lepton number, λ'' violates baryon number. Their simultaneous presence is therefore strongly constrained by nucleon-decay searches. For instance the experimental bound on $n \rightarrow K^+ e^-$ implies

$$|\lambda''_{i12}| \lesssim \left(\frac{4 \times 10^{-2}}{|\lambda_{i11}|} \right) \left(\frac{\tilde{m}}{200 \text{ GeV}} \right)^3 \left(\frac{175 \text{ GeV}}{m_{u_i}} \right) 10^{-25} . \quad (7)$$

Here \tilde{m} is the typical supersymmetry-breaking mass parameter in the \tilde{u}_i mass matrix. The presence of the quark mass m_{u_i} in eq. (7) is a product of the left-right squark mixing necessary to construct the $\Delta B = -\Delta L = -1$ four-fermion operator. Of course, in the case $i = 3$, it does not amount to any significant suppression.

From eq. (7) we see that the unification problem (or, in other words, the simultaneous presence of baryon and lepton number violation) poses a more severe difficulty than flavour. There can be hope that hierarchies between couplings analogous to those required by eqs. (3)–(5) can be explained in a complete theory of flavour. On the other hand, we prefer to believe that the observed suppression of nucleon decay is caused by the small ratio between the weak and the GUT scale rather than by some broken flavour symmetry, which generates hierarchies as a power expansion of some parameter like the Cabibbo angle. Examples of theories in which the baryon-number and R -violating interactions are suppressed by flavour symmetries exist [9], but are not embedded in a GUT. It is not clear how they can be unified and made consistent with the size of couplings suggested by the HERA data. For this reason, we believe that the solution should lie within the GUT dynamics.

3. We now want to embed the R -parity violating interaction of eq. (1) in a GUT, without running into the problem of baryon-number violation. To start, we choose the simplest example of $SU(5)$ and denote the matter content by $\mathbf{10}^i + \bar{\mathbf{5}}^i$ ($i = 1, 2, 3$) and the Higgs superfields by H and \bar{H} , respectively a $\mathbf{5}$ and $\bar{\mathbf{5}}$ of $SU(5)$. The Yukawa couplings are

$$W = h_{ij}(\Sigma) \mathbf{10}^i \mathbf{10}^j H + \bar{h}_{ij}(\Sigma) \mathbf{10}^i \bar{\mathbf{5}}^j \bar{H} . \quad (8)$$

Here h_{ij} and \bar{h}_{ij} are functions of the adjoint field Σ , which spontaneously breaks the $SU(5)$ symmetry. After Σ gets its vacuum expectation value (VEV), they reproduce the ordinary Yukawa couplings $h_{ij}^u, h_{ij}^d, h_{ij}^e$ at low energy.

The first attempt one can try is to include in the superpotential only the bilinear term²

$$W = \rho_i \bar{\mathbf{5}}^i H , \quad (9)$$

²The phenomenology of the non-GUT version of this term has been considered, for instance, in ref. [10].

where ρ_i are mass parameters smaller than the weak mass scale. The operators in eq. (9) could be generated by some mechanism similar to the one responsible for the Higgs-mixing μ term. By defining appropriate mass eigenstates, we can rotate the term in eq. (9) into some R -violating trilinear couplings. The \bar{D}_R^i states mix with the Higgs triplet contained in \bar{H} and give rise to the baryon-number violating coupling of eq. (6) with

$$\lambda''_{ijk} = \bar{h}_{ij} \frac{\rho_k}{M_H} . \quad (10)$$

Since the Higgs-triplet mass M_H is of the order of the GUT scale, we obtain a considerable suppression of the baryon-violating interaction. An even further suppression exists in models where the doublet-triplet splitting is obtained without a direct mass term $H\bar{H}$. The mixing between the lepton superfields L_L^i and the Higgs doublet generates couplings λ and λ' which are suppressed only by the ratio ρ_i/μ , where μ is the Higgs-mixing term of the order of the weak scale. However, in this case, $\lambda_{ijk} \propto h_{ij}^d \rho_k$ and the value of the R -violating coupling constant suggested by the HERA data is incompatible with the limit on the electron neutrino mass which implies [5]

$$|\lambda_{331}| < 5 \times 10^{-3} \left(\frac{m_{\bar{\nu}_e}}{200 \text{ GeV}} \right)^{\frac{1}{2}} . \quad (11)$$

An interesting possibility, which was first suggested in ref. [7], is that the only R -parity violation comes from an operator in the superpotential

$$\lambda_{ijk}^G(\Sigma) \mathbf{10}^i \bar{\mathbf{5}}^j \bar{\mathbf{5}}^k . \quad (12)$$

If $\lambda_{ijk}^G(\langle \Sigma \rangle)$ and $\bar{h}_{jk}(\langle \Sigma \rangle)$ (the Yukawa coupling defined in eq. (8)) are simultaneously diagonal in j, k for any i and for any $SU(5)$ index, then the interaction (12) generates nonvanishing λ , while λ' and λ'' identically vanish [7]. This is simply because λ'_{ijk} and λ''_{ijk} are antisymmetric in j, k , while λ_{ijk} has no symmetry properties. The coupling constants λ'_{ijk} vanish at the GUT scale, but small values are generated by the renormalization to the weak scale. It was also shown in ref. [7] that the renormalization to the weak scale gives only small further violation of flavour in the λ sector, and therefore the *ansatz* on the generation structure of λ_{ijk}^G specified above can render the R -breaking interpretation of the HERA data compatible with unification. However this *ansatz* may seem rather *ad hoc*. Also it seems to rely on flavour properties, which we find a disturbing aspect, as previously discussed.

We turn now to discuss how, with no reference to the flavour theory, a GUT can lead to R -parity violating couplings relevant for squark production at HERA together with vanishing baryon-number violating couplings. In terms of GUT representations, the R -violating interactions in eqs. (1) and (6) are written as³

$$\mathcal{O}_1^{ijk} \equiv \bar{\mathbf{5}}^j \cdot \bar{\mathbf{5}}^k \cdot \mathbf{10}^i \quad \rightarrow \quad \lambda_{ijk}, \lambda'_{ijk}, \lambda''_{ijk} \quad (13)$$

³This defines a complete set of operators. This can be understood by counting the number of gauge invariants. For N matter generations, the operators listed in eqs. (13)–(16) contain $N^2(N-1)/2$ or $N^2(N+1)/2$ flavour components, if the two $\bar{\mathbf{5}}$ are combined in a $\mathbf{10}$ or $\bar{\mathbf{15}}$, respectively. This makes a total of $N^2(2N-1)$ invariants and matches the number of invariants of the low-energy theory, which are described by N^3 couplings λ and $N^2(N-1)/2$ couplings λ' and λ'' .

$$\mathcal{O}_2^{ijk} \equiv (\bar{\mathbf{5}}^j \cdot \bar{\mathbf{5}}^k)_{\mathbf{10}} \cdot (\mathbf{10}^i \cdot \Sigma)_{\mathbf{10}} \quad \rightarrow \quad \lambda_{ijk}, \lambda'_{ijk}, \lambda''_{ijk} \quad (14)$$

$$\mathcal{O}_3^{ijk} \equiv (\bar{\mathbf{5}}^j \cdot \bar{\mathbf{5}}^k)_{\mathbf{15}} \cdot (\mathbf{10}^i \cdot \Sigma)_{\mathbf{15}} \quad \rightarrow \quad \lambda_{ijk} \quad (15)$$

$$\mathcal{O}_4^{ijk} \equiv \left((\bar{\mathbf{5}}^j \cdot \bar{\mathbf{5}}^k)_{\mathbf{10}} \cdot \Sigma \right)_{\mathbf{15}} \cdot (\mathbf{10}^i \cdot \Sigma)_{\mathbf{15}} \quad \rightarrow \quad \lambda_{ijk} . \quad (16)$$

Here we have specified the contractions of the $SU(5)$ indices with a pindex denoting the product representation. Operators with more powers of $\langle \Sigma \rangle$ reduce to combinations of the above since, for any n , $\langle \Sigma \rangle^n$ is a linear combination of $\langle \Sigma \rangle$ and the identity. We have also marked explicitly which of the couplings λ , λ' , or λ'' are generated by the various operators after Σ gets its VEV. If we select the operators \mathcal{O}_3 or \mathcal{O}_4 , at low energies we retain only the λ couplings. This can be easily implemented in a GUT, since specific operators can be selected by appropriately choosing the virtual states which generate them. This selection is a consequence of group-theoretical properties and does not rely on the flavour dynamics.

\mathcal{O}_3^{ijk} combines the two $\bar{\mathbf{5}}$ in a symmetric state. Therefore it is symmetric in the indices j and k and its contribution to λ' and λ'' identically vanishes. We will give later an example of how this case can be realized in GUTs. \mathcal{O}_4^{ijk} also selects just the λ coupling, although it is antisymmetric in the flavour indices j and k . Only the coupling λ survives because $(\mathbf{10}^i \cdot \Sigma)_{\mathbf{15}}$ is projected only onto Q_L^i , after GUT symmetry breaking.

Solving the problem of baryon-number and R -parity violating interactions with GUT dynamics may not be sufficient. We are now concerned with higher-dimensional operators suppressed by powers of the Planck mass M_P generated by the unknown dynamics of quantum gravity. These operators may have the most generic structure compatible with the unbroken symmetries, and therefore reintroduce the unwanted λ'' couplings in the low-energy effective theory. These couplings will only be suppressed by some powers of M_{GUT}/M_P , and therefore bounds like the one in eq. (7) require that these operators should not be present at least up to some high dimensionality.

The non-renormalization theorems [11] of supersymmetry can protect the GUT theory from this danger. If the operator \mathcal{O}_3 or \mathcal{O}_4 is generated only after some stage of symmetry breaking, this same symmetry together with the constraint of holomorphicity of the superpotential can forbid any dangerous higher-dimensional operator. This mechanism, in which a broken symmetry protects against the appearance of certain terms in the superpotential at all orders, has also been used in the constructions of flavour theories [12]. Of course the renormalization of the Kähler function is not under control and it can effectively generate new terms in the superpotential, once supersymmetry is broken. The size of these effects, which are proportional to some power of the ratio between the supersymmetry-breaking scale and M_P can be estimated in a given model and then compared with the experimental bound on nucleon decay.

4. Our mechanism is best illustrated by a simple example in $SU(5)$. To generate the desired operators \mathcal{O}_3 and \mathcal{O}_4 , we introduce some fields in the symmetric product of two fundamentals, $S + \bar{S}$, which transform as $\mathbf{15} + \mathbf{15}$. The presence of the field S together with its conjugate \bar{S} insures the cancellation of $SU(5)$ anomalies and allows a superheavy mass term. Let us consider the following interaction for the fields S and \bar{S} :

$$W = \bar{\mathbf{5}}^i \bar{\mathbf{5}}^j S + \bar{S} \mathbf{10}^i \Sigma + S \bar{S} \phi . \quad (17)$$

Here ϕ is a gauge singlet, which plays the rôle of a mass parameter, and $\langle\phi\rangle \equiv M_X$ is somewhat larger than the GUT scale. The effective theory below M_X , obtained by integrating out S and \bar{S} , contains the operator \mathcal{O}_3 . Just below M_{GUT} , λ_{ijk} is generated, but λ'_{ijk} and λ''_{ijk} vanish.

In order to explain the desired structure of R -breaking interaction, we also have to justify the absence of the renormalizable coupling $\mathbf{10}^i \bar{\mathbf{5}}^j \mathbf{5}^k$. As explained above, this is in general not sufficient, because the strong bounds on nucleon decay also require the absence of a large number of higher-dimensional operators. The simplest symmetry we can introduce is an abelian *flavour-independent* $U(1)$. The $U(1)$ charge assignment is completely determined by the requirement that the most general superpotential consistent with the $SU(5) \times U(1)$ symmetry is given by eqs. (8) and (17) together with terms responsible for the GUT symmetry breaking involving Σ , bilinears in $\bar{H}H$, and possibly other fields. We find the following $U(1)$ charges: $X(\mathbf{10}^i) = -1$, $X(\bar{\mathbf{5}}^i) = 3$, $X(H) = 2$, $X(\bar{H}) = -2$, $X(\Sigma) = 0$, $X(S) = -6$, $X(\bar{S}) = 1$, $X(\phi) = 5$. This $U(1)$ is anomalous, but the Green-Schwarz mechanism [13] can be invoked to cancel the gauge anomalies. It is interesting that an anomalous $U(1)$ group usually appears in the effective field theory derived from strings [14]. The effective theory then contains a Fayet-Iliopoulos term, equal to [14]

$$\xi = \frac{g^2 \text{Tr } X}{192\pi^2} M_P^2. \quad (18)$$

If the signs of $X(\phi)$ and $\text{Tr}X$ are opposite (and, in our example, this is true at least in the observable sector), then ϕ can get a VEV, given by

$$\langle\phi\rangle \equiv M_X = \sqrt{\frac{-\xi}{X(\phi)}}. \quad (19)$$

The theory has an accidental discrete symmetry, under which $\mathbf{10}^i$, $\bar{\mathbf{5}}^i$, \bar{S} , and ϕ are odd, while all other chiral superfields are even. This can be identified with the usual R parity, and it is broken by $\langle\phi\rangle$ at the scale M_X . The size of the low-energy R -parity violating coupling λ is $\lambda \sim \mathcal{O}(M_{GUT}/M_X)$.

We turn now to discuss the suppression of higher-dimensional operators. The property of holomorphicity of the superpotential and the $U(1)$ symmetry forbid all possible quantum-gravity derived operators, which give rise to R -parity breaking in the low-energy effective theory. Indeed terms of the generic form

$$\int d^2\theta H \bar{\mathbf{5}}^i f(\Sigma, \phi) \quad \text{or} \quad \int d^2\theta \mathbf{10}^i \bar{\mathbf{5}}^j \bar{\mathbf{5}}^k g(\Sigma, \phi) \quad (20)$$

cannot appear since holomorphicity requires only positive powers of ϕ in the functions f and g , while $U(1)$ invariance requires a negative power of ϕ . Of course this property depends on the particular charge assignment and it would not hold if there existed fields which acquire VEVs of the order of the GUT scale and have negative $U(1)$ charges.

Planck-mass suppressed operators can also affect the dynamics of the fields S and \bar{S} . In particular the most general interactions consistent with $SU(5) \times U(1)$ symmetry have the same form of those in eq. (17), with arbitrary insertions of Σ fields. The key point is that these

operators are not going to modify our mechanism, since they can generate \mathcal{O}_4 , but never \mathcal{O}_1 or \mathcal{O}_3 . We can understand this result differently by considering the $SU(3) \times SU(2) \times U(1)$ content of the **15**. The **15** does not contain any standard model representation with the correct quantum numbers to mediate effective interactions $\bar{U}_R \bar{D}_R \bar{D}_R$ or $\bar{E}_R L_L L_L$. Only a colour triplet, weak doublet can be propagated between the $\bar{\mathbf{5}}^i \bar{\mathbf{5}}^j$ and $\mathbf{10}^k \Sigma$ states. From this point of view, our mechanism is analogous to the missing partner mechanism [15] used to split the masses of the Higgs doublet and triplet belonging to the same GUT representation⁴.

The Kähler function can contain terms of the kind

$$\int d^4\theta H \bar{\mathbf{5}}^i \phi^\dagger F(Z, Z^\dagger) \quad \text{and} \quad \int d^4\theta \mathbf{10}^i \bar{\mathbf{5}}^j \bar{\mathbf{5}}^k \phi^\dagger G(Z, Z^\dagger), \quad (21)$$

where Z is the spurion superfield which parametrizes supersymmetry breaking and has a non-zero VEV of the auxiliary field. The terms in eq. (21) give rise to the low-energy parameters $\rho \sim (m_{\tilde{G}} M_X)/M_P$ and $\lambda \sim \lambda' \sim \lambda'' \sim (m_{\tilde{G}} M_X)/M_P^2$, where $m_{\tilde{G}}$ is the gravitino mass. In theories where the breaking of supersymmetry is communicated to the observable sector by gravity [17] $m_{\tilde{G}}$ is of the order of the weak scale. Then, a comparison with the bound in eq. (7) shows that a certain degree of suppression coming either from the flavour theory or from quantum gravity is necessary. Lacking much knowledge of either of the two theories, we cannot exclude this case. On the other hand, in theories where supersymmetry breaking is communicated by particles lighter than M_P , $m_{\tilde{G}}$ is smaller than the weak scale and it can efficiently suppress any Kähler-induced R -parity violation.

The last comment we wish to make about this model concerns the flavour structure of the R -parity violation. Although our mechanism does not address the question of flavour, as a byproduct of this model, we obtain a restriction on the generation structure of the λ_{ijk} couplings which, just below the GUT scale, have the form

$$\lambda_{ijk} = A_i B_{jk}. \quad (22)$$

This is a consequence of our assumption that a single state $S + \bar{S}$ mediates the effective interaction which generates the operator \mathcal{O}_3 . Only within a complete theory of flavour can we hope to understand the hierarchical structure of the vector A and the matrix B .

5. The model discussed in the previous section is certainly not the only possibility to realize our mechanism. Instead of an anomalous $U(1)$ symmetry, one could use an R symmetry which easily protects against higher-dimensional terms. However, in this case, one should specify the whole model, including the GUT symmetry breaking sector, which it was left undetermined in our previous example.

Another possibility is to consider different GUT groups, flipped $SU(5)$ [18] being a very interesting option. The GUT group is $SU(5) \times U(1)$, with matter transforming as $(\mathbf{10}, -1) + (\bar{\mathbf{5}}, 3) + (\mathbf{1}, -5)$, and the usual Higgs doublets embedded in $H = (\mathbf{5}, 2)$ and $\bar{H} = (\bar{\mathbf{5}}, -2)$. The Yukawa couplings are

$$W = h_{ij}^u \mathbf{10}^i \bar{\mathbf{5}}^j \bar{H} + h_{ij}^d \mathbf{10}^i \mathbf{10}^j H + h_{ij}^e \bar{\mathbf{5}}^i \mathbf{1}^j H, \quad (23)$$

⁴The use of an anomalous $U(1)$ to implement to all orders the missing partner mechanism was studied in ref. [16].

where we have denoted the matter superfields by their $SU(5)$ content. The GUT symmetry breaking is triggered by the VEVs of the fields $K = (\mathbf{10}, -1)$ and $\bar{K} = (\bar{\mathbf{10}}, 1)$, which allow a simple implementation of the missing-partner mechanism [18] through the interactions

$$W = HKK + \bar{H}\bar{K}\bar{K} . \quad (24)$$

An interesting feature of flipped $SU(5)$ is that renormalizable R -parity interactions are forbidden by gauge invariance. The low-energy R -violating operators in eqs. (1) and (6) are generated by the following GUT operators:

$$O_1^{ijk} \equiv (\mathbf{10}^j \cdot \mathbf{10}^k)_{\bar{\mathbf{5}}} \cdot (\bar{\mathbf{5}}^i \cdot K)_{\bar{\mathbf{5}}} \quad \rightarrow \quad \lambda_{ijk} \quad (25)$$

$$O_2^{ijk} \equiv (\bar{\mathbf{5}}^i \cdot \mathbf{10}^j)_{\bar{\mathbf{5}}} \cdot (\mathbf{10}^k \cdot K)_{\bar{\mathbf{5}}} \quad \rightarrow \quad \lambda_{ijk}, \lambda''_{ijk} \quad (26)$$

$$O_3^{ijk} \equiv \left((\mathbf{10}^j \cdot \mathbf{10}^k)_{\bar{\mathbf{45}}} \cdot (K \cdot \bar{K})_{\mathbf{24}} \right)_{\bar{\mathbf{5}}} \cdot (\bar{\mathbf{5}}^i \cdot K)_{\bar{\mathbf{5}}} \quad \rightarrow \quad \lambda_{ijk} \quad (27)$$

$$O_4^{ijk} \equiv \bar{\mathbf{5}}^j \cdot \bar{\mathbf{5}}^k \cdot \mathbf{1}^i \cdot K \quad \rightarrow \quad \lambda'_{ijk} . \quad (28)$$

This is a complete set of operators. Of course one can rearrange the contractions of the $SU(5)$ indices or take linear combination of the various operators. By doing this, we can construct an operator which selects only the λ'' couplings:

$$O_5^{ijk} \equiv \left((\mathbf{10}^j \cdot K)_{\bar{\mathbf{5}}} \cdot (\mathbf{10}^k \cdot K)_{\bar{\mathbf{5}}} \right)_{\mathbf{10}} \cdot (\bar{\mathbf{5}}^i \cdot \bar{K})_{\mathbf{10}} \quad \rightarrow \quad \lambda''_{ijk} \quad (29)$$

This operator is generated by exchange of states with flipped $SU(5)$ quantum numbers $(\mathbf{10}, 4)$, $(\bar{\mathbf{10}}, -4)$.

If we select only O_1 or O_3 , we obtain the R -violating coupling invoked for an interpretation of the HERA data, but forbid the baryon-number violating ones. The operator O_1 is generated by the virtual exchange of the superfields $S + \bar{S}$, transforming as $(\mathbf{5}, 2) + (\bar{\mathbf{5}}, -2)$, with interactions

$$W = \mathbf{10}^i \mathbf{10}^j S + \bar{S} \bar{\mathbf{5}}^i K + S \bar{S} \phi , \quad (30)$$

where ϕ is a gauge singlet such that $\langle \phi \rangle = M_X$. However we have to forbid the couplings $\mathbf{10}^i K S + \bar{S} \bar{\mathbf{5}}^i \mathbf{10}^j$, which generate the operator O_2 and introduce the baryon-number violating couplings λ'' . All unwanted structures are eliminated in renormalizable and non-renormalizable interactions by a discrete R -parity and an anomalous $U(1)$. The superfields $\mathbf{10}^i$, $\bar{\mathbf{5}}^i$, $\mathbf{1}^i$, \bar{S} , and ϕ are odd under R -parity, while all others are even. The charge assignment of the anomalous $U(1)$, up to a linear combination with the $U(1)$ of flipped $SU(5)$, is $X(\bar{\mathbf{5}}^i) = -X(\mathbf{1}^i) = -X(\bar{H}) = 2X(\bar{K}) = -X(\bar{S}) = X(\phi) = 1$, while all other fields are neutral. Both the R -parity and the anomalous $U(1)$ are spontaneously broken the VEVs of ϕ and \bar{K} . There are no fields with GUT scale VEV and negative charge X , and therefore holomorphicity and symmetry invariance protect against operators with unwanted structures of R -parity breaking. As in the previous example, higher-dimensional operators involving the fields S and \bar{S} do not affect the mechanism, since only a single component of the S field (a colour triplet, weak doublet) is propagated in the exchange between the $\mathbf{10}^i \mathbf{10}^j$ and $\bar{\mathbf{5}}^k K$ states. We also notice that, in the flipped $SU(5)$ case, the flavour structure of the R -violating coupling λ_{ijk} factorizes between quark and lepton indices

$$\lambda_{ijk} = A_{ij} B_k , \quad (31)$$

in contrast to the case of eq. (22).

Finally notice that the operator O_4 can be generated by S -exchange through the couplings $\bar{5}^i K \bar{5} + S \bar{5}^i \mathbf{1}^j$. The simultaneous presence of λ and λ' interactions is constrained by flavour and lepton violating processes. However, since flipped $SU(5)$ is not based on a simple group and the unification is not complete, the couplings λ and λ' are here unrelated. It is therefore possible that some flavour symmetry suppresses λ' without affecting the coupling λ invoked to explain the HERA anomaly.

In the case of $SO(10)$, left-right symmetry forbids R -violating dimension-four operators. We can then consider dimension-five operators of the kind $\mathbf{16}^i \mathbf{16}^j \mathbf{16}^k \mathbf{16}_H$, where $\mathbf{16}^i$ describes a generation of matter and $\mathbf{16}_H$ is the Higgs representation which breaks $SO(10)$ into $SU(5)$. Since these operators are $SU(5)$ invariant, they predict $\lambda = \lambda' = \lambda''$. It is possible to construct higher-dimensional operators which select only the λ coupling. For instance

$$(\mathbf{16}^j \cdot \mathbf{16}_H)_{\mathbf{10}} \cdot (\mathbf{16}^k \cdot \mathbf{16}_H)_{\mathbf{10}} \cdot \mathbf{54} \quad \text{and} \quad \mathbf{16}^i \cdot (\mathbf{45} \cdot \mathbf{54})_{\mathbf{45}} \bar{\mathbf{16}}_H, \quad (32)$$

which are mediated by a heavy $\mathbf{45}$ and $\mathbf{10}$ respectively, can give the right structure of R -violation, once the heavy $\mathbf{54}$ is exchanged. The model is obviously rather involved and depends on the unspecified dynamics which characterize the symmetry-breaking pattern.

6. We have shown that the R -parity violating interaction suggested by the HERA data and the observed absence of fast nucleon decay is compatible with the idea of grand unification. The interactions between matter and the heavy fields which break the GUT symmetry can split the different R -violating interactions. By choosing an appropriate field content, it is possible to select a preferred pattern of R -parity breaking. This pattern is protected against higher-dimensional operators with different R -violating structures by the combined effect of a symmetry, broken at the GUT scale, and the holomorphicity of the superpotential.

Our mechanism relies on purely group-theoretical properties and does not require any assumption about flavour structure. We have shown specific examples of GUT models in which our idea can be realized. Only in the context of a complete theory of flavour will it be possible to address the question of flavour-changing neutral current processes and lepton-family transitions.

We wish to thank G. Altarelli, G. Dvali, and M. Mangano for useful discussions.

References

- [1] C. Adloff *et al.*, H1 collaboration, DESY preprint 97-24, hep-ex/9702012.
- [2] J. Breitweg *et al.*, ZEUS collaboration, DESY preprint 97-25, hep-ex/9702015.
- [3] D. Buskulic *et al.*, ALEPH collaboration, *Z. Phys.* **C71** (1996) 179,
G. Cowan, for the ALEPH collaboration, CERN seminar, Feb. 25th, 1997.

- [4] W.D. Schlatter, for the LEP Working Group on Four-Jet Final States, CERN seminar, Feb. 25th, 1997.
- [5] D. Choudhury and S. Raychaudhuri, CERN preprint TH/97-26, hep-ph/9702392;
G. Altarelli, J. Ellis, G.F. Giudice, S. Lola, and M.L. Mangano, CERN preprint TH/97-40, hep-ph/9703276;
H. Dreiner and P. Morawitz, hep-ph/9703279;
J. Kalinowski, R. Ruckl, H. Spiesberger, and P.M. Zerwas, preprint BI-TP-97-07, hep-ph/9703288;
K.S. Babu, C. Kolda, J. March-Russell, and F. Wilczek, preprint IASSNS-HEP-97-04, hep-ph/9703299;
T. Kon and T. Kobayashi, preprint ITP-SU-97-02, hep-ph/9704221.
- [6] M. Carena, G.F. Giudice, S. Lola, and C.E.M. Wagner, *Phys. Lett.* **B395** (1997) 225.
- [7] R. Barbieri, A. Strumia, and Z. Berezhiani, Pisa preprint IFUP-TH 13/97, hep-ph/9704275.
- [8] V. Barger, G.F. Giudice, and T. Han, *Phys. Rev.* **D40** (1989) 2987.
- [9] V. Ben-Hamo and Y. Nir, *Phys. Lett.* **B339** (1994) 77;
C.D. Carone, L.J. Hall, and H. Murayama, *Phys. Rev.* **D53** (1996) 6282;
C.D. Carone, L.J. Hall, and H. Murayama, *Phys. Rev.* **D54** (1996) 2328.
- [10] L.J. Hall and M. Suzuki, *Nucl. Phys.* **B231** (1984) 419;
T. Banks, Y. Grossman, E. Nardi, and Y. Nir, *Phys. Rev.* **D52** (1995) 5319;
H.-P. Nilles and N. Polonsky, *Nucl. Phys.* **B484** (1997) 33.
- [11] M.T. Grisaru, W. Siegel, and M. Rocek, *Nucl. Phys.* **B159** (1979) 429.
- [12] T. Banks, Y. Nir, and N. Seiberg, Proc. of 2nd IFT Workshop on Yukawa Couplings and the Origins of Mass, Gainesville 1994, hep-ph/9403203.
- [13] M. Green and J. Schwarz, *Phys. Lett.* **B149** (1984) 117.
- [14] M. Dine, N. Seiberg, and E. Witten, *Nucl. Phys.* **B289** (1987) 585;
J. Atick, L. Dixon, and A. Sen, *Nucl. Phys.* **B292** (1987) 109;
M. Dine, I. Ichinose, and N. Seiberg, *Nucl. Phys.* **B293** (1987) 253.
- [15] S. Dimopoulos and F. Wilczek, ITP Santa Barbara preprint UM HE 81-71 (1981);
A. Masiero, D.V. Nanopoulos, K. Tamvakis, and T. Yanagida, *Phys. Lett.* **B115** (1982) 380;
B. Grinstein, *Nucl. Phys.* **B206** (1982) 387.
- [16] Z. Berezhiani and Z. Tavartkiladze, *Phys. Lett.* **B396** (1997) 150.
- [17] R. Barbieri, S. Ferrara, and C.A. Savoy, *Phys. Lett.* **B119** (1982) 343;
P. Nath, R. Arnowitt, and A. Chamseddine, *Phys. Rev. Lett.* **49** (1982) 970.

- [18] S. Barr, *Phys. Lett.* **B112** (1982) 219;
I. Antoniadis, J. Ellis, J. Hagelin, and D.V. Nanopoulos, *Phys. Lett.* **B194** (1987) 231;
for a review, see J. Lopez and D.V. Nanopoulos, Proc. of Int. School on Subnuclear Physics,
Erice, Italy, July 1995, hep-ph/9511266.