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CERN-TH.6111/91

Discrete Gauge Symmetries and the Origin of Baryon and Lepton Number Conservation in Supersymmetric Versions of the Standard Model

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ABSTRACT

In the supersymmetric standard model operators of dimension 4 and 5 generically violate B and L number. One usually assumes the presence of some discrete symmetry ("matter parities") in order to forbid dangerous operators which may lead otherwise to unacceptable violations of B and L . We give a general classification of such discrete Z_N symmetries (and R-symmetries) and show that the number of independent possibilities is substantially reduced by equivalences. We argue that normal discrete symmetries may be expected to be violated by quantum gravity effects and hence are not enough to inhibit nucleon decay. On the other hand, gauge (either discrete or continuous) symmetries are stable under quantum gravity effects and we discuss how such symmetries may eliminate the dangerous B or L violating operators. We find that the massless fermion content of models with discrete "gauge" symmetries is strongly constrained by the cancellation of "discrete gauge anomalies". We show that there are only two Z_N symmetries which are discrete anomaly free with the minimal light matter content. One of them is the standard R-parity whereas the other is a unique Z_3 symmetry allowing for lepton number violation. We argue that from the point of view of arranging for proton stability without fine-tuning the second option should be preferred. The differences in the phenomenology of the various supersymmetric models dictated by the different symmetries are discussed.

1. Introduction

The construction of a supersymmetric version of the standard model [1] immediately runs into difficulty for, unlike the standard model, there are terms of dimension 4 allowed by the $SU(3) \times SU(2) \times U(1)$ gauge symmetry which violate baryon- and lepton-number [2], [3]. If present these terms will normally lead to unacceptably fast rates of nucleon decay for they appear unsuppressed by any small inverse mass scale. For this reason these terms must be avoided and it is assumed this can only be done in a satisfactory manner if there are additional symmetries, either further gauge symmetries or discrete symmetries. For example the minimal supersymmetric standard model (MSSM) assumes a "matter parity" to eliminate both the baryon- and lepton- number violating terms [2], [3].

An immediate consequence of this is that the MSSM conserves what is known as R-parity under which all the superpartners of the standard model states are R-parity odd while all the standard model states are R-parity even. As a result the new supersymmetric states can only be produced in pairs and a supersymmetric state cannot decay only to conventional states. This has a profound effect [4], [5], [6] on the phenomenology of such states; in particular all experimental searches for the new supersymmetric states rely on pair production and most searches involve missing energy as a signal for the production of the "LSP" the lightest supersymmetric state which must be stable and neutral (for cosmological reasons). However matter parity (or equivalently R-parity) need not be conserved in supersymmetric extensions of the standard model. In view of the significant changes in the phenomenology of supersymmetry resulting from alternative schemes we feel it important to consider in detail the alternatives. Furthermore, one would like to understand the physical origin (presumably a symmetry) of such an important property as the stability of the proton. In this paper we will systematically explore the possible additional discrete and/or gauge symmetries in supersymmetric models which lead to acceptable rates for nucleon decay.

It is normally considered sufficient to eliminate only the leading dimension-4 operator contributions to nucleon decay for the higher dimension terms are suppressed by inverse powers of the "Grand Unified" mass at which the baryon- and lepton- number violating processes occur. However in 4 dimensional string theories all terms not suppressed by local gauge or discrete symmetries typically occur at order 1. Hence the expectation is that the higher dimension operators will at most be suppressed by inverse powers of the compactification scale which is itself bounded by the Planck scale. Even though this is very large, of order 10^{19} GeV , dimension-5 operators which are suppressed by only a single inverse power of mass may still generate nucleon decay at an unacceptable rate. Thus in a

generic string theory it is necessary to eliminate these terms too, requiring the introduction of symmetries other than the usual matter parity.

In Section 2 we discuss which baryon- and lepton- number operators of dimension-4 and dimension-5 may occur consistent with the gauge symmetries of the standard model. We summarise the constraints on the coefficients of dimension-4 operators and discuss the equivalent constraints on the coefficients of dimension-5 operators. We also briefly consider the constraints following from the requirement of a baryon number asymmetry in the early universe. In Section 3 we consider the discrete symmetries that may be invoked to select a suitable subset of the allowed operators, paying particular regard to symmetries commuting with flavour. A general analysis is presented for the case there is a single Z_N symmetry and the set of possible symmetries are determined for the simplest Z_2 and Z_3 symmetries. This leads to the identification of a small number of viable supersymmetric extensions of the standard model. In the case that only the dimension-4 contributions to nucleon decay are forbidden (and the minimal particle content is assumed) there are just three Z_2 (including the usual R parity) and four Z_3 symmetries allowed. In the case that the dimension-5 contributions are also required to be excluded there are only two Z_2 and three Z_3 symmetries allowed, a set which does not include the usual R-parity. In section 4 the analysis is extended to the case of Z_N symmetries not commuting with supersymmetry, i.e., Z_N R-symmetries, and it is again shown that only a reduced number of Z_2 and Z_3 symmetries are consistent with the phenomenological constraints.

In sections 5 and 6 we turn to a discussion of the possible origin of these symmetries. An important consideration is that the symmetries should be protected against large corrections coming from other parts of the theory and it has been suggested that only symmetries related to underlying gauge symmetries will be robust enough to eliminate large gravitational corrections (wormholes etc.) generating fast proton decay [7]. We consider the possibility that additional gauge symmetries lead to the suppression of the unwanted operators. Those gauge symmetries may be discrete or continuous. The case of *discrete gauge symmetries* [8], [9] has been only recently addressed [10]. Discrete gauge symmetries have the additional appeal of being naturally present in string theories. We observe that the cancellation of "*discrete gauge anomalies*" [10] strongly constrains the possible discrete gauge symmetries. Remarkably we find that there are only two anomaly-free flavour-blind gauge Z_N symmetries, one of them being the usual R-parity. The other symmetry is a Z_3 *Baryon parity* allowing for lepton number violation. The latter symmetry is in fact more satisfactory concerning proton stability since it forbids the presence of dangerous dimension-5 operators leading to proton decay. In section 6 we

briefly discuss the case of standard continuous gauge symmetries constraining the lepton and baryon number violation. The minimal model with a gauged $U(1)$ coupling to lepton number is presented. It is only anomaly-free for three quark-lepton generations and may lead to a model with a gauge Z_3 symmetry upon symmetry breaking which, in agreement with the discrete anomaly constraints, has additional right-handed neutrinos-like states. We also derive a $U(1)_B$ model which, upon symmetry breaking gives rise to the favoured Z_3 Baryon parity model demonstrating how this model can have only the minimal light multiplet content. In section 7 some issues concerning generation dependent generalized parities are discussed and a specific example presented. Finally in Section 8 we present our summary and conclusions and discuss some phenomenological aspects of the “baryon parity” model.

2. Baryon- and Lepton- number violating operators.

2.1 Dimension 3 and 4 operators

We start with a discussion of the renormalisable terms of dimension ≤ 4 allowed by the standard model gauge symmetry with the minimal particle content of the supersymmetric standard model. The MSSM is built from the minimal spectrum adding the Yukawa and associated scalar particle couplings needed to give masses to the quarks and leptons [1]. These come from the superpotential terms

$$[h_{iL} \bar{L}_i H e_k + h'_{iL} \bar{Q}_i H d_k + h''_{iL} \bar{Q}_i \tilde{H} u_k]_F \quad (1)$$

where L and e (Q and u , d) are the left-handed components of lepton doublet and antilepton singlet (quark doublet and antiquark singlet) chiral superfields respectively and $H, \tilde{H}$ are Higgs superfields. In this equation $SU(3)$ and $SU(2)$ indices have been suppressed. The terms of eq.(1) gives rise to the Yukawa couplings involving the component states. For example the first term gives

$$h_{iL} (\tilde{L}_i \tilde{H} e_k + L_i H e_k + L_i \tilde{H} \tilde{e}_k) + h.c. \quad (2)$$

where we denote by a supertwiddle the scalar partners to the quarks, leptons and Higgs namely the squarks, sleptons and Higgsinos. Note that the supersymmetric states only occur in pairs in eq.(2) consistent with the R-parity of the model. The terms in eqs.(1)

and (2) are of course allowed by the $SU(3) \times SU(2) \times U(1)$ gauge invariance but as noted above they are not the only renormalisable terms allowed. The most general gauge invariant superpotential giving a renormalisable Lagrangian consists of those in eq.(1) together with the terms

$$[\mu H \tilde{H} + m_i L_i H + \lambda_{ijk} \bar{L}_i L_j e_k + \lambda'_{ijk} \bar{L}_i Q_j d_k + \lambda''_{ijk} \bar{u}_i d_j d_k]_F \quad (3)$$

We will denote the last three operators in the above equation by N_1 , N_2 and N_3 respectively. The second term gives rise to mixing between the lepton and Higgs superfields and may be eliminated by a mass diagonalisation; doing this the terms of eq.(1) give rise to the third and fourth operators of eq.(3). We assume henceforth that this has been done. The last three terms of eq.(3) violate baryon- and lepton- number and, if all are present in the Lagrangian, they generate an unacceptably large amplitude for proton decay suppressed only by the inverse supersymmetry breaking mass scale squared [2], [3]. For this reason the MSSM requires an additional discrete symmetry called matter parity to forbid them. Under this symmetry the quark and lepton superfields appearing in the superpotential change sign while the Higgs superfields are left invariant. Thus the last four terms of eq.(3) change sign under this symmetry and are forbidden while the terms of eq.(1) are invariant and allowed. Using this symmetry the allowed couplings of the MSSM are those of eq.(1) and (cf eq.(2)) the new supersymmetric states only couple in pairs in accord with the R-parity of the model. In addition the first term of eq.(3) is allowed and indeed is needed for an acceptable pattern of electroweak symmetry breaking.

However it is not necessary to forbid the last four terms of eq.(3) to stabilise the nucleon [4]. The graph of Fig.(1) is responsible for nucleon decay and it may be seen from it that provided at least one of the last two baryon- or lepton- number violating vertices of eq.(3) is absent the nucleon will be stable. In this case the remaining operators will violate matter- and R- parity. For example the third term gives rise to the Yukawa couplings

$$\lambda_{ijk} (\bar{L}_i L_j \tilde{e}_k + L_i \tilde{L}_j e_k + \tilde{L}_i L_j e_k) + h.c. \quad (4)$$

involving a coupling to *single* sparticle states. The effect of these terms, below the threshold for superparticle production, is to induce corrections to the standard model via virtual superparticle exchange and this constrains some of the possible couplings [4]. Allowing for the most general family structure in eq.(3) there are 45 possible new dimension 4 terms and the associated couplings (assuming the absence of combinations leading to nucleon

decay) are constrained [4]. The bounds, particularly for the heavier generations, are not very severe except for the uds operator whose coupling is strongly constrained by heavy nuclei decay to be less than 10^{-9} . Except for this case the allowed magnitude of couplings is of the same range as the allowed magnitude of Yukawa couplings.

From a purely phenomenological point of view, therefore, the first priority in building a viable supersymmetric version of the standard model is to ensure the absence of combinations of dimension 4 terms leading to nucleon decay for they are constrained to be much smaller than each individual operator constraints just discussed of any of the “allowed” operators of eq.(3). A calculation of the graph of Fig. 1 together with an estimate of the associated operator matrix elements gives the constraint [3]

$$\lambda'\lambda'' < M_{SUSY}^2/(10^{16} GeV)^2 \sim O(10^{-26}) \quad (5)$$

where the bound refers to operators involving light quark external states and M_{SUSY} is the mass of the intermediate squark which is expected to be $O(1 TeV)$. It is of interest therefore to explore the symmetries capable of achieving this suppression before discussing the (family and flavour) symmetries which may give rise to the observed pattern of fermion masses. This will be discussed in Section 3 but first we discuss whether the suppression of the dimension 4 operator contribution to nucleon decay is sufficient to stabilise the proton.

2.2 Dimension 5 baryon- and lepton- number violating terms

Although dimension 4 terms are the most dangerous for nucleon decay they are not the only ones that must be suppressed. Dimension 5 operators may lead to nucleon decay suppressed by $1/M$ where M is the GUT mass scale associated with new baryon- and lepton- number violating physics beyond the (supersymmetric) standard model. In a compactified string theory this scale may be the compactification scale (of order the Planck scale) or a lower scale associated with the breaking of a GUT left unbroken after compactification. Even if we assume the scale is as high as the Planck scale some of the dimension 5 operators will generate proton decay at an unacceptable rate if the coefficients of the operators occur at order 1, so it is important to tabulate the possible operators and to extend to them the analysis of symmetries capable of eliminating their contribution to nucleon decay.

The dimension 5 baryon and lepton number violating operators allowed by the $SU(3) \times$

$SU(2) \times U(1)$ gauge symmetry are (suppressing family as well as gauge indices)

$$\begin{aligned} O_1 &= [QQQL]_F \\ O_2 &= [uude]_F \\ O_3 &= [QQQH]_F \\ O_4 &= [QucH]_F \\ O_5 &= [LL\bar{H}\bar{H}]_F \\ O_6 &= [LH\bar{H}\bar{H}]_F \\ O_7 &= [\bar{H}\bar{H}\epsilon^*]_D \\ O_8 &= [\bar{H}^*H\epsilon]_D \\ O_9 &= [QuL^*]_D \\ O_{10} &= [ud^*e]_D \end{aligned} \quad (6)$$

The contribution of these operators to nucleon decay occur only in radiative order for they all involve supersymmetric partners of standard model states. Thus, for example, nucleon decay due to the first operator occurs through the graph of Fig. 2 in which gaugino exchange converts the supersymmetric states back to normal states. Up to couplings the amplitude for this process is suppressed relative to that of eq.(5) coming from Fig. 1 by the loop factor and the factor (M_{SUSY}/M) . Using the Planck scale for M , and taking $M_W \sim M_W$ the bound on the coefficient η_1 of O_1 is $\eta_1 < 10^{-7}$ [11] for the operator involving the light quarks.

Not all the dimension five operators are equally dangerous. For example the operators O_1 involving heavy quarks will have contributions to nucleon decay suppressed by powers of mixing angles. This will be of importance when considering flavour symmetries capable of bringing nucleon decay within observed bounds. The other operators $O_2, \dots, O_{10}$ have varying contributions to nucleon decay. O_2 involves superfields which do not couple to the Winos and so the graph of Fig. 2 vanishes. The leading contribution will involve small Yukawa couplings and the resulting bound on its coefficients are rather mild and comparable to those on the Yukawa couplings responsible for masses. The operator O_3 only violates baryon number and must be combined with a lepton number violating operator to contribute to nucleon decay. If the lepton number violating dimension four operator $[LQ\bar{d}]_F$ is allowed the dominant contribution comes from the O_3 operator component involving the vacuum expectation value, v , of the Higgs field H . In this case the diagram giving proton decay is similar to that shown in Fig. 1 but the amplitude has an additional

suppression factor (v/M) so that the resultant bound is $\eta_3 \chi < 10^{-10}$ for $v \sim 100 \text{ GeV}$, still a significant constraint. On the other hand if the lepton number violating operator first occurs at dimension five through the operator O_4 the contribution to nucleon decay still comes from Fig. 1 but the amplitude is suppressed by two powers of (v/M) and imposes no significant constraint on the dimension 5 couplings. (This structure follows simply from a count of operator dimensions for nucleon decay driven by two dimension five operators must be inversely proportional in amplitude to two powers of M). Thus the only significant constraints which apply to the coefficients of the operators $O_3, \dots, O_{10}$ which violate *either* baryon- or lepton- number are those following from nucleon decay processes driven by a combination of dimension 4 and dimension five operators. In addition to the constraints on the pair $O_3 N_2$ discussed above there will be similar constraints on the products of the couplings of the pairs $O_4 N_3, O_9 N_3, O_{10} N_3$. We remind the reader that N_1, N_2, N_3 correspond to the last three operators in eq.(3).

The bounds on dimension five operators could be satisfied without appealing to (approximate) symmetries if, for some reason, dimension five couplings were anomalously small. In compactified string theories there is a possible explanation for the smallness of the coefficients of operators dimension five and above for such higher dimension terms only arise through non-perturbative string world-sheet effects and are suppressed by a factor proportional to $\exp(-a M_{\text{String}}^2 / M_{\text{Compactification}}^2)$ where a is a constant of order one. However the compactification scale and the string scale must be quite close together [12] so it is unlikely that this suppression factor will be small enough to suppress the dimension 5 terms to be consistent with the mentioned bounds. Thus it is likely that a realistic theory must have additional symmetries capable of suppressing such terms and for this reason we will extend our analysis of symmetries restricting the rate of nucleon decay to include a consideration of those symmetries capable of eliminating at the compactification scale the operators of equations (6) to be consistent with proton stability. We note that the pattern of fermion masses also requires that some Yukawa couplings of eq.(1) be quite small, possibly as the result of (approximate) symmetries too. However the scale of suppression required for these Yukawa couplings is significantly smaller than that required for nucleon stability so we consider it a reasonable first step to require a pattern of symmetries which give the latter while imposing no constraint on the former. The elimination of both dimension four and dimension five (plus four) contributions to nucleon decay seems to us a reasonable constraint to demand of a viable theory descending from the string but, in view of the uncertainties associated with the physics at such scales, we will also consider the weaker constraint that only dimension 4 operator contributions to nucleon decay should

be eliminated.

2.3 Baryogenesis constraints

There is a further necessary ingredient of a viable supersymmetric version of the standard model, namely it should be capable (at least in principle) of producing the observed baryon asymmetry of the universe. This is an interesting constraint for, unlike the above considerations, it requires the presence of baryon-number violation. At first sight this might seem to favour the presence of new sources of baryon-number violation coming from dimension-4 operators but as has been shown in references [13], [14] quite the opposite is true. In most schemes baryon number production proceeds through the decay, out of equilibrium, of heavy states coupled to both leptons and quarks. If a new source of baryon- or lepton- number violation is present in the low-energy Lagrangian then the excess of baryon or lepton number produced by these decays will be *reduced* by these new interactions. As discussed in refs. [14] the dimension-4 operators generate interactions which, for coefficients of $O(1)$, remain in equilibrium down to temperatures of order the scale of the new supersymmetric states and so would effectively erase any baryon- or lepton-number excess produced at temperatures above this scale. When coupled with the non-perturbative (B-L) violating processes of the standard model the effect of *any* dimension-4 baryon- or lepton- number operator is to eliminate *all* baryon- and lepton- number excess independently of the flavour component of this excess [13], [14]. Thus the authors of ref.[14] establish stringent bounds on dimension 4 operators.

However there are many alternative possibilities for baryon- number production which do not lead to a constraint on dimension- 4 operators. In particular it is possible that the out-of- equilibrium baryon-number generating decays occur at a low temperature and the reheating temperature does not reach the scale of masses of the supersymmetric states [15]. In this case the dimension-4 operator interactions are not in equilibrium after baryogenesis and do not significantly reduce the baryon excess. Alternatively baryogenesis may proceed via non-perturbative processes such as the appearance of squark condensates [16] or through the non-perturbative baryon-number processes [17], [18] of the standard model itself. Again these mechanisms produce a baryon number excess at low temperatures, of order of the electroweak breaking scale, at which the dimension-4 operators are not efficient at reducing the asymmetry. It is even possible that baryogenesis occurs at high temperature without stringent bounds on all flavours of dimension-4 operators. This follows because the flavour violation of the quark mass matrix is necessary to ensure the erasure of the baryon number excess in any flavour channel starting with a specific flavour

structure of the dimension-4 flavour operator. However this flavour violation may disappear at high temperatures eliminating the effectiveness of these operators in reducing the baryon number excess. (Whether this happens depends on the details of the electroweak breaking sector [19]).

In view of these uncertainties we feel it premature to use baryogenesis to argue against the presence of a significant component of new dimension-4 baryon- or lepton- number violating operators. Perhaps a better view to take is to note that if such a dimension-4 operator should be found it would limit the possible mechanisms for baryogenesis. In this context we note that the structure of dimension-5 operators also limits the mechanisms for baryogenesis. If baryogenesis is due to out-of-equilibrium heavy particle decay then on integrating out the heavy fields to obtain the effective low energy lagrangian there must be left baryon number violating operators, typically of dimension-5. If there are symmetries in the effective low energy theory which forbid all such operators there can not have been the heavy fields with baryon number couplings needed to give a baryon number asymmetry, unless the decay of the heavy state is via a high dimension (> 3) term in the superpotential for then the corresponding operator in the low energy theory will also have a large dimension (> 5).

3. Generalized Z_N Parities

In the previous section we have discussed the different experimental limits on the possible additional terms one may add to the effective Lagrangian through dimension four and five operators. Of course, the presence of some operators but not of others is rather artificial and unnatural unless there is some underlying symmetry which justifies it. We are going to study in this chapter the simplest such type of symmetries, the Z_N (generation independent) generalized parities (GP). Such type of symmetries have been studied in references [20]. Here we give a general classification of all GMP. We show that any possible GP Z_N symmetry acting in a generation independent way may be written in terms of only three independent Z_N generators. The explicit form of all independent Z_2 and Z_3 GP are given.

Let us consider a generation-blind Z_N symmetry acting on the five independent standard model superfields Q, u, d, L and e as

$$\Psi_j \rightarrow e^{\frac{i\alpha_j 2\pi}{N}} \Psi_j \quad j = Q, u, d, L, e. \quad (7)$$

One can describe each choice of Z_N by giving a vector of integers:

$$\vec{\alpha} = (\alpha_Q, \alpha_u, \alpha_d, \alpha_L, \alpha_e) \quad (8)$$

The charges of Higgs superfields $H, \bar{H}$ are determined by the fact their standard couplings to quarks and leptons should be allowed. What is the form of the most general such Z_N symmetry? Of course, in principle one may multiply by independent Z_N phases each of the five superfields and hence there are to start with $N^5 - 1$ non-trivial possibilities. This corresponds to a Z_N^5 freedom altogether. However, many of those are not independent and the number is reduced by the following considerations:

i) All Z_N elements which are related by discrete hypercharge rotations $\exp(i2\pi Y/N)$ are equivalent. Here Y corresponds to the integer hypercharge assignments (six times the usual ones). This corresponds to an invariance under shifts

$$\vec{\alpha} \rightarrow \vec{\alpha} + (1, -4, 2, 3, -6) \mod N \quad (9)$$

of the corresponding discrete symmetries. In the case of Z_2 this is equivalent to a simultaneous flipping of the Z_2 assignments of Q and L . In the case of Z_3 it corresponds to simultaneous multiplication of Q, u, d by $\alpha, \alpha^2, \alpha^2$ ($\alpha = \exp(i2\pi/3)$) respectively. The equivalence under (9) allows one to choose one of the phases (e.g., the one of Q) trivial.

ii) Since the Higgs fields H with hypercharge $+1/2$ should couple simultaneously to Qd and to $L\bar{e}$ (in order that both have masses), one must have

$$\alpha_Q + \alpha_d = \alpha_L + \alpha_e. \quad (10)$$

This means that one of the four remaining independent assignments α_j may be written in terms of the other three.

iii) There may also be extra accidental equivalences if we restrict ourselves to operators of certain dimensions. In particular, superpotentials of order M (leading to operators of dimension $M+1$) are invariant under a Z_M symmetry which multiplies all superfields by the same $\exp(i2\pi/M)$ phase. This corresponds to an equivalence under the Z_M shifts

$$\vec{\alpha} \rightarrow \vec{\alpha} + (1, 1, 1, 1, 1) \quad (11)$$

which is only valid for superpotential terms of order M . However, Z_M invariances of this type are broken by operators of dimension $\neq M$. Then, Z_N symmetries which are

equivalent under the shift (11) do not necessarily lead to the same physics, since their superpotential may differ by terms of order $\neq M$. We will thus consider symmetries related by such operations as inequivalent in what follows.

Due to the restrictions i) and ii), the number of independent phases is only N^3 , corresponding to a Z_N^3 freedom. It is now obvious that any Z_N symmetry acting on a single family may be written in terms of three independent discrete generators. For reasons which will become evident later on we will choose three generators R_N , A_N and L_N whose generating charges $\bar{\alpha}_N$ are shown in Table 1. They may in fact be written as discrete Z_N elements of some particular continuous $U(1)$ symmetries:

$$R_N = e^{\frac{2\pi i R}{N}}, \quad A_N = e^{\frac{2\pi i Y_A}{N}}, \quad L_N = e^{\frac{2\pi i L}{N}}. \quad (12)$$

Here I_3^R is (up to a minus sign) the third component of a right-handed weak-isospin (normalized to $\text{Tr}(I^R)^2 = 2$) and L is the standard (anti-)lepton number. The generator Y_A couples to quarks and leptons like the E_6 generator corresponding to the Cartan subalgebra of the $SU(2)$ appearing in the decomposition $E_6 \supset SU(6) \times SU(2)$. In the rest of this section we will ignore the possible connection of our discrete generators with any gauge symmetry. Any element g_N of a Z_N symmetry of the standard model can be written

$$g_N = R_N^m \times A_N^n \times L_N^p, \quad m, n, p = 0, 1, \dots, N-1 \quad (13)$$

The symmetries so obtained allow for different combinations of couplings in the effective Lagrangian. According to the type of dimension four operators allowed one can divide the symmetries into four classes:

a) Generalized Matter Parities (GMP). These are symmetries forbidding *both* baryon and lepton number violation in dimension four operators. The associated α_j charges must satisfy

$$\begin{aligned} \alpha_u + 2\alpha_d &\neq 0 \pmod{N} \\ 2\alpha_L + \alpha_e &\neq 0 \pmod{N}. \end{aligned} \quad (14)$$

Due to eq.(10) an additional condition is not required (the same is true for the rest of the generalized parities).

b) Generalized Lepton Parities (GLP). These are symmetries allowing for baryon but

not lepton violation. Their charges satisfy

$$\begin{aligned} \alpha_u + 2\alpha_d &= 0 \pmod{N} \\ 2\alpha_L + \alpha_e &\neq 0 \pmod{N}. \end{aligned} \quad (15)$$

c) Generalized Baryon Parities (GBP). These are symmetries allowing for lepton but not baryon violation. Their charges satisfy

$$\begin{aligned} \alpha_u + 2\alpha_d &\neq 0 \pmod{N} \\ 2\alpha_L + \alpha_e &= 0 \pmod{N} \end{aligned} \quad (16)$$

d) The fourth type

$$\begin{aligned} \alpha_u + 2\alpha_d &= 0 \pmod{N} \\ 2\alpha_L + \alpha_e &= 0 \pmod{N} \end{aligned} \quad (17)$$

allows for both baryon and lepton violation and is experimentally excluded due to the observed proton stability.

The conditions (14)-(17) lead to different constraints on the powers m, n, p in eq.(13) for the different generalized parities. Specifically, one can write the different generators in the general form of eq.(13) with the m, n, p constrained as follows

a) A generic GMP element (for any N) may be written

$$g_{GMP} = (R^2 A)^n R^{(m-2n)} L^p, \quad m-2n \neq 0; \quad m-2n-p \neq 0 \quad (18)$$

b) A GLP element:

$$g_{GLP} = (R^2 A)^n L^p, \quad p \neq 0 \quad (19)$$

c) A GBP element:

$$g_{GBP} = (R^2 A)^n (RL)^p, \quad p \neq 0 \quad (20)$$

d) A generic element g_{PQ} which does not forbid any dimension four operator at all is given by

$$g_{PQ} = (R^2 A)^n. \quad (21)$$

Notice that in all cases there is an $(R^2 A)^n$ factor. This corresponds to the generators leaving invariant all the dimension four operators of the theory, whatever they are. The

charges $\tilde{\alpha}PQ$ corresponding to this R^2A factor are also shown in Table 1. If this symmetry is promoted to a continuous one, one can see that it is just a Peccei-Quinn symmetry (notice it has QCD anomalies). This PQ symmetry is a symmetry of the dimension four operators of any supersymmetric standard model with minimal particle content. This continuous symmetry should not be an exact symmetry of the full Lagrangian since the Higgs superfields $H, \bar{H}$ are also charged with respect to it and lead on spontaneous symmetry breaking to a Weinberg-Wilczek axion in the spectrum, a possibility which is ruled out experimentally. In the minimal SUSY-standard model this PQ symmetry is explicitly broken by the Higgsino mass term $\mu H\bar{H}$. In models with extra particles (e.g., an extra singlet coupling to $H, \bar{H}$ replacing the $\mu H\bar{H}$ term) one has to check that the continuous PQ symmetry is explicitly broken. However there is no problem with a residual discrete version of this symmetry like $(R^2A)^n$ (apart from the possible cosmological creation of domain walls which may be diluted by appropriate inflation).

As we have noted the different generalized parities fall into conjugacy classes with respect to the subgroup generated by the “axionic” discrete symmetry $(R^2A)^n$. All GLPs are generated by a single generator L (modulo $g\bar{P}Q$) and all the GBP are given by a single generator RL (also modulo $g\bar{P}Q$). Then it is obvious that, for any given Z_N GLP or GBP the number of independent possibilities is $N \times (N-1)/2$ for N odd and $N^2/2$ for N even. In the same way one gets for the number of independent GMPs $N(N-1)^2/2$ for N odd and $N^2(N-1)/2$ for N even. Thus, for any given N there are altogether

$$\begin{aligned} \#GP &= \frac{N(N+1)(N-1)}{2} \quad (\text{odd } N) \\ \#GP &= \frac{N^2(N+1)}{2} \quad (\text{even } N) \end{aligned} \quad (22)$$

non-trivial generalized parities (plus the N elements generated by the “axionic” subgroup $(R^2A)^n$). These numbers are surprisingly small, taking into account the initial arbitrariness.

We show in Table 2 the independent Z_2, Z_3 generalized parities. The reader may easily find the equivalent results for higher N by using equations (13) and (18)-(21). In Table 3 we show what dimension five operators are allowed or forbidden by each of the Z_2 and Z_3 generalized parities. We also show which of them allow for a Higgs mass term $\mu H\bar{H}$. Let us make some general comments about the table. There are two Z_2 and six Z_3 GMP. The first of them, R_2 , is equivalent through a Z_2 hypercharge rotation to the standard R -parity of the minimal SUSY standard model. Furthermore, there are five GLPs and five GBPs. It is

important to note that out of these symmetries, only three GMPs, two GLPs and another two GBPs are consistent with the minimal particle content of the SUSY standard model. The reason is that the other generalized parities forbid the Higgs mass term $\mu H\bar{H}$. This coupling is needed to allow for an acceptable pattern of electroweak symmetry breaking in which both Higgs fields acquire vevs without an (experimentally excluded) light Higgs. Only the Z_2 symmetries R_2, L_2, R_2L_2 and the Z_3 symmetries $R_3, R_3L_3^2, L_3$ and R_3L_3 are consistent with the existence of such a mass term. Note that, for any N , it is the generators including non-vanishing powers of the “axionic” elements R^2A which forbid the Higgsino mass terms. Thus for each Z_N , a fraction $(N-1)/N$ of the generalized parities do forbid the needed mass terms.

The symmetries forbidding the $\mu H\bar{H}$ term can be made consistent with appropriate electroweak symmetry breaking by adding an extra singlet chiral field S to the spectrum with a coupling

$$\lambda_S SH\bar{H}. \quad (23)$$

If S acquire a vacuum expectation value the $\mu H\bar{H}$ is generated. This scheme by itself is not enough because, as we discussed above, in the absence of dimension two terms, the R^2A symmetry is promoted to a full global continuous Peccei-Quinn symmetry. Furthermore, dimension five operators are typically suppressed too much to help in giving a mass to the would-be axion. Thus new couplings breaking the continuous PQ symmetry while leaving the discrete symmetry intact should be added (for example a term MS^2 for $N=2$ or a cubic λS^3 term for $N=3$).

Notice that the GMP R_2 and R_3 respect $B-L$ and allow for proton decay through the dimension five operators $QQQL$ and $uude$. This is in fact true for all the R_N symmetries since R_N can be considered as a discrete version of gauged $B-L$ (see below). As we discussed in the previous section the coefficient of these operators has to be $\leq 10^{-7}$. If one wishes that this degree of fine-tuning should be avoided, one should then discard all the R_N discrete GMPs. In general, the conditions for these dimension five operators to be absent are $n+p \neq 0, \text{ mod } N$ (for $QQQL$) and $n-p \neq 0, \text{ mod } N$ (for $uude$). Here n, p are defined by equation (13).

Other seemingly innocent dimension five operators may induce too fast proton decay. Thus, for example, the Z_3 GBPs generated by $R_3A_3L_3^2$ allows for the presence of the dimension five operator $QQQH$ as well as for the lepton violating QdH . The combination of these two operators leads to too fast proton decay, as discussed in section 2. The symmetry $R_3A_3L_3^2$ is thus ruled out of the list of possible GPs.

If we adopt the more stringent requirement of absence of dimension five nuclear decay operators, the number of possibilities is much reduced. From table 3 we see that only one GMP (i.e., $R_3 L_3^2$, not the usual R-parity!), and two each of the GLP (L_2, L_3) and GBP ($R_2 L_2, R_3 L_3$) allow the $H\bar{H}$ term and eliminate nuclear decay coming from the dimension-5 terms. These favoured supersymmetric extensions of the supersymmetric standard model do not include the MSSM so their phenomenology may differ significantly from that normally considered. Notice that the above unique GMP is perhaps too effective: it forbids all dimension 5 operators violating B or L and hence it could be problematic to generate a baryon number asymmetry within that model.

Another interesting question is the presence of the operator $LL\bar{H}\bar{H}$. This is important since, once the Higgs fields get a vev, operators of this type are the ones which lead to Majorana neutrino masses for the left-handed neutrinos. The general condition (for any N) for this operator to be present is $2(m-n-p) = 0, \text{ mod } N$. Obviously, all Z_2 GPs allow for such operators. That is not the case for the Z_3 (and in general for the Z_N) GPs. If the presence of neutrino masses was confirmed (as might be indicated by the solar neutrino data) there should be a violation of lepton number at some level of the theory. GPs like the R_N for $N \geq 3$ would then be ruled out since they do not allow for lepton violation at sufficient rate. This may be seen to be the case also for many of the Z_3 symmetries in table 3: the symmetries $R_3, R_3 A_3 L_3, R_3 L_3^2, A_3, L_3$ and $R_3 A_3^2 L_3$ all forbid neutrino masses through the dimension five operator $LL\bar{H}\bar{H}$. The case of the GBPs is more tricky: these symmetries allow for lepton violating couplings QdL and LLe which induce neutrino masses through one loop diagrams [21]. Thus GBPs symmetries are perfectly consistent with the existence of neutrino masses. Furthermore, it has been recently shown that there might exist effective neutrino oscillations in the solar core mediated by lepton violating couplings of this type, even in the absence of any neutrino vacuum mixing [22].

We have also studied the case of the Z_4 parities but we refrain from presenting the relevant tables here in order not to overwhelm the reader. Let us just mention that the general structure of the obtained symmetries is similar to the one of Z_2 or Z_3 . There are altogether 24 GMPs, 8 GLPs and 8 GBPs. Out of them only 6 GMPs, 2 GLPs and 2 GBPs allow for Higgsino mass terms. Some of them do in fact correspond to Z_2 symmetries so that there are in fact only 5 GMPs, 1 GLP and 1 GBP new symmetries coming from Z_4 and allowing for Higgsino mass terms. This again shows how the various equivalences severely limit the number of independent possibilities.

In principle one could go to Z_N symmetries with arbitrarily large N . Apart from the lack of aesthetic appeal of such large Z_N symmetries, one may argue, at least in the

context of strings, that not very large N s are expected. The reason is as follows. Discrete symmetries in strings are related to the level of the conformal field theories used in the string compactification, which is itself limited. On the other hand, the discrete symmetry may be a remnant of a spontaneously broken continuous gauge symmetry. In this case the order of the discrete symmetry is limited by the Higgs field responsible for the breaking. The representation of the Higgs field itself is constrained to be small in string theories.

4. Discrete R-symmetries .

In the previous section we have considered discrete Z_N symmetries commuting with supersymmetry. Those require e.g. the invariance of the superpotential W (and the superspace measure $d\theta^2$) with respect to the given symmetry. We will study now the case of the Z_N R-symmetries. In this case the superspace chiral measure and the superpotential transform as

$$\begin{aligned} d\theta^2 &\longrightarrow e^{\frac{i2\pi\alpha}{N}} d\theta^2 \\ W &\longrightarrow e^{-\frac{i2\pi\alpha}{N}} W \end{aligned} \quad (24)$$

where α is an integer. We hope the reader will not confuse the notion of R-symmetry here defined with the notion of R-parity described in the previous section.

An analysis along the lines of the previous section can be easily done. It is obvious that a general (generation blind) Z_N R-symmetry of the susy-SM may be written as

$$\hat{g}_N = P_N \times R_N^m \times A_N^n \times L_N^p \quad (25)$$

where P_N is defined by $P_N d\theta^2 = e^{\frac{i2\pi}{N}} d\theta^2$ (leaving invariant the matter superfields). All these R-symmetries can be classified according to what dimension four Yukawa couplings they allow or forbid. The possibility of discrete hypercharge rotations and eq. (10) considerably reduce the number of possibilities. We further assume that the standard Yukawa couplings to quarks and leptons are always allowed, i.e. by definition one chooses

$$\begin{aligned} \alpha_H &= -\alpha_Q - \alpha_d - 1 = -\alpha_L - \alpha_e - 1 \\ \alpha_{\bar{H}} &= -\alpha_Q - \alpha_u - 1 \end{aligned} \quad (26)$$

Notice that this sometimes implies non-conventional assignments of discrete quantum numbers for the Higgs multiplets. In doing a general classification is then important the fol-

lowing observation. The generator

$$I \equiv P_N \times R_N^{-1} \quad (27)$$

always leaves invariant the dimension four couplings coming from the superpotential, as the reader may check. In fact this is nothing but a discrete subgroup of the continuous R-symmetry which is always present in globally supersymmetric theories with only dimension four couplings. We also know from the previous chapter that the generators $(R_N^2 A_N)^n$ also leave invariant those couplings and hence one can easily convince oneself that the general R-symmetry

$$\tilde{g}_{PQ} \equiv P_N \times R_N^{-1} \times g_{PQ}^r ; \quad r = 1, \dots, N \quad (28)$$

is the most general one leaving invariant the dimension four couplings. These elements play the role for R symmetry that g_{PQ} played in the previous case of the Z_N symmetries commuting with supersymmetry. As before the Z_N R-symmetries may be classified according to which terms are allowed or forbidden:

a) R-symmetries forbidding both baryon and lepton-violating dimension four superpotential couplings (Generalized R-Matter Parities, GRMPs). Their general form is

$$\tilde{g}_{GRMP} = \tilde{g}_{PQ} R^{(m-2n)} L^p, \quad m - 2n \neq 0 ; \quad m - 2n - p \neq 0. \quad (29)$$

b) R-symmetries allowing for baryon but not lepton violation (Generalized R-Lepton Parities, GRLPs). A general element is

$$\tilde{g}_{GRLP} = \tilde{g}_{PQ} L^p, \quad p \neq 0. \quad (30)$$

c) R-symmetries allowing for lepton but not baryon violation (Generalized R-Baryon Parities, GRBPs). A general element is

$$\tilde{g}_{GRBP} = \tilde{g}_{PQ} (RL)^p, \quad p \neq 0. \quad (31)$$

In all these expressions $\tilde{g}_{PQ}$ is given by any of the generators obtained giving values to r in eq. (28).

Due to the presence of the P_N generator in eq.(28) the number of independent Z_N R-symmetries for a given N is bigger than the corresponding number of ordinary Z_N symmetries. Indeed there are $N \times (N-1)$ for N odd (N^2 for N even) independent GRLPs and GRBPs. Concerning the GRMPs one gets $N(N+1)(N-1)$ for N odd ($N^2(N+1)$ for N even). As is obvious from eqs. (29),(30),(31), they are related to the usual Z_N symmetries by multiplication by the element I in eq.(27).

For the convenience of the reader we give the list of independent Z_2 and Z_3 R-symmetries in tables 4-a and 4-b respectively. We also show in those tables the dimension five operators allowed or forbidden by each Z_2 or Z_3 R-symmetry. It is also shown whether each given symmetry allows for a Higgs mass term. We refrain from given an exhaustive analysis of all the possible phenomenological aspects of these R-symmetries. Let us just make a couple of general comments. One can easily check that in order for a mass term $H\bar{H}$ to be allowed, the corresponding R-symmetry must have $n = 1$, i.e. it must contain the generator A_N . Only ten out of the thirty possible Z_2 and Z_3 R-symmetries allow for Higgs mass terms. Another interesting observation is that there is a unique GRMP which allows for Higgs mass term but forbid the dimension five operators $QQQL$ and $uude$. It is the one given by the generator $P_3 R_3^2 A_3 L_3^2$. Finally, let us remark that four out of the six Z_3 GRBPs and one of the Z_2 are ruled out because they allow simultaneously for the existence of the operators QdL and $QQQH$ (or QeH) which combined give rise to fast proton decay.

5. Discrete gauge anomaly cancelations and "anomaly free" generalized parities.

We have shown in the previous two sections how the discrete symmetries and R-symmetries may be classified in terms of only three independent generators. Although the number of independent possibilities turns out to be much smaller than expected, there is still a substantial arbitrariness in choosing which symmetry or R-symmetry could be the one which guarantees sufficient proton stability in the supersymmetric standard model. Of course, the fundamental reason of this arbitrariness is the lack of any a priori restriction on discrete symmetries. Gauge symmetries on the contrary are substantially constrained by the requirement of cancellation of anomalies.

We have recently remarked [10] that, if the discrete symmetry is a remnant of a spontaneously broken continuous gauge symmetry, it should also obey certain "discrete anomaly" cancellation conditions. For the sake of completeness we briefly review here these "discrete

gauge symmetries" and their corresponding "anomalies". The simplest example of a discrete gauge symmetry is provided by an Abelian Higgs model with a $U(1)$ gauge symmetry broken by a vev of a scalar field ϕ with charge Nq [8]. If the rest of the particles in the spectrum have charges quantized in terms of q , there will be a residual unbroken discrete Z_N group. This is a discrete gauge symmetry. The "discrete gauge anomalies" discussed in ref.[10] are non-perturbative like Witten's $SU(2)$ anomaly. Thus, e.g., if there is a mixed $Z_N \times SU(3) \times SU(3)$ anomaly you expect that QCD instanton effects will generate terms violating the discrete Z_N symmetry explicitly.

As pointed out by Krauss and Wilczek [8], [9], "discrete gauge symmetries" have the interesting property that are stable under quantum gravity (wormhole) effects. This is unlike usual non-gauge discrete symmetries which are explicitly broken by these quantum gravity fluctuations. This is important since, e.g., if the observed stability of the proton in a given model depends on the existence of a non-gauge discrete symmetry (like the ones studied in previous chapters), one generically expects that quantum gravity effects will violate it and will induce fast proton decay [7]. It then would be desirable the use of discrete gauge symmetries to guarantee that the structure of allowed and forbidden couplings is respected by the ever present quantum gravity fluctuations.

An additional appeal of discrete gauge symmetries is that discrete symmetries in four-dimensional string models have typically a gauge origin. That is for example the case of the abundant Z_N symmetries in orbifold models which are subgroups of $U(1)$ s which are spontaneously broken to Z_N s when one moves away from the "multicritical" (enhanced symmetry) points [23]. This phenomenon does also occur in other type of four-dimensional strings like those based in the tensoring of $N = 2$ superconformal field theories (and their associated Calabi-Yau compactifications). Thus it seems reasonable to expect the existence of residual unbroken gauge Z_N symmetries in the low energy effective actions of four-dimensional strings.

The interesting point is that if we insist on all the relevant discrete symmetries (like the generalized parities) being of gauge origin, it can be shown that the light fermions must obey some restrictive "discrete anomaly cancellation conditions" [10]. Let us assume there is an effective gauge Z_N symmetry under which the massless fermions of the theory transform with charges q_i . There are several types of "discrete anomalies" which must be cancelled by appropriately choosing the chiral fermion content of the theory. These are the following:

i) Cubic Z_N^3 anomaly cancellation condition:

$$\sum_i (q_i)^3 = r N + \eta \frac{N^3}{8}, \quad r, s \in \mathbb{Z}. \quad (32)$$

where $\eta = 1, 0$ for N =even, odd.

ii) Mixed Z_N -gravitational anomalies:

$$\sum_i (q_i) = r' N + \eta \frac{s'}{2}, \quad r', s' \in \mathbb{Z}. \quad (33)$$

iii) Mixed Z_N - $SU(M)$ - $SU(M)$ anomalies:

$$\sum_i T_i(q_i) = \frac{1}{2} r'' N, \quad r'' \in \mathbb{Z}. \quad (34)$$

Here T_j is the quadratic $SU(M)$ Casimir corresponding to each given representation (the normalization is such that the Casimir of an M -plet is $= 1/2$). The sums are over the light sector of the theory (those fields which do not get a mass when the symmetry breaking $U(1) \rightarrow Z_N$ takes place).

In the case there are extra unbroken abelian factors, $U(1)_a$ s, in the theory one can write down further constraints associated with mixed $U(1)_a^2 Z_N$ and $U(1)_a Z_N^2$ anomalies. However [10] they turn out not to be very useful without the precise knowledge of the underlying theory before the symmetry breaking $U(1) \rightarrow Z_N$ took place. If, on the other hand, there are two discrete symmetries Z_N and Z_M present, there are extra conditions from the cancellation of mixed $Z_N^2 Z_M$ anomalies. These yield

$$\sum_i (q_i)(p_j)^2 = t N + u M + s'' \eta_q \eta_p \frac{NM^2}{8}, \quad t, u, s'' \in \mathbb{Z} \quad (35)$$

where η_q (η_p) $= 1, 0$ for M (N) even, odd. Here q_j and p_j correspond to the Z_N and Z_M charges respectively.

All the above conditions look like a discretized version (i.e. mod N) of the usual continuous anomaly cancellation conditions. There is however an important difference provided by the last term in equations (32), (33) and (35). This term is associated with the possibility that at the symmetry breaking stage $U(1) \rightarrow Z_N$ some fermions get a large *Majarana* mass and gives the contribution of these massive fermions to the anomaly. Thus these extra terms may in some sense be interpreted as "discrete Wess - Zumino"-like terms.

An interesting question is which (if any) of the discrete symmetries described in the previous chapters fulfill the criterion of being “discrete anomaly free” and hence stable under quantum gravity effects. Already from the outset one can rule out a complete class of discrete symmetries. The R-symmetries studied in the previous chapter cannot have a gauge origin because any gauged continuous version of a discrete R-symmetry necessarily violates $(N = 1)$ supersymmetry. Particles belonging to the same multiplet would couple differently to the gauge bosons and hence gauge invariance would not commute with supersymmetry. Although this is possible in extended supersymmetries it is not possible in the $N = 1$ case we are considering here. (The situation concerning a gauge origin for the discrete R-symmetries may change in the case of effective theories from strings, as we discuss in the next section).

It is very easy to find the constraints on a generic discrete symmetry $g = R_N^m A_N^n L_N^p$ for any N coming from the above constraints. Let us consider the conditions coming from cancelation of mixed gravitational, mixed Z_N - $SU(3) - SU(3)$ and mixed $Z_N - SU(2) - SU(2)$ anomalies. We choose these because they are linear in the Z_N charges. Using the definitions in table 2 the reader may easily arrive at the following three conditions:

$$\begin{aligned} m N_g + n N_D &= \eta s \frac{N}{2}; \quad s \in \mathbb{Z} \\ n N_g &= 0 \\ p N_g - n N_D &= 0, \end{aligned} \quad (36)$$

which should be obeyed modulo N . Here the integers m, n, p correspond to the exponent of the generators R_N, A_N, L_N in the general expression of the symmetry element eq. (13). N_g is the number of quark-lepton generations involved in the cancelation of the anomalies and N_D is the number of Higgs sets $H, \bar{H}$.

Let us first look for solutions in which the cancelation of discrete anomalies takes place generation by generation (i.e., $N_g = 1$). In this case necessarily $n = p = 0, \text{ mod } N$. Only the generator R_N is left and hence one can choose $m = 1$. The first of the equations (36) may only be true for $N = 2$. Thus we reach the first important conclusion that there is a unique anomaly free discrete symmetry in which cancelation of anomalies takes place generation by generation. It is just R_2 , the standard R-parity of the minimal supersymmetric model. The reader can check that the last (discrete Wess-Zumino-like) terms in equations (32) and (33) play a crucial role in the anomaly cancelation for R_2 .

The other class of solutions correspond to having $N = N_g$. The phenomenologically interesting case has $N_g = 3$ so that in this case one is necessarily limited to Z_3 generalized

parities. The two first equations (36) are automatically verified but the third is only true if either $n = 0$ or the number of doublets is at least three (since necessarily $N_D \geq 1$). Let us neglect for the moment the latter possibility since it implies an enlargement of the particle content of the SUSY standard model (we will consider the extensions below). One is thus left with the Z_3 symmetries $R_3, R_3 L_3^2, L_3$ and $R_3 L_3$. In fact computing the cubic Z_3^3 anomaly one finds that *only* the GBP $R_3 L_3$ is anomaly free at this level (with the minimal low energy particle content of the SUSY standard model)*. Thus one can state the following:

“ There are only two generalized parities which are discrete anomaly free with the minimal particle content of the SUSY standard model. One is the standard Z_2 R-parity and the other is the Z_3 symmetry generated by $B_3 = R_3 L_3$ which allows for lepton number violation ”

This degree of uniqueness is remarkable. If combined with the requirement that nucleon decay mediated by dimension-5 operators should be suppressed (cf. section 3 and table 3) we find a *unique* possibility, namely the lepton violating symmetry $B_3 \equiv R_3 L_3$. For this reason we think that this “*Baryon parity*” symmetry B_3 should also be taken seriously and their phenomenological implications fully studied. The phase assignment of matter fields with respect to B_3 are $q(Q,u,d,L,e) = (1, \alpha^2, \alpha, \alpha^2, \alpha^2)$.

More anomaly free generalized parities exist if we enlarge the particle content of the minimal SUSY standard model by adding extra massless chiral multiplets. But in this case the phenomenology of the model will be altered and the effect of the new particles should be taken into account. The simplest modification is the addition of extra singlets and/or $SU(2)$ doublets. Let us briefly discuss the new anomaly-free discrete symmetries one may obtain in this manner:

a) Extra right-handed neutrino states. If an extra right-handed neutrino state per generation exists, the GMPs generated by R_N^m are all anomaly free generation by generation ($N_g = 1$) and for any N . This is trivial to check from the above equations. These are the only anomaly free Z_N parities which exist for $N \geq 4$. These possibilities have however the shortcoming that relatively large neutrino masses are in general expected through the see-saw mechanism. From this point of view R_2 (the standard R-parity) is safer since it allows for large Majorana masses for the $\nu_{\mu\tau}$. These Majorana masses are forbidden for $N \geq 3$.

* In checking that the other possibilities are anomalous, it is crucial to realize that in the Z_3 case the coefficient r in equation (32) must be necessarily a multiple of three, as shown in ref.[10].

Apart from these solutions which are anomaly free generation by generation, there are two other Z_3 which are anomaly free if the three generations are considered ($N_g = 3$) and a singlet per generation is added. They are the GLP generated by L_3 and the GMP generated by $R_3 L_3^2$.

b) Extra sets of Higgs multiplets $H, \bar{H}$. Examining again equations (36) one sees that for $N = N_g = 3$ there are further solutions with $N_D = 3$, i.e., three sets of Higgses $H, \bar{H}$ and $n \neq 0$. One can see there are altogether two such Z_3 parities: one $GBP, A_3 L_3$ and one $GLP, R_3 A_3^2 L_3$. Models with three sets of Higgses have however a couple of shortcomings. First, one typically has problems with flavour changing neutral currents, unless the extra Higgses are substantially heavier than the standard ones. Second, assuming the "big desert" hypothesis, the prediction for the Weinberg angle (presently in amazingly good agreement with recent LEP data [24]) is completely spoiled.

c) Both extra singlets and Higgs multiplets. One can easily convince oneself that the rest of the Z_3 symmetries can be made anomaly free by the addition of both extra singlets and extra Higgs doublets. These are $R_3 A_3, A_3 L_3^2, R_3 A_3^2 L_3, A_3 L_3, R_3 A_3 L_3, A_3$ and $R_3 A_3^2 L_3$.

Notice that the additional singlets we are discussing are charged with respect to the Z_N symmetries in order that they can contribute in cancelling the anomalies. There could be other extra singlets inert with respect to Z_N or even pairs of singlets with conjugate Z_N charges not contributing to the anomalies. These new singlets could be needed for other phenomenological reasons, not connected to anomaly cancellation.

Let us remark that the above list is the complete list of anomaly free Z_N generalized parities. Notice that there are only two of them (R_3 and B_3) which are anomaly free with the minimal particle content. The rest of the Z_3 symmetries can be made anomaly free by adding extra singlets and/or doublets. But this enlargement of the spectrum may lead to some phenomenological problems with neutrino masses, flavour changing neutral currents or the Weinberg angle prediction. For $N \geq 4$ only the R_N symmetries can be made anomaly free by adding some right-handed neutrinos. Of course, by the addition of a sufficiently large number of exotic particles one will be able to "gauge" more and more discrete symmetries but then it will be hard to understand the apparent success of the minimal standard model.

Another important point is that in all the anomaly free Z_3 symmetries discussed above the number three is indeed the number of generations, i.e., one should really write Z_{N_g} . In this way the stability of the proton and the presence of R-parity violation is connected to the number of generations in an intriguing way. We will show below examples of gauged

$U(1)$ symmetries which are only anomaly free for three generations and lead when broken to a Z_3 symmetry of the type considered above.

6. Baryon and lepton number violation and gauged $U(1)$ s.

The idea of a discrete gauge Z_N symmetry which is the remnant of an original anomaly-free gauged $U(1)$ leads us naturally to the question of the possible role of standard continuous gauge symmetries in guaranteeing proton stability. * This problem was addressed in ref.[23] for the case of $U(1)$ s forbidding all terms violating lepton or baryon number (i.e. continuous generalizations of the Generalized Matter Parities). We will consider in more detail here the case in which e.g. only lepton or baryon number are conserved.

An additional interest in considering continuous $U(1)$ s is trying to find out which $U(1)$ s correspond to each of the Z_N gauge discrete symmetries considered in the previous chapter. Let us recall that the anomaly cancellation conditions (32),(33),(34) are *necessary* but not sufficient conditions that any discrete gauge symmetry coming from a broken $U(1)$ must fulfill. In some cases, as we shall describe below, it is easy to identify the existence of anomaly-free $U(1)$ extensions of certain gauged Z_N s.

The extension of the most general (generation blind) Z_N symmetry to a gauged $U(1)$ is trivial. The most general gauged $U(1)$ coupling in a generation independent manner to the standard model can be written

$$\Theta = aY + bI_3^R + cY_A + dY_L \quad (37)$$

where a, b, c, d are real coefficients and Y is the Weak hypercharge. The other three generators are the ones mentioned in eq.(12) and Table 1. The first three generators can be embedded into an E_6 group and have been abundantly discussed in the literature. Let us recall the (sometimes forgotten) fact that Weak hypercharge is the unique anomaly-free $U(1)$ which one can couple to one SM generation. In order to couple any of the other three generators in eq. (37) in an anomaly-free manner one has to add appropriate extra matter fields.

* This was first considered by S. Weinberg in ref.[3]. However, in that early reference the cancellation of mixed $U(1)$ -gravitational anomalies was not imposed.

In the case of I_3^R it is well known that anomaly cancellation requires the addition of one right-handed neutrino ν_L^c per generation. All the R_N discrete symmetries considered in the previous chapters are nothing but discrete versions of the $U(1)$ coupled to I_3^R , the third component of a right-handed weak isospin. In particular, as we remarked elsewhere, the usual R-parity of the minimal susy model is nothing but a Z_2 subgroup of this $U(1)$. Thus R-parity can be obtained as a discrete Z_2 gauge symmetry by breaking spontaneously I_3^R through the vevs of appropriate (lepton number = ± 2) scalars. In this case these scalars can give a mass to the right-handed neutrinos ν_L^c whose existence is required for I_3^R to be anomaly-free. Notice that this is not the case for R_N , $N \geq 3$ and that is why the standard R-parity is the only Z_N gauge symmetry which is anomaly free generation by generation and with the minimal content of the susy-SM (no light ν_L^c s required).

The generator Y_A may be embedded inside E_6 as the diagonal generator of $SU(2)$ in the branching $E_6 \supset SU(6) \times SU(2)$. In order to make it anomaly free the particle content of each SM generation has to be enlarged by pairs of fermions $L', \bar{L}', D, \bar{D}$ with the quantum numbers of L and d (and their complex conjugates). It is however not obvious what extra fermions and scalars one can add in order to obtain an unbroken discrete gauge symmetry once the scalars get a vev and break Y_A .

The case of a gauged lepton number generator Y_L has not, as far as we know, been studied in the literature and hence we will discuss it in this chapter in more detail. How can one promote the discrete generator L to a continuous anomaly-free $U(1)$ Y_L ? In the SM only the leptons L, e have non-vanishing values 1, -1 for this charge. Then it is obvious that there will be mixed $Y_L \times SU(2) \times SU(2)$ anomalies unless extra $SU(2)$ doublets are introduced to cancel them. The simplest solution is to add the multiplets (the Y_L charge is given in parenthesis)

$$L'(-1); e'(1); \bar{L}'(0); \bar{e}'(0); N(1); \bar{N}(-1) \quad (38)$$

to each quark-lepton generation. This also guarantees the cancellation of hypercharge anomalies. Giving vevs $\langle N \rangle = \langle \bar{N} \rangle > 0$ one can give masses to the extra matter fields while breaking the gauge symmetry. Unfortunately the dimension five operator $QLd\bar{N}$ and the dimension four coupling udd are allowed and would lead to fast proton decay unless $\langle \bar{N} \rangle > /M_{\text{Planck}} \leq 10^{-26}$ (or the coefficients of the operators are ridiculously small).

The origin of the problem is that the Y_L gauge symmetry is *completely* broken; there is no residual Z_N gauge symmetry. In fact the possibility in eq.(38) is not even the most economical one since it requires the introduction of six extra $SU(2)$ doublets. The minimal

model in which the lepton number may be gauged is the following. One introduces one right-handed neutrino ν_L^c per generation (lepton number = -1) plus the multiplets (in a single copy, no family replication)

$$L'(-3); e'(3); \bar{L}'(0); \bar{e}'(0); N(3); \bar{N}(-3); N'(3) \quad (39)$$

where $N, \bar{N}, N'$ are SM singlets. The reader can check that this model is anomaly-free provided the number of quark-lepton generations is three. There is a flat direction $\langle N \rangle = \langle \bar{N} \rangle > 0$ in which Y_L is spontaneously broken to a Z_3 discrete gauge symmetry. Furthermore the gauge symmetry allows for Yukawa couplings which give large masses to $L', e', \bar{L}', \bar{e}'$. The massless spectrum is just that of the standard susy model apart from the ν_L^c s and non-interacting singlets (N' and a linear combination of $N, \bar{N}$). The residual unbroken gauge Z_3 symmetry is just the one generated by L_3 in the previous section. Notice that, as we discussed at the end of the previous section, L_3 requires the presence of extra singlets in order to be anomaly-free. These are the three right-handed neutrinos. Notice also that in the cancellation of anomalies it is crucial that the number of generations is three, i.e. this is an example of the $Z_{N'}$ type of anomaly-free discrete gauge symmetries discussed in the previous section.

We consider this model as an interesting gauge extension of the SM model which deserves attention in its own right, in addition of providing us with a nice example of a model in which the number of generations and lepton number conservation are directly correlated. It has however the problem that cancellation of Z_3 anomalies forces the existence of relatively light right-handed neutrinos whose coupling to the ordinary neutrinos through a Dirac mass is not prevented by any symmetry.

Another interesting question is whether one can construct a model with an extra $U(1)_B$ forbidding all couplings violating baryon number and leading, upon symmetry breaking, to the preferred "baryon parity" $B_3 = R_3 L_3$ model discussed in the previous section. The answer is yes, although we have found more than one possible realization of this mechanism. We show for completeness a specific example of such type of model which is probably not the simplest one can find with the above characteristics but suffices as an existence proof. The $U(1)_B$ charges of the standard particles are given by $\alpha(Q, u, d, L, e) = (0, -1, 1, -1, -1)$. In order to cancel the mixed $U(1)_B \times SU(2)^2$ one adds extra heavy leptons $L'(3), \bar{L}'(0), e'(3), \bar{e}'(0)$ (the $U(1)_B$ charges are given in brackets). The cancellation of mixed $U(1)_B \times U(1)_Y^2$ and $U(1)_B^2 \times U(1)_Y$ anomalies may be achieved by adding leptons with $\pm 1/3$ electric charges: $2 \times F^{1/3}(3); 1 \times F^{-1/3}(3)$ plus $2 \times F^{-1/3}(0); 1 \times F^{1/3}(0)$. Finally,

the cancellation of Z_3^3 anomalies and the multiplets required for $U(1)_B$ symmetry breaking are given by $5 \times N^0(-3); N^0(1); N^0(2); N^0(3)$. The superindices indicate the hypercharge assignments. By giving vevs to $N^0(\pm 3)$ $U(1)_B$ is spontaneously broken to the preferred discrete 'baryon parity' B_3 . The particle content of the low energy theory is that one of the minimal susy SM since the fields $N^0(-3)$ may give masses to all extra fields through the allowed couplings $L\bar{L}N^0(-3)$; $e'\bar{e}'N^0(-3)$; $F^{1/3}F^{-1/3}N^0(-3)$; $N^0(1)N^0(2)N^0(-3)$. Thus in the present B_3 model (and unlike the lepton model with L_3 symmetry) no extra particles (like right-handed neutrinos) are required in the massless spectrum. This is again an explicit check of the anomaly cancellation arguments of the previous section. Let us comment that leptons with electric charge $\pm 1/3$ occur naturally in $SU(3)^3$ unified (or string models). The standard charges are obtained for particles in multiplets of the form $(1, 3, \bar{3})$ (and permutations) whereas particles with $\pm 1/3, \pm 2/3$ charges come in multiplets of type $(1, 1, 3)$.

The above two examples of gauged $U(1)$ s forbidding either lepton or baryon number violations (along with the ones forbidding *both* which were studied in ref.[23]) show that continuous gauged $U(1)$ symmetries are a viable means of guaranteeing proton stability in low energy supersymmetry. However, from the above analysis it is obvious that discrete Z_N gauge symmetries are most economic since they do not require (at least for R_2 and B_3) any extension of the massless particle content of the susy SM.

The case of discrete symmetries not commuting with supersymmetry (R-symmetries) is different. From the four-dimensional field theory point of view there is no way they could be gauge symmetries since the corresponding $U(1)$ s would not commute with $N = 1$ supersymmetry. On the other hand, if the four dimensional field theory is obtained from higher dimensions (as may be the case in Kaluza-Klein and string theories), there could be R-symmetries which are discrete remnants of the rotation group in the compactified dimensions (e.g., discrete subgroups of $SO(6)$ in the case of compactified heterotic strings). The cancellation of gravitational anomalies in the extra dimensions will then give rise to anomaly cancellation constraints analogous to the ones discussed in section 5 for the effective low energy R-symmetries. Furthermore, if an underlying string theory exists there could exist in the effective theory discrete versions of the Green-Schwarz anomaly cancellation mechanism. Also, string theories contain other (higher spin) continuous gauge symmetries (like those associated to two-index antisymmetric B_{mn} tensor) which could be spontaneously broken to discrete subgroups. The structure of discrete gauge symmetries in string models is intimately connected to the "target space duality" symmetries present in these models and are known to have important consequences in constraining the low

energy effective action from strings. We postpone these interesting aspects of discrete gauge symmetries in string models and the consideration of anomaly constraints on discrete R-symmetries to future work.

7. Generation dependent generalized parities.

Most of the emphasis in the previous chapters has been put on the case of *generation-blind* discrete and/or gauge symmetries. A similar analysis considering flavour dependent symmetries is obviously quite cumbersome since e.g. the number of independent Z_N generators goes like $N^3 \times N_g$, N_g being the number of generations. However it is obvious that generation dependent discrete symmetries are of potential importance for understanding the structure of quark-lepton mass matrices (Yukawa couplings). In this context the following scenario (to some extent realized in some specific string models) seems plausible. At a large (GUT or string) scale, M , the structure of Yukawa couplings is governed by a certain number of gauged symmetries (e.g. $U(1)_1 \times U(1)_2 \times \dots \times U(1)_m$). Some or all of those symmetries may in fact be discrete gauge symmetries (e.g. $Z_{N_1} \times \dots \times Z_{N_m}$). At a scale M_I slightly below M there are scalar fields which break those $U(1)$ s or Z_N s down to a "diagonal" gauged Z_N symmetry of the type discussed in the previous chapters guaranteeing sufficient proton stability. Since there is a hierarchy of scales M_I/M in the breakdown

$$U(1)_1 \times \dots \times U(1)_m \times Z_{N_1} \times \dots \times Z_{N_m} \longrightarrow Z_N \quad (40)$$

in addition to the unbroken discrete gauge symmetry Z_N there will survive *approximate* gauge Z_N symmetries (valid up to order M_I/M corrections). The latter will be reflected in the structure of Yukawa couplings and hence on the fermion mass matrices. A general analysis of this scenario and the constraints coming from anomaly cancellation is certainly very interesting but we postpone it for future research. Here we would like just to show one simple example of generation dependent gauge symmetry leading to an unbroken gauged $Z_2 \times Z_2$ symmetry. It is an interesting example because the massless particle content is just that of the susy standard model yet some new physics is present.

A motivation for this model is the development of the attempt discussed in the last section to find a model leading to an anomaly free gauged lepton parity, i.e., an anomaly free Z_2 under which leptons are odd and the rest of the spectrum is even. We furthermore require that the particle spectrum is the one of the susy SM (*without* right-handed neutrinos). In the previous section we have shown that this is not possible for a single Z_N discrete factor (except in the unphysical case $N_g = 2$). However, it is possible to find a

model with these characteristics and a $Z_2 \times Z_2$ gauge symmetry. Consider a $U(1) \times U(1)$ theory under which the leptons of each generation transform:

$$\begin{aligned} e &\longrightarrow L_e(1, 0); e(-1, 0); \nu_e^c(-1, 0); N_e^0(2, 0); \bar{N}_e^0(-2, 0); N_e^{0'}(2, 0) \\ \mu &\longrightarrow L_\mu(0, 1); \mu(0, -1); \nu_\mu^c(0, -1); \bar{N}_\mu^0(0, 2); \bar{N}_\mu^{0'}(0, -2); N_\mu^{0'}(0, 2) \\ \tau &\longrightarrow L_\tau(1, 1); \tau(-1, -1); \nu_\tau^c(-1, -1); N_\tau^0(2, 2); \bar{N}_\tau^0(-2, -2); N_\tau^{0'}(2, 2) \end{aligned} \quad (41)$$

where we have introduced extra singlets to obtain anomaly cancellation and do the symmetry breaking down to $Z_2 \times Z_2$. The cancellation of mixed anomalies require the addition of (one copy of) extra heavy leptons

$$L'(-2, -2); e'(2, 2); \bar{L}'(0, 0); \bar{e}'(0, 0) \quad (42)$$

In these equations the charges under $U(1) \times U(1)$ are shown in brackets. All the rest of the spectrum (quarks, Higgses) are singlets. The reader may check that all anomalies cancel. By giving vevs to the $N, \bar{N}$ fields the symmetry breaking $U(1) \times U(1) \rightarrow Z_2 \times Z_2$ is achieved. Furthermore, the extra states in eq.(42) get massive through their allowed couplings to $N^0, \bar{N}^0$. Notice also that the right-handed neutrinos (unlike the Z_3 example of previous chapter) may also acquire masses through their allowed Majorana type couplings to the $N^{0'}$ s. Thus the massless spectrum is just the one of the minimal susy standard model. The flavour structure of this particular model is as follows. There are no constraints on the quark mass matrices since quarks do not have lepton numbers and there is no constraint either on the udd couplings which are allowed in this model. The structure of the lepton mass matrices is diagonal and all lepton number violating dimension four couplings are forbidden except for the term

$$L_e \times L_\mu \times \tau_L^c. \quad (43)$$

The existence of this coupling is essentially unconstrained by present experimental limits. It can however lead to some amusing effects. To start with, of course, R-parity is not preserved and the Lightest Supersymmetric Particle (LSP) may decay like $LSP \rightarrow \tau^\pm \mu^\pm \nu_e$ or $LSP \rightarrow \tau^\pm e^\pm \nu_\mu$. For a sufficiently large coupling in eq. (43) the decays will occur inside the detectors. The sleptons have also direct decays

$$\begin{aligned} \tilde{e}_L &\longrightarrow \tau + \nu_\mu \\ \tilde{\mu}_L &\longrightarrow \tau + \nu_e \\ \tilde{\tau}_R &\longrightarrow e + \nu_\mu \text{ or } \mu + \nu_e \end{aligned} \quad (44)$$

which compete with the standard channel $\tilde{l} \rightarrow l LSP$. Couplings like that in eq.(43) can also lead to enhanced neutrino oscillations in solar matter even in the absence of vacuum

neutrino mixing (see ref.[22]). We hope this model exemplifies some of the interesting features which may be present in models with generation dependent discrete gauge symmetries.

8. Summary and conclusions.

In this paper we have investigated the symmetries, both continuous and discrete, capable of inhibiting proton decay in supersymmetric theories. Our motivation has been to construct viable supersymmetric extensions of the standard model, paying particular regard to the question whether the MSSM is the most reasonable version of the theory since the latter has been used almost exclusively in phenomenological studies. In the case of discrete symmetries it has been possible to study the general class of theory with a Z_N symmetry including the possible R-symmetries. Surprisingly there are very few possible theories once one demands the minimal light spectrum with the terms needed for masses. The matter-parity of the standard model is capable of eliminating nucleon decay coming from dimension-4 operators but allows dimension-5 operator contributions which, arguably, are still unacceptably large. There are discrete symmetries capable of eliminating both dimension-4 and dimension-5 contributions but these do not respect the R-parity of the MSSM.

There is a potential problem for any theory, supersymmetric or not, coming from large gravitational corrections which are thought likely to break any non-gauge symmetry. If this happens the dangerous dimension-4 and dimension-5 operators giving rise to fast proton decay will reappear, even if forbidden by a discrete symmetry, unless the discrete symmetry is a "discrete gauge symmetry", a remnant of a continuous gauge symmetry. We have investigated the possibility that the discrete symmetries may be left exact after the breaking of a continuous gauge symmetry. In the case of non-R-symmetries we found a "discrete gauge anomaly cancellation" condition which is necessary for a discrete gauge symmetry. Applied to the Z_N symmetries discussed above we found only two possible discrete gauge symmetries with the minimal supersymmetric multiplet content, the R-parity of the MSSM and a new "baryon parity". Only the latter eliminates the large dimension-5 contributions to nucleon decay.

We have also investigated the continuous gauge symmetries capable of leaving unbroken, on spontaneous symmetry breaking, these discrete gauge symmetries. This involves the construction of the massive sector of the theory and at this stage there is considerable ambiguity. We found examples in every case studied of continuous symmetries which did break to the desired discrete factor and which leave light just the fields needed by the

discrete anomaly cancellation condition. Of course there is also the possibility that the proton is stabilised by a continuous gauge symmetry which does not leave an unbroken discrete gauge factor. In this case, however, there are severe constraints on the scale of gauge symmetry breaking which in some cases are unacceptably low. We briefly reviewed these constraints and constructed examples of viable schemes which avoid the pitfalls.

Although we have emphasised the gauge symmetries, both discrete and continuous, we have also considered non-gauge discrete symmetries as a possible way to build a consistent supersymmetric model. In addition to the normal symmetries discussed above we have investigated R-symmetries as a way to stabilise the proton. We were able to classify models constructed using a single Z_N R-symmetry and found examples of viable models. Although apparently not gauge symmetries it is possible that they too descend from a gauge symmetry in a compactified string theory and we briefly discussed this possibility in section 6. If true, R-symmetries would be as attractive as other discrete gauge symmetries in stabilising the proton and would merit further study.

Clearly there are many possible ways to construct a supersymmetric version of the standard model, many with just the minimum supermultiplet content. The differences between these models lies in the allowed couplings and in this respect the MSSM is unusual in having no couplings which allow the LSP to decay. Given this diversity we consider it prudent, at the very least, to allow for the possibility of other versions when considering phenomenological signals for supersymmetry [5], [6]. In models in which the LSP decays within the detector the invisible energy of the (neutral) LSP is converted into visible energy in the form of its decay products. Thus the standard missing energy signals for supersymmetry no longer apply. Fortunately for models which include the $[LLe]_F$ and $[LQd]_F$ couplings the decay products typically contain charged leptons so there are new characteristic lepton-number-violating signals for supersymmetry which prove to be at least as visible as those of the MSSM [6]. In the case of models which include the $[udd]_F$ couplings the missing energy is converted into hadronic jets, but even in this case good signals for supersymmetry may be found by considering the cascade decays of the produced supersymmetric states [6]. In both cases, however, it is important to extend the experimental search patterns to include the new signals if one is to find these versions of the supersymmetric extension of the standard model.

Once one allows for R-parity violating models there is an embarrassingly large class of possible models. Our analysis of the possible symmetries capable of stabilising the proton has reduced this number considerably. From a purely aesthetic point of view two models appear most attractive being based on discrete gauge symmetries. They are

the MSSM using a matter parity to stabilise the proton and a new model which instead uses a discrete baryon parity. The latter has the advantage of eliminating dangerous dimension-5 contributions to nucleon decay and has the interesting property that single sparticle production and decay is possible with associated lepton-number-violating signals. Given the theoretical advantages of the discrete baryon parity model we think it should be considered at least on a par with the, by now, conventional MSSM when discussing the possible phenomenological signals of supersymmetry.

REFERENCES

- [1] H.P. Nilles, Phys.Rep. C110 (1984) 1;
G.G. Ross, "Grand Unified Theories", Benjamin Inc., (1984);
L.E. Ibáñez, "Beyond the Standard Model (yet again)", CERN-TH.5982/91, to appear in the proceedings of the 1990 CERN summer school (Mallorca).
- [2] S. Dimopoulos and H. Georgi, Nucl. Phys. B193 (1981) 150.
- [3] S. Weinberg, Phys. Rev. D26 (1982) 287;
N. Sakai and T. Yanagida, Nucl. Phys. B197 (1982) 533.
- [4] L. Hall and M. Suzuki, Nucl. Phys. B231 (1984) 419;
F. Zwirner, Phys. Lett. 132B (1983) 103;
R. Mohapatra, Phys.Rev.D34 (1986) 3457;
R. Barbieri and A. Masiero, Nucl. Phys. B267 (1986) 679;
S. Dimopoulos and L. Hall, Phys. Lett. 196B (1987) 135
V. Barger, G.F. Giudice and T. Han, Phys.Rev.D40 (1989) 2987.
- [5] S. Dimopoulos, R. Esmailzadeh, L. Hall and G. Starkman, Phys.Rev.D41 (1990) 2099;
S. Dawson, Nucl. Phys. B261 (1985) 297.
- [6] H. Dreiner and G.G. Ross, Oxford preprint OUTP-91-15P
P. Binetruy et al., Proceedings of the ECFA Large Hadron Collider (LHC) Workshop, Aachen, 1990. Vol.I, CERN report CERN 90-10.
- [7] G. Gilbert, Nucl. Phys. B328 (1989) 159 and references therein.
- [8] L. Krauss and F. Wilczek, Phys.Rev.Lett. 62 (1989) 1221.
- [9] T. Banks, Nucl. Phys. B323 (1989) 90;
L. Krauss, Gen.Rel.Grav. 22 (1990) 50;
M. Alford, J. March-Russell and F. Wilczek, Nucl. Phys. B337 (1990) 695;
J. Preskill and L. Krauss, Nucl. Phys. B341 (1990) 50;
M. Alford, S. Coleman and J. March-Russell, preprint HUTP-90/A040 (1990).
- [10] L.E. Ibáñez and G.G. Ross, CERN-TH.6000/91 (1991), to appear in Phys.Lett.B.
- [11] J. Ellis, J. Hagelin, D. Nanopoulos and K. Tamvakis, Phys. Lett. 124B (1983) 484.
- [12] V. Kaplunovsky, Phys.Rev.Lett. 55 (1985) 1036;
M. Dine and N. Seiberg, Phys. Lett. 162B (1985) 299.
- [13] M. Fukugita and T. Yanagida, Phys. Lett. 174B (1986) 45; Phys.Rev.D42 (1990) 1285;
A. Bouquet and P. Salati, Nucl. Phys. B284 (1987) 557;
J. Harvey and M. Turner, Phys.Rev.D42 (1990) 3344;
A. Nelson and S. Barr, Phys. Lett. 246B (1990) 141.
- [14] B. Campbell, S. Davidson, J. Ellis and K. Olive, CERN-TH.5926/90;
W. Fischler, G. Giudice, R. Leigh and S. Paban, Texas preprint UTTG-44-90 (1990).
- [15] S. Dimopoulos and L. Hall, Phys. Lett. 196B (1987) 135;
J. Cline and S. Raby, Ohio preprint DOE/ER/01545-444 (1990).
- [16] I. Affleck and M. Dine, Nucl. Phys. B249 (1985) 361;
D. Morgan, Texas preprint UTTG-04-91 (1991).
- [17] V. Kuzmin, V. Rubakov and M. Shaposhnikov, Phys. Lett. 155B (1985) 36.
- [18] P. Arnold and L. McLerran, Phys. Rev. D37 (1988) 1020;
S. Khlebnikov and M. Shaposhnikov, Nucl. Phys. B308 (1988) 885;
D. Grigoriev, V. Rubakov and M. Shaposhnikov, Phys. Lett. 216B (1989) 172.
- [19] H. Dreiner and G.G. Ross, Oxford preprint, in preparation.
- [20] L. Hall and M. Suzuki, Nucl. Phys. B231 (1984) 419;
M. Bento, L. Hall and G.G. Ross, Nucl. Phys. B292 (1987) 400.
- [21] R. Barbieri, M. Guzzo, A. Masiero and D. Tommasini, Phys. Lett. 252B (1990) 251.
- [22] M. Guzzo, A. Masiero and S. Petcov, SISSA preprint 16/91 EP (1991);
E. Roulet, Fermilab preprint 91/18-A (1991).
- [23] A. Font, L.E. Ibáñez and F. Quevedo, Phys. Lett. 228B (1989) 79.
- [24] J. Ellis, S. Kelley and D.V. Nanopoulos, Phys. Lett. 249B (1990) 441; Phys. Lett. 260B (1991) 131;
P. Langacker, "Precision tests of the standard model", Pennsylvania preprint UPR-0435T, (1990);
U. Amaldi, W. de Boer and H. Fürstenau, CERN-PPE/91-44 (1991);
P. Langacker and M. Luo, Pennsylvania preprint UPR-0466T, (1991).

Fig. 1-. Graph generating nucleon decay via dimension-4 operators. Solid lines denote fermions and the dotted line denotes the squark propagator.

Fig. 2-. Typical graph generating nucleon decay via the dimension 5 operator O_1 . The cross denotes a Wino mass insertion.

TABLE 1

	Q	u	d	L	e	H	$\bar{H}$
$\tilde{\alpha}_R$	0	-1	1	0	1	-1	1
$\tilde{\alpha}_A$	0	0	-1	-1	0	1	0
$\tilde{\alpha}_L$	0	0	0	-1	1	0	0
$\tilde{\alpha}_{PQ}$	0	-2	1	-1	2	-1	2

Table 1 : The charges of the three independent Z_N generators R_N, A_N, L_N and the “axionic” generator $R^2 A$.

TABLE 2

$Z_2 - GMP$	R_2	$A_2 R_2$				
$Z_2 - GLP$	L_2	$A_2 L_2$				
$Z_2 - GBP$	$R_2 L_2$	$A_2 R_2 L_2$				
$Z_3 - GMP$	R_3	A_3	$R_3 A_3$	$R_3 A_3 L_3$	$R_3 L_3^2$	$A_3 L_3^2$
$Z_3 - GLP$	L_3	$R_3 A_3^2 L_3^2$	$R_3 A_3^2 L_3$			
$Z_3 - GBP$	$R_3 L_3$	$A_3 L_3$	$R_3 A_3 L_3^2$			

Table 2 : The independent generalized Z_2 and Z_3 parities.

	$H\bar{H}$	$QQQL$	$uude$	$QQQH$	$QueH$	$LL\bar{H}\bar{H}$	$LH\bar{H}\bar{H}$	$\bar{H}\bar{H}e^*$	$\bar{H}^*H e$	QuL^*	ud^*c
GMP :											
R_2	✓	✓	✓			✓					
$A_2 R_2$				✓	✓	✓	✓		✓	✓	✓
R_3	✓	✓	✓								
$R_3 A_3$				✓	✓	✓			✓	✓	✓
$R_3 A_3 L_3$			✓	✓			✓				
$R_3 L_3^2$	✓										
A_3							✓	✓			
$A_3 L_3^2$		✓			✓	✓			✓	✓	✓
GLP :											
L_2	✓			✓		✓					
$A_2 L_2$		✓	✓		✓	✓	✓		✓	✓	✓
L_3	✓			✓							
$R_3 A_3^2 L_3$		✓					✓				
$R_3 A_3^2 L_3^2$			✓		✓	✓			✓	✓	✓
GBP :											
$R_2 L_2$	✓				✓	✓	✓	✓	✓	✓	✓
$A_2 R_2 L_2$		✓	✓	✓		✓		✓			
$R_3 L_3$	✓				✓	✓	✓	✓	✓	✓	✓
$A_3 L_3$											
$R_3 A_3 L_3^2$		✓		✓							

Table 3 : Dimension 2 and 5 operators allowed by each of the Z_2 and Z_3 generalized parities.

	$H\bar{H}$	$QQQL$	$uude$	$QQQH$	$QueH$	$LL\bar{H}\bar{H}$	$LH\bar{H}\bar{H}$	$\bar{H}\bar{H}e^*$	$\bar{H}^*He$	QuL^*	ud^*e
GRMP :											
P_2				✓	✓		✓	✓	✓	✓	✓
P_2A_2	✓	✓	✓								
GRLP :											
$P_2R_2L_2$		✓	✓		✓		✓	✓	✓	✓	✓
$P_2R_2A_2L_2$	✓			✓			✓	✓			
GRBP :											
P_2L_2		✓	✓	✓							
$P_2A_2L_2$	✓				✓			✓	✓	✓	✓

Table 4 – a : Dimension 2 and 5 operators allowed by the Z_2 R-symmetries.

	$H\bar{H}$	$QQQL$	$uude$	$QQQH$	$QueH$	$LL\bar{H}\bar{H}$	$LH\bar{H}\bar{H}$	$\bar{H}\bar{H}e^*$	$\bar{H}^*He$	QuL^*	ud^*e
GRMP :											
P_3				✓	✓				✓	✓	✓
P_3R_3											
P_3A_3	✓	✓	✓			✓					
$P_3L_3^2$			✓	✓							
$P_3R_3L_3$		✓			✓				✓	✓	✓
$P_3R_3A_3^2$					✓	✓			✓	✓	✓
$P_3R_3L_3^2$							✓	✓			
$P_3R_3A_3L_3^2$		✓									
$P_3R_3^2A_3^2L_3^2$	✓					✓					
$P_3R_3^2A_3^2$				✓	✓		✓	✓	✓	✓	✓
$P_3R_3^2A_3$	✓	✓	✓				✓	✓			
$P_3R_3^2A_3^2L_3$			✓	✓							
GRLP :											
$P_3R_3^2L_3^2$			✓		✓				✓	✓	✓
$P_3R_3^2L_3$		✓									
$P_3A_3^2L_3$			✓		✓		✓	✓	✓	✓	✓
$P_3A_3^2L_3^2$		✓				✓					
$P_3R_3A_3L_3$	✓			✓		✓					
$P_3R_3A_3L_3^2$	✓			✓			✓	✓			
GRBP :											
$P_3R_3^2A_3L_3$	✓				✓				✓	✓	✓
$P_3R_3^2A_3^2L_3^2$		✓		✓			✓	✓			
P_3L_3		✓		✓		✓		✓			
$P_3A_3L_3^2$	✓				✓				✓	✓	✓
$P_3R_3L_3^2$			✓			✓		✓			
$P_3R_3A_3^2L_3$							✓	✓			

Table 4 – b : Dimension 2 and 5 operators allowed by each of the Z_3 independent R-symmetries.