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LEFT-RIGHT SYMMETRY BREAKING SCALE AND
SUPERSYMMETRY IN UNIFIED THEORIES

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A B S T R A C T

We propose a class of supersymmetric grand unified models where the decoupling of parity and $SU(2)_R$ breaking scales are widely separated and compatible with a low lying mass for the right-handed gauge boson W_R . The intermediate symmetry $SU(4)_C \times SU(2)_L \times SU(2)_R$ and Higgs content are uniquely fixed if $m_{W_R} < 10^9$ GeV. The unification scale lies within an order of magnitude below the Planck mass.

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1. Introduction

The physics beyond the standard model has been the subject of a great deal of attention during the past few years, but we are still very far from a unique theory which is likely to explain how the world might look at energies above the one hundred GeV scale. In view of the well known natural stability of supersymmetry¹⁾ against radiative corrections leading to a solution of the gauge hierarchy problem, supersymmetric theories provide at present the most popular scenarios one can think of when going beyond the standard $SU(3)_C \times SU(2)_L \times U(1)_Y$. Despite the theoretical beauty of supersymmetry, the one to one correspondence between bosons and fermions it implies necessarily predicts the existence of supersymmetric partners of known particles. Their search has so far proved unsuccessful²⁾.

On the other hand, left right symmetric models³⁾ based on the group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ provide a direct and interesting extension of the Weinberg Salam theory. They give rise to a new energy scale M_R and new phenomenology associated with the breaking of the $SU(2)_R \times U(1)_{B-L}$ gauge symmetry down to $U(1)_Y$. The possibility that this scale lies within one or two orders of magnitude above the weak scale is particularly appealing, since right handed current signatures associated to W_R could be accessible in the next generation of experiments. Furthermore a low m_{W_R} theory can be embedded in a grand unified theory such as $SO(10)$ ⁴⁾. Two alternative scenarios have been proposed which are consistent with a low-lying right handed scale: the decoupling of parity and $SU(2)_R$ breaking scales⁵⁾ and mass split fermion representations⁶⁾. In the second case the detailed mechanisms for the mass splittings which were introduced⁷⁾ break however $SU(2)_R$ down to $U(1)_R$ so that only a Z_R (isosinglet) gauge boson is predicted and left right symmetry is lost. The supersymmetric extension of this situation has been investigated, the intermediate group being $SU(3)_C \times U(1)_{B-L} \times SU(2)_L \times U(1)_R$. The right handed scale M_R in this case related to the breaking $U(1)_{B-L} \times U(1)_R \rightarrow U(1)_Y$ could be as low as 1 TeV⁷⁾. Supersymmetry breaking effects were not however considered. The mass splitting mechanisms referred to above shift the $SU(2)_R$ breaking scale to energies $\sim 10^{14}$ GeV both in the $SU(3)_C \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$ ⁷⁾ and in the $SU(4)_C \times SU(2)_L \times SU(2)_R$ supersymmetric cases⁸⁾. Therefore, as far as a low-lying $SU(2)_R$ breaking scale is concerned (W_R gauge triplet), we are left with scenarios involving the decoupling of parity and $SU(2)_R$ breaking scales.

In this paper we propose a "natural" supersymmetric version of the decoupling of parity and $SU(2)_R$ breaking scale scenarios. No fine tunings are used except for the standard ones allowing for the different gauge breaking scales and, as long as supersymmetry remains unbroken, the two scales (parity and $SU(2)_R$) are kept far apart. A low-lying W_R appears to be possible only if the intermediate gauge group is $SU(4)_C \times SU(2)_L \times SU(2)_R$ which immediately implies a lower bound on its mass⁹⁾, as it is simultaneously associated to the breaking of $SU(4)_C$ symmetry. The unification scale comes within one order of magnitude of the Planck mass. The paper is divided as follows: in section 2 we expound our supersymmetric decoupling scenario and then present several variants of the model relating them to m_{W_R} (M_R) predictions. It will be seen that the one which is consistent with a low-lying W_R prediction uses two pairs of spinor representations (16, $\overline{16}$) and a 45 all undergoing mass splittings. In section 3 we analyze the direct supersymmetric extension of the decoupling scenario proposed by Chang, Mohapatra, Parida⁵⁾ and finally we look at the fully left-right symmetric models - the ones where $SU(2)_R$ and parity breaking occur simultaneously. It is found that these models when supersymmetrized, fail to produce low-lying W_R predictions consistent with $\sin^2 \theta_W$ ¹⁰⁾. Our renormalization group equations will "feel" the supersymmetry breaking through the mass parameter m_S describing the splitting within the supermultiplets which we take in the range $10^2 \text{ GeV} < m_S < 10^4 \text{ GeV}$. In section 4 we draw some conclusions and comments.

2. "Natural" Parity Decoupling Scenarios in Supersymmetry

The idea of decoupling D-parity and $SU(2)_R$ breaking scales was first proposed by Chang, Mohapatra and Parida⁵⁾ and relies on the introduction of a vacuum expectation value (VEV) of a field η which is odd under parity and invariant under $SU(2)_R$. Thus, in the context of $SO(10)$ unification

$$SO(10) \xrightarrow{M_X} G \times D \xrightarrow{\langle \eta \rangle = M_D} G \xrightarrow{M_R} SM \quad (1)$$

where D denotes the discrete parity symmetry, G is some subgroup of $SO(10)$ containing $SU(2)_R$ and SM stands for standard model. The D parity is defined by

$$D = \sigma_{67} \sigma_{23} \quad (2)$$

where σ_{ij} ($i, j = 1, \dots, 10$) denote the $SO(10)$ generators. So the field η must satisfy

$$D \eta D^{-1} = - \eta. \quad (3)$$

Examples of fields satisfying this condition are the $(15,1,1)$ and $(1,1,1)$ submultiplets* of 45 and 210 respectively of $SO(10)$. On the other hand, given the decompositions of 45 and 16,

$$\begin{aligned} 45 &= (15,1,1) + (1,3,1) + (1,1,3) + (6,2,2) \\ 16 &= (4,2,1) + (\bar{4},1,2) \end{aligned} \quad (4)$$

the D-parity transforms $(1,1,3) \longleftrightarrow (1,3,1)$ and e.g. $(4,2,1)$ of 16 into $(\bar{4},1,2)$ of 16 and vice-versa.

2.1 - $SU(3)_C \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$

Suppose that we start with a locally supersymmetric theory at the unification scale M_x , introduce in the superpotential the coupling

$$T(210) \phi(16) \bar{\phi}(\bar{16}) \quad (5)$$

and let the 210 acquire VEV's along the $(15,1,1)$ and $(1,1,1)$ directions thus breaking at M_x $SO(10)$ down to $SU(3)_C \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$. Since D-parity is a symmetry of $SO(10)$ the coupling (5) must be even under D. Furthermore, since the submultiplets $(15,1,1)$ and $(1,1,1)$ are even and odd respectively under D, the only way to expand the coupling (5) in $SU(4)_C \times SU(2)_L \times SU(2)_R$ representations is (in the vacuum):

$$\begin{aligned} <15,1,1>_T [(\bar{4},2,1)_{\bar{\phi}} (4,2,1)_{\phi} + (4,1,2)_{\bar{\phi}} (\bar{4},1,2)_{\phi}] + \\ &+ <1,1,1>_T [(\bar{4},2,1)_{\bar{\phi}} (4,2,1)_{\phi} - (4,1,2)_{\bar{\phi}} (\bar{4},1,2)_{\phi}]. \end{aligned} \quad (6)$$

The requirement that global supersymmetry is preserved at M_x scale implies the vanishing of all F terms in the vacuum, in particular

$$\begin{aligned} F_{(4,1,2)_{\bar{\phi}}} &= 0 = (<15,1,1>_T - <1,1,1>_T) (\bar{4},1,2)_{\phi} \\ F_{(\bar{4},1,2)_{\phi}} &= 0 = (<15,1,1>_T - <1,1,1>_T) (4,1,2)_{\bar{\phi}}. \end{aligned} \quad (7)$$

*The numbers in the brackets denote $SU(4)_C \times SU(2)_L \times SU(2)_R$ representations.

The components of ϕ , $\bar{\phi}$ in (7) must have non vanishing VEV in order to break the intermediate gauge group down to the standard model, so

$$\langle 15, 1, 1 \rangle_T = \langle 1, 1, 1 \rangle_T. \quad (8)$$

This leaves the $SU(2)_L$ singlets, $SU(2)_R$ doublets of ϕ , $\bar{\phi}$ massless at M_x scale and allows their mass to occur only at a much lower scale. On the other hand the $SU(2)_L$ doublets, $SU(2)_R$ singlets receive a mass contribution of the order of M_x from (6). This mass splitting, intrinsically connected to the VEV of an odd D-parity singlet, is obviously a non left-right symmetric effect although $SU(2)_R$ remains unbroken below M_x . It provides the appearance of an inequality

$$\alpha_{2L} > \alpha_{2R} \quad (9)$$

between the $SU(2)_L$ and $SU(2)_R$ coupling constants as will now become evident from our renormalization group analysis.

We start with the one loop renormalization group equations (RGE) for the $SU(3)_C$, $SU(2)_L$, $U(1)_Y$ couplings at energy μ

$$\begin{aligned} \frac{1}{\alpha_3} &= \frac{1}{\alpha_G} - \frac{b_3^-}{2\pi} \ln \frac{m_S}{\mu} - \frac{b_{3S}^-}{2\pi} \ln \frac{M_R}{m_S} - \frac{b_{3S}^+}{2\pi} \ln \frac{M_x}{M_R} \\ \frac{1}{\alpha_{2L}} &= \frac{1}{\alpha_G} - \frac{b_{2L}^-}{2\pi} \ln \frac{m_S}{\mu} - \frac{b_{2LS}^-}{2\pi} \ln \frac{M_R}{m_S} - \frac{b_{2LS}^+}{2\pi} \ln \frac{M_x}{M_R} = \frac{\sin^2 \theta_w}{\alpha} \\ \frac{1}{\alpha_1} &= \frac{5}{3} \frac{1}{\alpha_G} - \frac{5}{3} \frac{1}{2\pi} \left[b_1 \ln \frac{m_S}{\mu} + b_{1S} \ln \frac{M_R}{m_S} + \left(-\frac{2}{5} b_{1CS} + \frac{3}{5} b_{2RS} \right) \ln \frac{M_x}{M_R} \right] = \\ &= \frac{\cos^2 \theta_w}{\alpha}. \end{aligned} \quad (10)$$

Here the coefficients of the β functions are given for $SU(N)$ in the non-super-symmetric case by

$$b_N = \frac{11}{3} N - \frac{4}{3} N_G - \frac{2}{3} \sum_f T_f(R) - \frac{1}{6} \sum_s T_s(R) \quad (11)$$

where N_G is the number of generations (we assume $N_G=3$) and the third and fourth terms denote two component fermions and real Higgs scalar contribution respectively. $T(R)$ is the appropriate Casimir invariant for representation R . For the $U(1)_Y$ generator

$$b_1 = -\frac{4}{3} N_G - \frac{2}{5} \sum_f T_f(R) - \frac{1}{10} \sum_s T_s(R) \quad (12).$$

In the supersymmetric case we have

$$b_{NS} = 3N - 2 N_G - \sum_R n_R T(R) \quad (13)$$

and

$$b_{1S} = -2 N_G - \frac{3}{5} \sum_R n_R T(R) \quad (14)$$

where n_R is the number of R - representations. Superscripts $+$ and $-$ refer to representations above and below the intermediate scale M_R and the rest of the notation is standard. We take the low energy scale $\mu = 100$ GeV. The parameter m_S denotes the mass splitting within each supermultiplet due to supersymmetry breaking¹¹⁾

$$m_S = 0 \left(\frac{M_S^2}{M_P} \right) \quad (15)$$

where M_S is the primordial supersymmetry breaking scale. We demand m_S to run between its extreme values dictated by the lower mass bounds on supersymmetric partners of known particles and the gauge hierarchy preservation. Accordingly we take

$$10^2 \text{ GeV} < m_S < 10^4 \text{ GeV} \quad (16)$$

Below m_S no supersymmetry traces appear and above m_S every particle should have its own superpartner.

In addition to the $T(210)$, $\bar{\emptyset}(\bar{16})$, $\emptyset(16)$ representations referred to above we need an $H(10)$ representation containing the weak doublets (and colour triplets) undergoing the doublet/triplet splitting¹²⁾. The coupling $T(210) \cdot \Gamma(126) \cdot H(10)$ cannot however be used unless we introduce a different $T'(210)$ representation acquiring a VEV along the $(1,1,1)$ direction. This arises from the fact that $T(210)$ also has a VEV along $(15,1,1)$ giving through $\langle 15,1,1 \rangle_T (15,2,2)_\Gamma (1,2,2)_H$ an extralarge mass contribution to the weak doublets in $(1,2,2)$. A third 210 with zero VEV in all directions is still needed to provide an extraheavy mass for T, T' without giving rise to non zero F -terms of the superpotential in the vacuum. Thus we can have

$$\begin{aligned}
 W = & T \bar{\phi} \phi + T' \Gamma H + M_1 T T'' + M_2 T' T'' + M_3 \bar{\Gamma} \Gamma + A \bar{\phi}' \phi + A \bar{\phi} \phi' + M_4 \bar{\phi}' \phi' + \\
 & + \frac{\bar{\phi} \phi}{M_P} \phi_f \phi_f + H \phi_f \phi_f
 \end{aligned} \quad (17)$$

satisfying the vacuum condition $|W| \leq m_w M_P^2$.

The masses of the light components of $\phi, \bar{\phi}$ are signaled by the VEV's $\langle \phi \rangle, \langle \bar{\phi} \rangle$ and arise from the coupling to a singlet A and to the extraheavy $\phi', \bar{\phi}'$. The mass parameters in (17) are of the order of M_x or higher so there need not be any low energy ($< M_x$) traces of this complicated structure. The last two terms in (17) provide a Majorana mass for the right handed neutrino and the other fermion masses. Our collection of chiral superfields is then

$$\begin{aligned}
 & T, T', T'', \Gamma, \bar{\Gamma}, \phi, \bar{\phi} (\text{SU}(2)_L \text{ doublets}), \phi', \bar{\phi}', H (\text{colour triplets}) \text{ at scale } M_x \\
 & \phi, \bar{\phi} (\text{SU}(2)_R \text{ doublets}) \text{ at scale } M_R \\
 & \tilde{H} (\text{fermionic weak doublets}), \tilde{\phi}_f (\text{bosonic matter fields}) \text{ at scale } m_S \\
 & \psi (\text{ordinary fermions}), H (\text{weak doublets}) \text{ at scale } \mu.
 \end{aligned} \quad (18)$$

The corresponding b coefficients for eq. (10) are thus

$$\begin{aligned}
 b_1 = -\frac{41}{10}, \quad b_{2L}^- = \frac{19}{6}, \quad b_3^- = 7, \quad b_{1S} = -\frac{33}{5}, \quad b_{2LS}^- = -1, \quad b_{3S}^- = 3, \\
 b_{1CS} = -8, \quad b_{2LS}^+ = -1, \quad b_{2RS} = -5, \quad b_{3S} = 1
 \end{aligned} \quad (19)$$

leading to the following relation among the parameters M_x, M_R, m_S

$$M_x = C m_S^{-\frac{1}{6}} M_R^{-\frac{3}{7}} \quad \text{with} \quad C = \mu^{\frac{67}{42}} e^{-\frac{8\pi}{21} \left(\frac{1}{\alpha_3} - \frac{3}{8\alpha} \right)} \quad (20)$$

Note in (19) the inequality $b_{2LS}^+ > b_{2RS}$ implying $\alpha_{2L} > \alpha_{2R}$ and therefore a "pseudo" left right symmetry¹³⁾. Taking $\alpha = 0.00784$, $\alpha_3 = 0.101$ we obtain the values shown in Table I. As can be easily inferred from there, the lowest possible value for the right handed scale is $M_R \approx 10^9$ GeV which occurs for a rather high unification scale, $M_x \approx 5 \times 10^{18}$ GeV. Furthermore the insertion of an extra pair $16, \bar{16}$ undergoing the mechanism for parity decoupling described earlier (see eqs. (6), (7)) does not alter the mass relation (20): only $\sin^2 \theta_w$ is changed. A right handed scale which could be accessible to the next generation of experiments is out of question.

Let us now investigate the effect of the insertion of a 45 with the couplings

$$T(210) \Sigma(45) \Sigma(45) + S(54) \Sigma(45) \Sigma(45). \quad (21)$$

The S(54) is a 10x10 symmetric matrix which we assume to acquire a $VEV \sim M_x$ along its even parity (1,1,1) direction and T(210) is the one introduced earlier (exp. (5)). The relevant part of the vacuum decomposition of (21) is

$$\begin{aligned} & \langle 1,1,1 \rangle_T \left[(1,3,1)_{\Sigma} (1,3,1)_{\Sigma} - (1,1,3)_{\Sigma} (1,1,3)_{\Sigma} \right] + \\ & + \langle 1,1,1 \rangle_S \left[(1,3,1)_{\Sigma} (1,3,1)_{\Sigma} + (1,1,3)_{\Sigma} (1,1,3)_{\Sigma} \right]. \end{aligned} \quad (22)$$

As before the requirement of supersymmetry at the unification scale implies in the vacuum

$$F_{(1,1,3)_{\Sigma}} = 0 = (\langle 1,1,1 \rangle_S - \langle 1,1,1 \rangle_T) (1,1,3)_{\Sigma} \quad (23)$$

$$\text{and therefore } \langle 1,1,1 \rangle_S = \langle 1,1,1 \rangle_T \quad (24)$$

as the $SU(2)_R$ triplet of 45 may acquire a $VEV \sim M_R$. Consequently our superpotential will now be given by eq. (17) with in addition

$$T \Sigma \Sigma + S \Sigma \Sigma + \phi \bar{\phi}' \Sigma \quad (25)$$

The last term is inserted to account for the mass of $\Sigma(45)$ and ensures also $|W| \leq m_w M_P^2$. Recalling that we now have one extra $SU(2)_R$ triplet with mass $\sim M_R$, the b coefficients will be

$$\begin{aligned} b_1 &= -\frac{41}{10}, \quad b_{2L}^- = \frac{19}{6}, \quad b_3^- = 7, \quad b_{1S}^- = -\frac{33}{5}, \quad b_{2LS}^- = -1, \\ b_{3S}^- &= 3, \quad b_{1CS} = -8, \quad b_{2LS}^+ = -1, \quad b_{3S}^+ = 1, \quad b_{2RS} = -7 \end{aligned} \quad (26).$$

Again, we have an obvious asymmetry between left and right as seen from $b_{2LS}^+ > b_{2RS}$. The relation between M_x, M_R, m_S is now

$$M_x = C m_S^{-\frac{7}{48}} M_R^{-\frac{1}{4}} \text{ with } C = \mu^{\frac{67}{48}} e^{-\frac{\pi}{3} \left(\frac{1}{\alpha_3} - \frac{3}{8\alpha} \right)} \quad (27).$$

Numerical results are shown in Table II. It is seen that no M_R scale relevant for phenomenology is again possible in this case. The lowest possible value is $M_R \approx 10^{11}$ GeV occurring for $M_x \approx 10^{17}$ GeV. The trend shown in Table II shows a fact which can be immediately confirmed by direct calculation: for M_x closer to the Planck scale values of M_R in the range of a few TeV are possible, however at the price of an unrealistically low $\sin^2 \theta_w$ (0.18-0.19). This suggests us to consider the same Higgs content in the framework of the group $SU(4)_C$ instead of $SU(3)_C \times U(1)_{B-L}$.

2.2 - $SU(4)_C \times SU(2)_L \times SU(2)_R$

The VEV'S that break the unifying group $SO(10)$ must in this case obviously transform as singlets under $SU(4)_C$. The couplings to be used in the superpotential are now

$$T \Sigma \Sigma + S \Sigma \Sigma + T \bar{\emptyset} \bar{\emptyset} + S \bar{\emptyset} \bar{\emptyset} \quad (28)$$

where the $T(210)$ and $S(54)$ acquire VEV $\sim M_x$ along their $(1,1,1)$ components of the intermediate group. Imposing supersymmetry at M_x scale, by means of the mechanism described earlier the first two terms in (28) provide the mass splitting within Σ (45) (see eqs.(6)-(8)) and the last two generate the mass splitting within $\bar{\emptyset}$ (16), $\bar{\emptyset}$ (16). The one loop RGE are now

$$\begin{aligned} \frac{1}{\alpha_3} &= \frac{1}{\alpha_G} - \frac{b_3^-}{2\pi} \ln \frac{m_S}{\mu} - \frac{b_{3S}^-}{2\pi} \ln \frac{M_R}{m_S} - \frac{b_{4S}^-}{2\pi} \ln \frac{M_x}{M_R} \\ \frac{1}{\alpha_{2L}} &= \frac{1}{\alpha_G} - \frac{b_{2L}^-}{2\pi} \ln \frac{m_S}{\mu} - \frac{b_{2LS}^-}{2\pi} \ln \frac{M_R}{m_S} - \frac{b_{2LS}^+}{2\pi} \ln \frac{M_x}{M_R} = \frac{\sin^2 \theta_w}{\alpha} \\ \frac{1}{\alpha_1} &= \frac{5}{3} \frac{1}{\alpha_G} - \frac{5}{3} \frac{1}{2\pi} \left[b_1 \ln \frac{m_S}{\mu} + b_{1S} \ln \frac{M_R}{m_S} + \left(\frac{2}{5} b_{4S} + \frac{3}{5} b_{2RS} \right) \ln \frac{M_x}{M_R} \right] = \\ &= \frac{\cos^2 \theta_w}{\alpha} \end{aligned} \quad (29)$$

with the following b coefficients

$$\begin{aligned} b_1 &= -\frac{41}{10}, \quad b_{2L}^- = \frac{19}{6}, \quad b_3^- = 7, \quad b_{1S} = -\frac{33}{5}, \quad b_{2LS}^- = -1, \quad b_{3S}^- = 3, \\ &+ \\ b_{2LS}^+ &= -1, \quad b_{2RS} = -7, \quad b_{4S} = 4. \end{aligned} \quad (30)$$

Since we are now using the same Higgs content as in the previous case, the mass relation (27) is not altered. The weak mixing angle increases however too

much: for instance for $M_x = 10^{18}$ GeV, $M_R = 10^6$ GeV, $m_S = 5.2 \times 10^3$ GeV we would have $\sin^2 \theta_w = 0.296$ (compare with Table II). Furthermore it can be seen that no fitting is possible in this case for the parameters $M_x, m_S, \sin^2 \theta_w$ in the allowed regions at any value of M_R . If however we insert a second pair (16, $\overline{16}$) undergoing the same mass splitting as $\emptyset, \overline{\emptyset}$, then a low M_R is possible within the correct range for $m_S, \sin^2 \theta_w$ implying a rather high unification scale (see Table III). The b coefficients are the same as in (30) with the single exception

$$b_4 = 2. \quad (31)$$

We have therefore found a possibility for a low-lying W_R gauge boson in the context of supersymmetric unification. Its mass is connected to the scale of $SU(4)_c$, the quark lepton symmetry, and the unification scale is within one order of magnitude of the Planck mass in accordance with present ideas on superstring phenomenology¹⁴⁾. Upper and lower bounds on m_{W_R} existent in the literature¹⁵⁾ from left-right symmetric models are entirely compatible with the numbers obtained here.

3 - Supersymmetric scenario of Chang, Mohapatra, Parida and full left right symmetry

So far we have been concerned with model building which avoids fine-tunings except for the standard one connected to the existence of different gauge symmetry breaking scales. We now turn our attention to the supersymmetric extension of the Chang, Mohapatra, Parida⁵⁾ scenario, the first one to use the decoupling of parity and $SU(2)_R$ breaking scales. First we take the case where at scale M_x $SO(10)$ and parity are broken to $SU(3)_c \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$. To this end introduce in the superpotential (17) the coupling

$$\Sigma \overline{\Gamma} \Gamma \quad (32)$$

where the odd-parity (15,1,1) component of $\Sigma(45)$ acquires a VEV. The left handed triplets $\Delta_L \equiv (10,3,1), \overline{\Delta}_L \equiv (\overline{10},3,1)$ of $T(126), \overline{T}(\overline{126})$ receive therefore two positive mass contributions (from (32) and from the $M_3 \overline{\Gamma} \Gamma$ term). The right handed triplets $\Delta_R \equiv (10,1,3), \overline{\Delta}_R \equiv (\overline{10},1,3)$ receive a negative one from (32) and a positive one from the direct mass term. The mechanism described in section 2 is obviously available here as well but it would lead to a very

large number of light $SU(2)_R$ triplets, therefore upsetting the values of couplings and $\sin^2 \theta_w$. On the other hand, by appropriate fine tuning of the coefficient M_3 and since $SU(4)_c$ is also broken, only a pair of the right handed chiral triplets may be made massless at scale M_x . In the absence of any other criterion, simplicity will lead us to assume this triplet pair to lie close to the other scale of the model, M_R . Thus we have $\alpha_{2L} > \alpha_{2R}$ as before. The intermediate symmetry is broken by a pair $16, \overline{16}$ with no mass splitting. A suitable superpotential is

$$W = TTH + \Sigma \bar{T}T' + M_1 TT' + M_3 \bar{T}T' + A \bar{\theta} \theta + A \bar{\theta} \theta' + M_4 \bar{\theta} \theta' + \frac{\bar{\theta} \bar{\theta}}{M_P} \theta_f \theta_f + H \theta_f \theta_f \quad (33)$$

with notation as before (17). We find

$$b_1 = -\frac{41}{10}, b_{2L}^- = \frac{19}{6}, b_3^- = 7, b_{1S} = -\frac{33}{5}, b_{2LS}^- = -1, b_{3S}^- = 3, \\ b_{1CS} = -19, b_{2LS}^+ = -5, b_{2RS} = -9, b_{3S}^+ = -1 \quad (34)$$

leading to

$$M_x = C m_S - \frac{7}{72} M_R \frac{1}{6} \quad \text{with } C = \mu^{\frac{67}{72}} e^{-\frac{2\pi}{9}(\frac{1}{\alpha_3} - \frac{3}{8\alpha})} \quad (35)$$

Numerical results are displayed in Table IV for the only allowed values of m_S , $\sin^2 \theta_w$. A single region exists for the right handed scale, $10^{11} < M_R < 4 \times 10^{11}$ GeV. In unbroken $SU(4)$ (for the same Higgs spectrum) a region $30 \text{ TeV} < M_R < 60 \text{ TeV}$ exists but supersymmetry would essentially lose its relevance ($m_S \approx 10^{17}$ GeV) due to the large number of light states. We find no motivation to combine this class of models using extra fine tunings for Higgs masses with the "natural" models of section 2.

To conclude this section we briefly discuss the models with no decoupling of parity and $SU(2)_R$ breaking scales, in other words the fully left right symmetric ones. At the intermediate scale we simply now have the spinor pair $16, \overline{16}$. The rest of the Higgs spectrum is unchanged. We find

$$M_x = C m_S - \frac{1}{6} M_R - \frac{3}{7} \quad \text{with } C = \mu^{\frac{67}{42}} e^{-\frac{8\pi}{21}(\frac{1}{\alpha_3} - \frac{3}{8\alpha})} \quad (36)$$

and b coefficients

$$\begin{aligned}
 b_1^- &= -\frac{41}{10}, \quad b_{2L}^- = \frac{19}{6}, \quad b_3^- = 7, \quad b_{1S}^- = -\frac{33}{5}, \quad b_{2LS}^- = -1, \\
 b_{3S}^- &= 3, \quad b_{1CS}^- = -10, \quad b_{2LS}^+ = b_{2RS}^+ = -5, \quad b_{3S}^+ = 1
 \end{aligned}
 \tag{37}$$

Numerical results are shown in Table V. It is seen that the lowest possible value for the right handed scale is $M_R \approx 7 \times 10^9$ GeV for an unification scale $M_x \approx 10^{18}$ GeV. Again the $SU(4)_c$ case is totally uninteresting.

4 - Summary and Concluding Remarks

We have investigated the prospects for a low-lying charged right handed current scale in supersymmetric models. In a like manner as for the non-supersymmetric case, only the models involving the decoupling of parity and $SU(2)_R$ breaking scales seem to be consistent with a light W_R . We proposed a natural way to implement the decoupling scenarios in supersymmetry, through the vanishing of F-terms at the unification scale, which generates mass split representations. These mass splittings should be protected by supersymmetry with no need for extra fine tunings for the Higgs masses. In this way we are led to a unique model allowing for a low mass W_R where $SU(2)_R$ and $SU(4)_c$ breakings occur simultaneously. The Higgs content at intermediate scale M_R consists of sixteen $SU(2)_R$ chiral doublets and one $SU(2)_R$ chiral triplet. The supersymmetric extension of Chang, Mohapatra, Parida scenario and the fully left right symmetric scenario have also been analysed. They are uninteresting from the point of view of low mass W_R predictions. On the other hand the theoretical bounds¹⁵⁾ on m_{W_R} (upper and lower) together with a lower bound on the $SU(4)_c$ scale suggest that an m_{W_R} in the 10 TeV range is totally consistent with our $SU(4)_c$ model. The unification scale would then come in the range $(2-6) \times 10^{18}$ GeV, i.e. within an order of magnitude of M_P in agreement with current ideas on superstring phenomenology¹⁴⁾, implying unobservability of proton decay. But on the basis of these considerations, expecting observability of $N-\bar{N}$ oscillations¹⁶⁾, right handed currents and supersymmetric particles within the forthcoming generation of experiments would be far too optimistic.

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$M_x = 10^{19}$ GeV			$M_x = 10^{18}$ GeV		
M_R	m_S	$\sin^2 \theta_w$	M_R	m_S	$\sin^2 \theta_w$
3×10^7	1.6×10^4	0.206	8×10^9	9.27×10^3	0.214
10^8	726	0.210	10^{10}	5.22×10^3	0.215
3×10^8	43	0.213	10^{11}	14	0.222

Table I - The values of M_R (right handed scale), m_S (supersymmetry mass splitting) and $\sin^2 \theta_w$ for a given unification scale M_x in the $SU(3)_C \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$ scenario with "natural" parity decoupling and eight right handed chiral doublets at M_R scale. Units are in GeV.

$M_x = 10^{18}$ GeV			$M_x = 10^{17}$ GeV		
M_R	m_S	$\sin^2 \theta_w$	M_R	m_S	$\sin^2 \theta_w$
5×10^5	1.7×10^4	0.190	10^{10}	5.22×10^3	0.209
10^6	5.23×10^3	0.192	10^{11}	100	0.216
10^7	100	0.198	3×10^{11}	15	0.219

Table II - The values of M_R , m_S and $\sin^2 \theta_w$ in the $SU(3)_C \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$ scenario with "natural" parity decoupling, eight right handed $SU(2)_R$ chiral doublets and an $SU(2)_R$ chiral triplet at M_R scale. Units are in GeV.

$M_x = 5 \times 10^{18} \text{ GeV}$			$M_x = 3 \times 10^{18} \text{ GeV}$		
M_x	m_S	$\sin^2 \theta_w$	M_R	m_S	$\sin^2 \theta_w$
8×10^3	331	0.227	10^4	7.5×10^3	0.225
10^4	226	0.227	3×10^4	1.14×10^3	0.226
2×10^4	69	0.228	5×10^4	475	0.227

Table III - The values of M_R , m_S and $\sin^2 \theta_w$ in the $SU(4)_c \times SU(2)_L \times SU(2)_R$ scenario with "natural" parity decoupling, sixteen $SU(2)_R$ chiral doublets and $SU(2)_R$ chiral triplet at M_R scale. Units are in GeV.

$M_x = 10^{15} \text{ GeV}$		
M_R	m_S	$\sin^2 \theta_w$
10^{11}	100	0.227
3×10^{11}	663	0.218
4×10^{11}	1.1×10^3	0.215

Table IV - The values of M_R , m_S for the only self consistent range of parameters in the supersymmetric parity decoupling scenario of Chang, Mohapatra, Parida for the group $SU(3)_c \times U(1)_{B-L} \times SU(2)_L \times SU(2)_R$. At M_R scale there are eight $SU(2)_R$ chiral doublets and two $SU(2)_R$ chiral triplets.

$M_X = 10^{16} \text{ GeV}$			$M_X = 10^{18} \text{ GeV}$		
M_R	m_S	$\sin^2 \theta_w$	M_R	m_S	$\sin^2 \theta_w$
5×10^{14}	4.3×10^3	0.220	7×10^9	1.3×10^4	0.237
10^{15}	726	0.229	10^{10}	5.22×10^3	0.242

Table V - The values of M_R , m_S and $\sin^2 \theta_w$ in the $SU(4)_C \times SU(2)_L \times SU(2)_R$ fully left right symmetric scenario (no parity decoupling). At M_R scale there are eight $SU(2)_R$ chiral doublets.

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