SO(10) Grand Unification with a Low-Energy $SU(2)_R$-Symmetry-Breaking Scale M_R

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Contrary to commonly held belief, we show that one can obtain a low value for M_R, the $SU(2)_R$ breaking scale, in grand unification theories based on SO(10). This possibility emerges in the supersymmetric version of SO(10) with a judicious choice of Higgs content. The unification scale is found to be consistent with the constraint from proton decay. This result is first explicitly demonstrated using the one-loop renormalization group equations, and then a full two-loop analysis is carried out.

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It is commonly believed [1,2] that the group $SU(2)_R$ has to be broken at a large energy scale $M_R \sim 10^{10}$ GeV if it is to emerge from a grand unified symmetry such as SO(10). This is also assumed to be true for the supersymmetric (SUSY) version of SO(10) [1,2]. Consequently, additional gauge bosons that could possibly be produced at supercollider energies are thought to originate only from additional U(1) factors which lead to Z' bosons [3]. The phenomenology of new charged W' bosons at supercolliders is therefore less frequently investigated [4]. We will show in this Letter that although the above result is true for the simplest Higgs structure, if this sector is suitably enlarged, the scale for the right-handed gauge bosons, M_R, could be made arbitrarily low. We will consider only the supersymmetric version in detail, with some brief remarks on the nonsupersymmetric case given at the end of our discussion.

We investigate the break chain

$$\mathrm{SO(10)(SUSY)} \overset{M_U}{\rightarrow} 2_L 2_R 1_{B-L} 3_C(\mathrm{SUSY}) \overset{M_R}{\rightarrow} 2_L 1_Y 3_C(\mathrm{SUSY}) \overset{M_Z}{\rightarrow} 3_C 1_Q \ , \qquad (1)$$

where, as an example of our notation, 2_L represents $SU(2)_L$. Here, we have set the "effective" supersymmetry breaking scale to be M_Z, and will comment on this later.

In Ref. [5], it was shown that if both $2_L 2_R 1_{B-L} 3_C$ singlets of the **210** representation, together with the $2_L 2_R 1_{B-L} 3_C$ singlet of the **45** representation, acquire vacuum expectation values (VEVs) then this is sufficient to break supersymmetric SO(10) down to supersymmetric $2_L 2_R 1_{B-L} 3_C$ without D parity. The breaking at M_R can be performed either by the Higgs fields in the $\mathbf{126}\oplus\overline{\mathbf{126}}$ representation or in the $\mathbf{16} \oplus \overline{\mathbf{16}}$ representation and we consider both these possibilities in our discussion below. We further assume that ordinary electroweak breaking at the Z scale is achieved as usual by a complex **10** representation. For the purpose of generating fermion masses, we assume that the entire bi-doublet of the **10** representation has a mass at the scale of M_Z. [We remind the reader that a bi-doublet corresponds to the (2,2,0,1) representation of $2_L 2_R 1_{B-L} 3_C$.] Also, we assume that the $SU(2)_R$ triplets of the **126** and $\overline{\mathbf{126}}$ representations and the $SU(2)_R$ doublets of **16** and $\overline{\mathbf{16}}$ representations have masses at the scale M_R. All other Higgs multiplets are given masses of order M_U as follows from the survival hypothesis. We make the important observation that in this symmetry breaking pattern pseudo Goldstone bosons do not appear [5].

First let us examine the one-loop equations:

$$\alpha_{1Y}^{-1}(M_Z) = \alpha_U^{-1}(M_U) + \frac{b_{1Y}}{2\pi} R + \frac{1}{2\pi}\left(\frac{3b_{2R}}{5} + \frac{2b_{B-L}}{5}\right) \times (U - R)$$

$$\alpha_{2L}^{-1}(M_Z) = \alpha_U^{-1}(M_U) + \frac{b_{2L}}{2\pi} U \ , \qquad (2)$$

$$\alpha_{3C}^{-1}(M_Z) = \alpha_U^{-1}(M_U) + \frac{b_{3C}}{2\pi} U \ ,$$

where

$$R = \ln\frac{M_R}{M_Z}, \quad U = \ln\frac{M_U}{M_Z} \ . \qquad (3)$$

The b_i's are the one-loop beta functions, which for the supersymmetric case are given by

$$b_N^{\mathrm{SUSY}} = 2n_g - 3N + T(S_N) \ , \qquad (4)$$

for n_g generations, the gauge group $SU(N)$, and the complex Higgs fields contribution which is parametrized by $T(S_N)$. For U(1) gauge groups, $N = 0$ in the above equation and the gauge couplings are normalized as usual. Explicitly we find the Higgs boson contributions to be given by

$$T_{1Y} = \tfrac{3}{5} n_{10}, \quad T_{2L} = n_{10}, \quad T_{3C} = 0 \ ,$$
$$T_{2R} = n_{10} + n_{16} + 4n_{126}, \quad T_{1X} = \tfrac{3}{2} n_{16} + 9n_{126} \ , \qquad (5)$$

where the subscripts on the T's refer to the relevant gauge group. In the above, n_{10} is the number of complex **10** Higgs bi-doublets at the scale M_Z, and the n_{16} and n_{126} are the number of $\mathbf{16} \oplus \overline{\mathbf{16}}$ and $\mathbf{126} \oplus \overline{\mathbf{126}}$ Higgs boson pairs, respectively, which are used to break the intermediate gauge symmetry. Using Eqs. (2) and (5) together with the definitions

$$\alpha_{1Y}^{-1}(M_Z) = \frac{3}{5}\frac{1-\tilde{x}}{\alpha(M_Z)}, \quad \alpha_{2L}^{-1}(M_Z) = \frac{\tilde{x}}{\alpha(M_Z)}, \tag{6}$$

gives the relations

$$\begin{aligned} \frac{2\pi}{\alpha(M_Z)}\left(1 - \frac{8\alpha(M_Z)}{3\alpha_{3C}(M_Z)}\right) &= (C_1 - C_2)U + C_2 R, \\ \frac{2\pi}{\alpha(M_Z)}\left(1 - \frac{8}{3}\tilde{x}\right) &= \left(\frac{5}{3}C_3 - C_2\right)U + C_2 R, \end{aligned} \tag{7}$$

where the abbreviation $\tilde{x} \equiv \sin^2\theta_W(\overline{\text{MS}})$ ($\overline{\text{MS}}$ denotes the modified minimal-subtraction scheme) is used, and the C_i are given by

$$\begin{aligned} C_1 &\equiv b_{2L} + \tfrac{5}{3}b_{1Y} - \tfrac{8}{3}b_{3C} = 18 + 2n_{10}, \\ C_2 &\equiv \tfrac{5}{3}b_{1Y} - b_{2R} - \tfrac{2}{3}b_{1X} = 6 - 2n_{16} - 10n_{126}, \\ C_3 &\equiv b_{1Y} - b_{2L} = 6 - \tfrac{2}{5}n_{10}. \end{aligned} \tag{8}$$

We make the observation that if $C_2 = 0$ then the scale M_R is completely undetermined at the one-loop level. This gives us hope that when $C_2 = 0$, a solution with a low-energy M_R will exist. We can have $C_2 = 0$ only when $n_{16} = 3$ and $n_{126} = 0$. We then need only to require that the two equations in (7) give agreeing values of M_U within the level of accuracy of the one-loop approximation. The latest values [6] of the input parameters that we use in this analysis are

$$\begin{aligned} \alpha^{-1}(M_Z) &= 127.9 \pm 0.1, \quad \alpha_{3C}(M_Z) = 0.118 \pm 0.007, \\ \tilde{x}(M_Z) &= 0.2326 \pm 0.0011, \quad M_Z = 91.187 \pm 0.007\ \text{GeV}. \end{aligned} \tag{9}$$

Consistency of the two equations in (7) with $n_{16} = 3$, $n_{126} = 0$, and $n_{10} = 1$, and taking the central values of $\alpha^{-1}(M_Z)$ and $\tilde{x}(M_Z)$ from above implies $\alpha_{3C}(M_Z) =$ 0.112, which is within the experimentally allowed range given above. Of course we will have to perform a full two-loop analysis to ensure a solution exists with a low-energy M_R with $n_{16} = 3$, $n_{126} = 0$, and $n_{10} = 1$. Note that for this particular choice of Higgs representations the combination $\frac{3}{5}\alpha_{2R}^{-1} + \frac{2}{5}\alpha_{1X}^{-1}$ runs identically at the one-loop level as α_{1Y}^{-1} in the minimal supersymmetric standard model (MSSM). Since supersymmetric SU(5) grand unification is consistent with $M_S \simeq M_Z$ [1,2,7], one would naively expect that a two-loop analysis will show the SO(10) scenario to be equally consistent and we will now show this explicitly.

We numerically integrate the two-loop equations

$$\mu\frac{\partial\alpha_i(\mu)}{\partial\mu} = \frac{1}{2\pi}\left(b_i + \frac{b_{ij}}{4\pi}\alpha_j(\mu)\right)\alpha_i^2(\mu) \tag{10}$$

assuming the breaking pattern as given in Eq. (1) (with $n_{10} = 1$, $n_{16} = 3$, and $n_{126} = 0$) and we use the approximation that all superparticles have mass $M_S = M_Z$. Actually, the "average" sparticle masses could be somewhat different than this effective value [1,2,7]. We will also assume that the top quark and right-handed neutrino have masses $\simeq M_Z$. We use the appropriate two-loop matching conditions [8] at M_U that follow from dimensional reduction:

$$\alpha_U^{-1}(M_U) - \frac{C_U}{12\pi} = \alpha_i^{-1}(M_U) - \frac{C_i}{12\pi}, \tag{11}$$

where i represents the intermediate gauge groups 2_L, 2_R, 1_{B-L} or 3_C and C_G is the quadratic Casimir invariant for group G. Similarly at M_R we have

$$\alpha_{1Y}^{-1}(M_R) = \frac{3}{5}\left(\alpha_{2R}^{-1}(M_R) - \frac{C_2}{12\pi}\right) + \frac{2}{5}\alpha_{1B-L}^{-1}(M_R). \tag{12}$$

$C_G = N$ for SU(N) and $C_G = 0$ for U(1). In Eq. (10), we assume the MSSM below the scale M_R so that, with $i = 1_Y$, 2_L, 3_C, respectively, we use the following two-loop beta functions [9]:

$$b_{ij}^{\text{MSSM}} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & -24 & 0 \\ 0 & 0 & -54 \end{pmatrix} + n_g \begin{pmatrix} \frac{38}{15} & \frac{6}{5} & \frac{88}{15} \\ \frac{2}{5} & 14 & 8 \\ \frac{11}{15} & 3 & \frac{68}{3} \end{pmatrix} + n_{10} \begin{pmatrix} \frac{9}{25} & \frac{9}{5} & 0 \\ \frac{3}{5} & 7 & 0 \\ 0 & 0 & 0 \end{pmatrix}. \tag{13}$$

For the intermediate symmetry SUSY $2_L 2_R 1_{B-L} 3_C$, we derive from the generic two-loop expression in Ref. [9]

$$b_{ij}^{\text{int}} = \begin{pmatrix} -24 & 0 & 0 & 0 \\ 0 & -24 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & -54 \end{pmatrix} + n_g \begin{pmatrix} 14 & 0 & 1 & 8 \\ 0 & 14 & 1 & 8 \\ 3 & 3 & \frac{7}{3} & \frac{8}{3} \\ 3 & 3 & \frac{1}{3} & \frac{68}{3} \end{pmatrix} + n_{10} \begin{pmatrix} 7 & 3 & 0 & 0 \\ 3 & 7 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix} + n_{16} \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 7 & \frac{3}{2} & 0 \\ 0 & \frac{9}{2} & \frac{9}{4} & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}, \tag{14}$$

where $i, j = 2_L, 2_R, 1_{B-L}, 3_C$, respectively. In evaluating the above two equations for b_{ij}, we of course assume $n_g = 3$, $n_{10} = 1$, and $n_{16} = 3$. We treat M_R as a free parameter. From the one-loop calculation, we expect that the two-loop analysis will yield solutions for arbitrary values of M_R between M_Z and M_U. We have explored the possibility that M_R can take on a wide range of potential values and find the expectation above to be fulfilled. As an example of this and in particular to show that M_R can be low, we display in Fig. 1 the case $M_R = 1$ TeV. We find that $\alpha_U^{-1}(M_U) = 23.4 \pm 0.5$ and $M_U = 10^{16.2\pm0.4}$ GeV, which is sufficiently large to be consistent with the nonobservation of proton decay. We also note that at the scale M_R, $\alpha_{2L}/\alpha_{2R} \simeq 1.5$, thus implying the bound $M_{W_R} > 380$ GeV from muon decay [10], and 480 GeV from direct collider searches [11,12]. We have investigated the influence of heavy top quark Yukawa couplings on the two-loop evolution [7], and find the effect to be negligible. *Thus, at least for this choice of symmetry breaking, M_R can indeed be sufficiently low as to be of consequence for existing and future colliders.* The most important as well as immediate impact of low M_R is the predicted existence of new gauge bosons, Z', W' which can be copiously produced and easily detected at both the Superconducting Super Collider (SSC) and the CERN Large Hadron Collider (LHC) for masses as large as several TeV [3,4]. For example, after acceptance cuts, a 3 TeV Z' in this model can lead to more than 100 dilepton pair events in each flavor channel at the SSC. A somewhat larger number of W' events would be expected for the same mass.

We have now described the conditions under which M_R may be associated with a low scale in supersymmetric SO(10) grand unification. From the previous discussion, we can see that we need $M_S \simeq M_Z$ for our scenario to be realized as in the case of SUSY SU(5). What about nonsupersymmetric SO(10) grand unification? In this case, the equations analogous to (7) and (8) imply that for M_R to drop out of the one-loop equations we require that $5n_{126} + n_{16} = 22$, where n_{126} and n_{16} refer to the number of **126** and **16** representations, respectively, used to break the intermediate gauge symmetry $2_L 2_R 1_{B-L} 3_C$ to the standard model. Demanding consistency of the two one-loop equations analogous to Eq. (8) and using the values of the low-energy parameters as given above further implies $n_{10} = 4$, where n_{10} is the number of scalar bi-doublets at the scale M_Z. [The entire Higgs multiplet (2,2,0,1) within the **10** representation must have mass less than M_R, otherwise the one-loop equations will still depend on M_R.] This illustrates the result that in order to achieve grand unification in this nonsupersymmetric SO(10) case, as in the conventional SU(5) model [7], we would have to employ many Higgs doublets. In the SO(10) case, we then obtain unification with $M_U \simeq 10^{13.6}$ GeV, a value which is clearly inconsistent with limits on the proton lifetime.

We note that in our model small neutrino masses also arise from a variation of the usual seesaw mechanism. Since we do not use the **126** representation which permits the conventional seesaw mechanism with a low $B-L$ breaking scale and the two-loop radiative Witten mechanism [13] will not work with a low $B-L$ in a supersymmetric context [14], we must adopt one of the strategies used in superstring inspired models [4]. Mohapatra has demonstrated two scenarios that can be employed in our situation. One possibility involves gaugino mixing [15] while the second introduces grand unified theory singlet fermions [16]. The latter is preferred in three-generation SUSY models, and has been used successfully in the SO(10) context in the literature [17]. In our case with the **16** representation having a low mass scale it provides a natural seesaw mechanism. Furthermore, our renormalization group analysis is unaltered in this case.

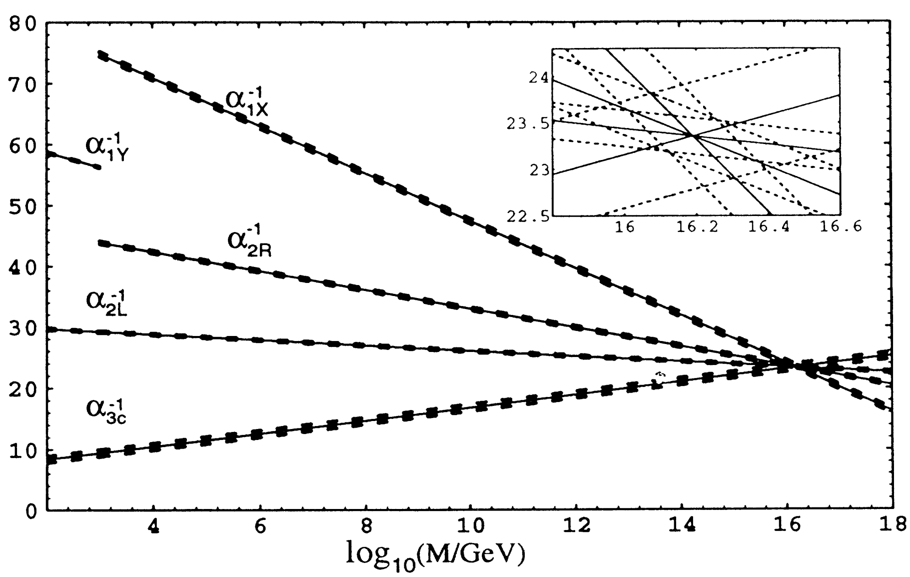

FIG. 1. Evolution of coupling constants for the symmetry breaking chain given in Eq. (1). We have used $n_{10} = 1$, $n_{16} = 3$, $n_{126} = 0$, and $M_R = 1$ TeV, where these quantities are defined in the text. The error bars we show arise from uncertainties in the low-energy parameters in Eq. (9). In the figure, α_i^{-1} is calculated via the dimensional reduction scheme at two-loop order.

As is well known, once the seesaw mechanism generates hierarchal neutrino masses, it is quite easy to accommodate existing bounds on neutrino properties from both oscillation experiments and astrophysics while simultaneously predicting a mass for the τ-neutrino in a cosmologically interesting range [18].

In summary, we have shown for the first time that a low $SU(2)_R$ breaking scale is compatible with SO(10) grand unification. The symmetry breaking at the scale M_R is accomplished by three generations of Higgs bosons in the **16** representation, not unlike the three families of quarks and leptons. Further consequences of this symmetry breaking scenario for fermion masses will be the subject of a future investigation.

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