Supersymmetric SO(10) Seesaw Mechanism with Low *B-L* Scale

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We propose a new seesaw mechanism for neutrino masses within a class of supersymmetric SO(10) models with broken D parity. It is shown that in such scenarios the B-L scale can be as low as TeV without generating inconsistencies with gauge coupling unification nor with the required magnitude of the light neutrino masses. This leads to a possibly light new neutral gauge boson as well as relatively light quasi-Dirac heavy leptons. These particles could be at the TeV scale and mediate lepton flavor and CP violating processes at appreciable levels.

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The origin of neutrino masses is the most well-kept secret of modern elementary particle physics. The basic dimension-five operator which leads to neutrino masses [1] can arise from physics at vastly different scales. One popular alternative is the seesaw mechanism, in which case the small neutrino masses are induced by the exchange of superheavy neutral fermions [2–5] or superheavy scalars, or both [6–8]. The light neutrino masses are given as

$$M_\nu \simeq -v^2 Y M^{-1} Y^T. \tag{1}$$

An alternative inverse seesaw scheme has been suggested [9] for theories which lack the representation required to implement the canonical seesaw, as happens in a class of string inspired models.

In addition to the normal neutrinos ν, such an inverse seesaw mechanism employs two sequential $SU(3) \otimes SU(2) \otimes U(1)$ singlets, ν^c and S (these are all left-handed two-component spinors [6]). The effective neutrino mass matrix has the following form:

$$M_\nu = \begin{pmatrix} 0 & Yv & 0 \\ Y^T v & 0 & M \\ 0 & M & \mu \end{pmatrix} \tag{2}$$

in the basis ν_L, ν_L^c, S_L. Here Y is the Yukawa matrix parametrizing the $Y^{ij}\nu_L^{iT} C^{-1} \nu_L^{jc} + \text{H.c.}$ interactions, while M and μ are $SU(3) \otimes SU(2) \otimes U(1)$ invariant mass entries. When $\mu \to 0$, a global lepton-number symmetry is exactly conserved and all three light neutrinos are strictly massless. Yet it has been shown [10,11] that lepton flavor and CP can be violated at appreciable levels even in the absence of supersymmetry, provided the scale M is sufficiently low. When $\mu \neq 0$, the mass matrix for the light eigenstates is given by

$$M_\nu \simeq -v^2 (Y M^{T-1}) \mu (M^{-1} Y^T). \tag{3}$$

One sees that neutrinos can be made very light, as required by oscillation data [12], even if M is very low, far below the unified scale $M_G (M \ll M_G)$, provided μ is very small, $\mu \ll M$. This scheme has a very rich and interesting phenomenology, since no new scales need to be added to generate the small neutrino masses, instead a small parameter μ is added. Note that in such a $SU(3) \otimes SU(2) \otimes U(1)$ inverse seesaw the smallness of μ is natural, in the 't Hooft sense [13], as the symmetry enhances when $\mu \to 0$. However, there is no dynamical understanding of this smallness.

Here we provide an alternative inverse seesaw realization consistent with a realistic unified SO(10) model. Such embedding brings several issues: (i) The $\nu^c S$ entry M [generated by the vacuum expectation value (VEV) of a Higgs multiplet χ_R with $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ quantum numbers $(1, 1, 2, -1)$] breaks the B-L symmetry, now gauged. The corresponding scale $\langle \chi_R \rangle$ must be compatible with gauge coupling unification. Together with the requirement of low-energy supersymmetry to stabilize the hierarchies, this places rather strong constraints on how we must fill the "desert" of particles below M_G. (ii) The need to justify the magnitude of the singlet μSS mass. (iii) It implies a nonzero νS entry in Eq. (2), proportional to the VEV of the L-R partner of χ_R, namely, $\chi_L \equiv (1, 2, 1, +1)$. Let us now discuss one by one these three points and show that there indeed exists a supersymmetric SO(10) model that addresses all these conditions in a satisfactory way and offers a new way to understand the smallness of neutrino masses.

First, note that there are several mechanisms that could be used to get rid of the SS term in Eq. (2). For example, we can treat the SO(10) embedding into E_6 where the fermionic singlet could be a member of a 27-dimensional irreducible representation with the familiar $\mathrm{SO}(10) \otimes \mathrm{U}(1)_X$ decomposition,

$$27_F = 1_F^4 \oplus 16_F^1 \oplus 10_F^{-2}. \tag{4}$$

If at the E_6 scale there is no $351'$ Higgs representation, the U(1)-charge of the $1_F 1_F$ matter bilinear is so large that it is very hard to saturate it. Thus, as long as the corresponding U(1) is unbroken we have $\mu = 0$. Even if we break the U(1) symmetry at some lower scale, it could be rather complicated to generate an effective *SS* entry, which brings further suppression, even at the level of effective operators. From now on we will neglect μ.

Now consider the νS term. A typical SO(10) superpotential contains the following terms:

$$W \ni M_{16} 16_H \overline{16}_H + \rho 16_H 16_H 10_H + \text{H.c.}$$

The fact that in "standard" supersymmetric SO(10) models there is a small induced VEV generated for the neutral component of $\chi_L = (1, 2, +1)_{SM} \in (1, 2, 1, +1)_{LR}$ of $\overline{16}_H$ once the *B-L* symmetry is broken can be seen from the structure of the F terms. For example, $F^\dagger_{(1,2,1,\pm 1)}$ is proportional to

$$M_{16}(1, 2, 1, \mp 1)_{16} + \rho(1, 1, 2, \mp 1)_{16}(1, 2, 2, 0)_{10} + \cdots.$$

After giving a nonzero VEV to the $(1, 1, 2, \mp 1)$ field (that subsequently breaks *B-L*) and the traditional doublet pair in 10_H [to break the standard model (SM)] the requirements to be in a supersymmetric vacuum lead to

$$\langle \chi_L \rangle \equiv v_L \simeq \langle (1, 2, 1, \mp 1)_{16} \rangle \simeq \rho \frac{v_R v}{M_{16}}. \tag{5}$$

Therefore, there is a new contribution to (2) coming from a term of the type $F^{ikl} \nu_L^{iT} C^{-1} S_L^k \chi_L^l$ + H.c. so that the neutrino mass matrix reads

$$M_\nu = \begin{pmatrix} 0 & Yv & Fv_L \\ Y^T v & 0 & \tilde{F} v_R \\ F^T v_L & \tilde{F}^T v_R & 0 \end{pmatrix} \tag{6}$$

instead of Eq. (2). Here $\tilde{F}$ is an independent combination of the VEVs of the $\chi_L^l s$, namely, $\tilde{F}^{ij} v_L = \sum_k F^{ijk} \langle \chi_L \rangle^k$, while $F^{ij} v_R = \sum_k F^{ijk} \langle \chi_R \rangle^k$. By inserting v_L from (5) into Eq. (6) one sees that the v_R scale drops out completely from the previous formula, leading to

$$M_\nu \simeq \frac{v^2}{M_G} \rho [Y(F\tilde{F}^{-1})^T + (F\tilde{F}^{-1}) Y^T] \tag{7}$$

so that the neutrino mass is suppressed by M_G *irrespective of how low the B-L breaking scale is*. This is a key feature of our mechanism, illustrated in Fig. 1. In contrast to both the canonical seesaw in Eq. (1) and the inverse seesaw in Eq. (3), this new seesaw is linear in the Dirac Yukawa couplings Y. Note also that, for given M_G and Y, the scale of neutrino masses can be adjusted by choosing appropriately the value of the cubic scalar sector coupling constant ρ, as well as the $F\tilde{F}^{-1}$ (the latter tends to 1 if there is just one copy of $16_H \oplus \overline{16}_H$). Note also that the current mechanism, apart from being unified, is also quite distinct from the left-right symmetric attempts in Refs. [14,15]. From this argument it follows that the *B-L* breaking scale can in principle be as low as few TeV.

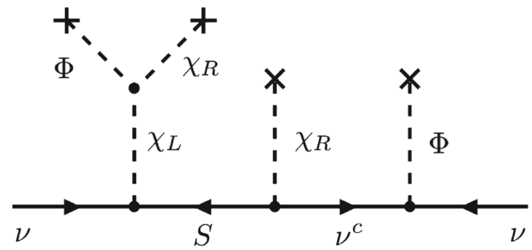

FIG. 1. The proposed supersymmetric SO(10) seesaw mechanism comes from this graph, up to transposition. The neutrino mass is suppressed by the unification scale, not by the *B-L* breaking scale, which can be low.

Concerning the stability of the texture zeros at the $\nu^c \nu^c$ and *SS* entries in formula (6) it can be protected as long as the $\mathrm{U}(1)_R$ and the $\mathrm{U}(1)_X$ [of $E_6 \supset \mathrm{SO}(10) \otimes \mathrm{U}(1)_X$ in E_6 inspired setups] are exact. Indeed, $\mathrm{U}(1)_X$ must be broken at the v_R scale [$16_H \oplus \overline{16}_H$ always has a $\mathrm{U}(1)_X$ charge]. However, the charge of χ_R's is such that the relevant operators arise only at higher orders and may be neglected.

Now we turn to gauge coupling unification. As was shown by Deshpande *et al.* [16] the scale at which the $\mathrm{SU}(2)_R \otimes \mathrm{U}(1)_{B-L}$ symmetry is broken to $\mathrm{U}(1)_Y$ can be arbitrarily low if we populate properly the desert from M_Z to M_G. In their case, this is achieved by putting *three* copies of $(1, 1, 2, +1) \oplus (1, 1, 2, -1)$ coming from $16_H \oplus \overline{16}_H$ right at the v_R scale. Then the (one-loop) minimal supersymmetric standard model (MSSM) running of the α_Y^{-1} can be "effectively" extended above the v_R scale ($\alpha_Y^{-1} = \frac{3}{5}\alpha_R^{-1} + \frac{2}{5}\alpha_{B-L}^{-1}$) by a conspiracy between the running of α_{B-L}^{-1} and α_R^{-1}. However, such a scheme is rather *ad hoc* as we need to push three identical copies of a Higgs multiplet to a very low scale, at odds with the "minimal fine-tunning."

Here we present a more compelling scheme in which the $\mathrm{SU}(2)_R$ breaking scale V_R is separated from the low $\mathrm{U}(1)_{B-L} \otimes \mathrm{U}(1)_R$ breaking scale v_R, $V_R \gg v_R$ in the chain $\mathrm{SU}(2)_R \otimes \mathrm{U}(1)_{B-L} \to \mathrm{U}(1)_R \otimes \mathrm{U}(1)_{B-L} \to \mathrm{U}(1)_Y$. At each step we assume just those multiplets needed to break the relevant symmetry. The first step is achieved by a light admixture of the $(1, 1, 3, 0)$ multiplets living in 45 and 210, while the second stage is driven by the light component of the $(1, 1, +\frac{1}{2}, -1) \oplus (1, 1, -\frac{1}{2}, +1)$ scalars [in $\mathrm{SU}(3)_c \otimes \mathrm{SU}(2)_L \otimes \mathrm{U}(1)_R \otimes \mathrm{U}(1)_{B-L}$ notation] of $16_H \oplus \overline{16}_H$. Note that to allow for such a *L-R* asymmetric setup the D parity of SO(10) must be broken.

Let us further specify the ingredients of our supersymmetric SO(10) model needed to implement the mechanism described above. As usual, we use three copies of 16_F^i to accommodate the SM fermions and for each of them we add a singlet fermion 1_F^i to play the role of S_L. A realistic fermionic spectrum requires more than one copy of 10_H Higgs multiplet. Moreover, we assume one (or more) copy of $16_H \oplus \overline{16}_H$ to implement our new supersymmetric seesaw mechanism. To prevent fast proton decay via dimension 4 operators, we assign the matter fermions in 16_F and 1_F with a discrete matter parity that forbids the mixing of 16_F and 16_H. Finally, we add a 45_H and 210_H to trigger the proper symmetry breaking pattern with no D parity below the unified scale [17–19]. The SO(10) invariant Yukawa superpotential then reads

$$W_Y = Y_{aij} 16_F^i 16_F^j 10_H^a + F_{ijk} 16_F^i 1_F^j \overline{16}_H^k. \qquad (8)$$

We do not impose other discrete symmetries to reduce the number of parameters that might, however, be welcome in connection with the doublet-triplet splitting problem in a more detailed analysis. The Higgs superpotential is

$$\begin{aligned} W_H = {} & M_{16}^{kl} 16_H^k \overline{16}_H^l + M_{10}^{ab} 10_H^a 10_H^b + M_{45} 45_H 45_H \\ & + M_{210} 210_H 210_H + \rho_{klm} 16_H^k 16_H^l 10_H^m \\ & + \overline{\rho}_{klm} \overline{16}_H^k \overline{16}_H^l 10_H^m + \sigma_{kl} 16_H^k \overline{16}_H^l 45_H \\ & + \omega_{kl} 16_H^k \overline{16}_H^l 210_H + \lambda 45_H^3 + \kappa 45_H^2 210_H \\ & + \xi 45_H 210_H^2 + \zeta 210_H^3. \end{aligned} \qquad (9)$$

The components of 210_H and 45_H that receive unified-scale VEVs and trigger the breaking of SO(10) to $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ are (in Pati-Salam language) $210_H \ni (15, 1, 1) \oplus (1, 1, 1)$ and $45_H \ni (15, 1, 1)$ As shown in [17–19], this pattern can accommodate the desired D-parity breaking allowing for an intermediate L-R symmetric group with an asymmetric particle content, leading to distinct g_L and g_R below M_G. The subsequent $SU(2)_R \to U(1)_R$ breaking at V_R is induced by the VEV of a light superposition of $(1, 1, 3, 0)_{210}$ and $(1, 1, 3, 0)_{45}$ that can mix below the unified scale. Next, the $U(1)_R \otimes U(1)_{B-L}$ is broken down at v_R by the VEVs of the light component of type $(1, 1, +\frac{1}{2}, -1) \oplus (1, 1, -\frac{1}{2}, +1)$ coming from $(1, 1, 2, -1) \oplus (1, 1, 2, 1)$ of $16_H \oplus \overline{16}_H$. The final SM breaking step is as usual provided by the VEVs of the $(1, 2, 2, 0)$ bidoublet components. Note that unlike the example given in [16] there is no artificial redundancy in the number of light states living at intermediate scales.

Let us finally inspect the one-loop gauge coupling unification. Using the normalization convention $2\pi t(\mu) = \ln(\mu/M_Z)$ we have (for $M_A < M_B$)

$$\alpha_i^{-1}(M_A) = \alpha_i^{-1}(M_B) + b_i(t_B - t_A)$$

in the ranges $[M_Z, M_S]$, $[M_S, v_R]$, and $[V_R, M_{\rm GUT}]$, M_S is the supersymmetry breaking scale taken at $\sim$1 TeV. Between v_R and V_R the two U(1) factors mix and the running of α_R^{-1} and α_{B-L}^{-1} requires separate treatment. The Cartans obey the traditional formula (with "physically" normalized B-L and Y_W) $Y_W = 2T_3^R + (B - L)$. Note that the SO(10) normalization of b_{B-L} is $b'_{B-L} = \frac{3}{8} b_{B-L}$.

Once the D parity is broken below M_G, we have $g_L \neq g_R$. The Higgs sector in the stage down to V_R is as follows: $1 \times (1, 1, 3, 0)$, $1 \times (1, 1, 2, +1) \oplus (1, 1, 2, -1)$, and $1 \times (1, 2, 2, 0)$. This gives rise to the b coefficients $b_3 = -3$, $b_L = 1$, $b_R = 4$, and $b_{B-L} = 20$.

At the subsequent stage from V_R to v_R we keep only the weak scale bidoublet $(1, 2, 2, 0)$ [that below V_R splits into a pair of L doublets with the quantum numbers $(1, 2, +\frac{1}{2}, 0) \oplus (1, 2, -\frac{1}{2}, 0)$ under the $SU(3)_c \otimes SU(2)_L \otimes U(1)_R \otimes U(1)_{B-L}$ group] and a part of $(1, 1, 2, +1) \oplus (1, 1, 2, -1)$ that is needed to break $U(1)_R \otimes U(1)_{B-L}$ to $U(1)_Y$—a pair of the χ_R fields $(1, 1, +\frac{1}{2}, -1) \oplus (1, 1, -\frac{1}{2}, +1)$. Since these fields are neutral with respect to all SM charges, the position of the v_R scale does not affect the running of the "effective α_1^{-1}" (given by the appropriate matching condition) and the only effects arise from the absence of the right-handed W bosons at this stage. Using the $SU(2)_R$ normalization of the $U(1)_R$ charge the matching condition at V_R is trivial. The relevant b coefficients of $SU(3)_c \otimes SU(2)_L$ and the matrix of anomalous dimensions of the mixed $U(1)_R \otimes U(1)_{B-L}$ couplings are $b_3 = -3$, $b_L = 1$, and

$$\begin{pmatrix} \gamma_{11} & \gamma_{12} \\ \gamma_{21} & \gamma_{22} \end{pmatrix} = \begin{pmatrix} \frac{15}{2} & -1 \\ -1 & 18 \end{pmatrix}. \qquad (10)$$

Below the v_R scale the model is the ordinary MSSM with the b coefficients $b_3 = -3$, $b_L = 1$, and $b_Y = 33/5$, and finally, the b coefficients for the SM stage below the M_S scale are $b_3 = -7$, $b_L = -3$, and $b_Y = 21/5$. The v_R-scale matching condition reads $\alpha_Y^{-1}(v_R) = \frac{3}{5}\alpha_R^{-1}(v_R) + \frac{2}{5}\alpha_{(B-L)'}^{-1}(v_R)$. Recalling that $\alpha_1^{-1}(M_Z) = \frac{3}{5}(1 - \sin^2\theta_W)\alpha^{-1}(M_Z)$ and $\alpha_2^{-1}(M_Z) = \sin^2\theta_W \alpha^{-1}(M_Z)$ the initial condition (for central values

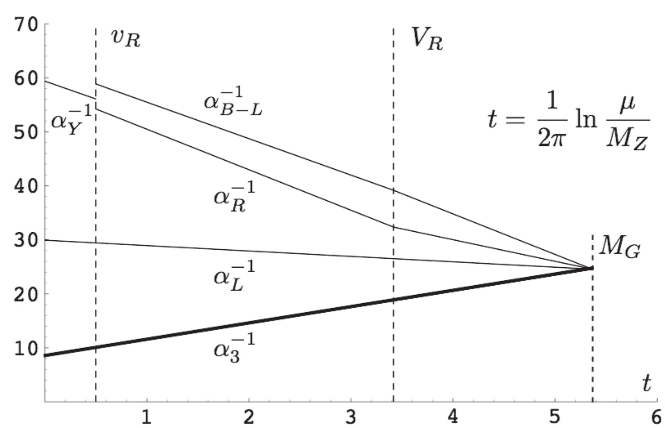

FIG. 2. The one-loop gauge coupling unification in the model described in the text. The D parity is broken at M_G, and the intermediate scales V_R and v_R correspond to $SU(2)_R \to U(1)_R$ and $U(1)_R \otimes U(1)_{B-L} \to U(1)_Y$ breaking, respectively.

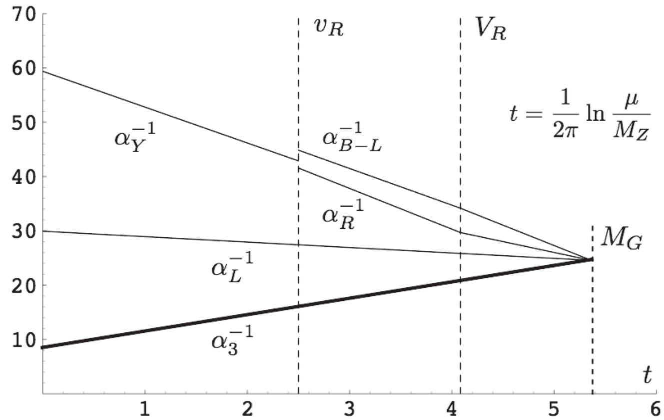

FIG. 3. Same as in Fig. 2 for the case of higher v_R. As expected, the prediction for $\alpha_1^{-1}(M_Z)$ does not depend on the position of the v_R scale.

of the input parameters) is $\alpha_1^{-1}(M_Z) \doteq 59.38$, $\alpha_2^{-1}(M_Z) \doteq 29.93$, and $\alpha_3^{-1}(M_Z) \doteq 8.47$ [20].

Inspecting the results of the numerical analysis (Figs. 2 and 3) one confirms that the v_R scale does not affect the predicted value of $\alpha_1^{-1}(M_Z)$ and remains essentially free at the one-loop level. Thus, the unification pattern is fixed entirely by the interplay of M_S and V_R. The lower bound $V_R \gtrsim 10^{14}$ GeV is consistent with the standard minimally fine-tuned supersymmetry SO(10) behavior; see for instance [21].

In summary, we have proposed a variant supersymmetric SO(10) seesaw mechanism, which involves a dynamical scale v_R that can be rather low, as it is not strongly restricted either by gauge coupling unification or neutrino masses. The smallness of neutrino masses coexists with a light *B-L* gauge boson, possibly at the TeV scale, that can be produced at the Large Hadron Collider, by the Drell-Yan process. Moreover, the "heavy" neutrinos involved in the seesaw mechanism [see Eq. (6)] get masses at v_R and can therefore be sufficiently light as to bring a rich set of testable implications. For example, their exchange can induce flavor violating processes, such as $\mu \to e\gamma$ with potentially very large rates, similar to the inverse seesaw model of Eq. (3) [10,11]. We conclude that, even in an unified seesaw context, the dynamics underlying neutrino masses may have observable effects at accelerators and in the flavor sector.

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Note Added.—Recently, we came across related papers by Steve Barr and collaborators and by Fukuyama *et al.* [22]. These indeed have strong elements in common with the early work in Refs. [9,10]. However, the mechanism we now propose differs crucially from all of these in that our *B-L* scale can be very low, in contrast to theirs. We show how our key new feature not only accounts for the observed neutrino mass scale, but also fits with the gauge unification condition in SO(10).

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