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Utilitarian supersymmetric gauge model of particle interactions

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A remarkabale $U(1)$ gauge extension of the supersymmetric standard model was proposed 8 years ago. It is anomaly free, has no μ term, and conserves baryon and lepton numbers automatically. The phenomenology of a specific version of this model is discussed. In particular, leptoquarks are predicted, with couplings to the heavy singlet neutrinos, the scalar partners of which may be components of dark matter. The Majorana neutrino mass matrix itself may have two zero subdeterminants.

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I. INTRODUCTION

The transition from the standard model (SM) of particle interactions to the minimal supersymmetric standard model (MSSM) is fraught with two well-known shortcomings. (1) Whereas baryon number B and lepton number L are automatically conserved in the SM, they are conserved in the MSSM only by the imposition of R parity, i.e. $R \equiv (-1)^{2j+3B+L}$. (2) There exists the term $\mu\phi_1\phi_2$ in the MSSM superpotential, where $\phi_{1,2}$ are the two Higgs superfields which spontaneously break the electroweak gauge symmetry. Since this term is allowed by the gauge symmetry and the supersymmetry, there is no understanding of why μ should be of the order of the electroweak breaking scale, rather than some very large unification scale.

It is clearly desirable and useful to have a single mechanism which solves both problems. One such utilitarian proposal was made 8 years ago [1], using a new $U(1)_X$ gauge symmetry. Let the quark and doublet superfields transform as n_1 and n_4 respectively under $U(1)_X$. Requiring the absence of anomalies, two classes of solutions for the other superfield assignments are then obtained as functions of n_1 and n_4. In this paper, the particularly simple choice of $n_1 = 0$ in solution (A) is discussed, together with some of its phenomenology, relevant to the operating Large Hadron Collider (LHC).

II. MODEL

Consider the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X$ with the particle content of Ref. [1]. For $n_1 = 0$ in solution (A), the various superfields transform as shown in Table I. There are three copies of Q, u^c, d^c, L, e^c, N^c, S_1, S_2; two copies of U, U^c, S_3; and one copy of ϕ_1, ϕ_2, D, D^c. The only allowed terms of the superpotential are thus trilinear, i.e.

$$Qu^c\phi_2, \quad Qd^c\phi_1, \quad Le^c\phi_1, \quad LN^c\phi_2, \quad S_3\phi_1\phi_2, \quad N^cN^cS_1, \tag{1}$$

$$S_3UU^c, \quad S_3DD^c, \quad u^cN^cU, \quad u^ce^cD, \quad d^cN^cD, \quad QLD^c, \quad S_1S_2S_3. \tag{2}$$

The absence of any bilinear term means that all masses come from soft supersymmetry breaking, thus explaining why the $U(1)_X$ and electroweak symmetry breaking scales are not far from that of supersymmetry breaking. As $S_{1,2,3}$ acquire nonzero vacuum expectation values (VEVs), the exotic (U, U^c) and (D, D^c) fermions obtain Dirac masses from $\langle S_3\rangle$, which also generates the μ term. The singlet N^c fermion gets a large Majorana mass from $\langle S_1\rangle$, so that the neutrino ν gets a small seesaw mass in the usual way. The singlet $S_{1,2,3}$ fermions themselves get Majorana masses from their scalar counterparts $\langle S_{1,2,3}\rangle$ through the $S_1S_2S_3$ terms. The only massless fields left are the usual quarks and leptons. They then become massive as $\phi^0_{1,2}$ acquire VEVs, as in the MSSM.

Because of $U(1)_X$, the structure of the superpotential conserves both B and $(-1)^L$, with $B = 1/3$ for Q, U, D, and $B = -1/3$ for u^c, d^c, U^c, D^c; $(-1)^L$ odd for L, e^c, N^c, U, U^c, D, D^c, and even for all others. Hence the exotic U, U^c, D, D^c scalars are leptoquarks and decay into ordinary quarks and leptons. The R parity of the MSSM is defined here in the same way, i.e. $R \equiv (-)^{2j+3B+L}$, and is con-

TABLE I. Particle content of proposed model.

Superfield	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_X$
$Q = (u, d)$	3	2	1/6	0
u^c	3^*	1	−2/3	3/2
d^c	3^*	1	1/3	3/2
$L = (\nu, e)$	1	2	−1/2	1
e^c	1	1	1	1/2
N^c	1	1	0	1/2
ϕ_1	1	2	−1/2	−3/2
ϕ_2	1	2	1/2	−3/2
S_1	1	1	0	−1
S_2	1	1	0	−2
S_3	1	1	0	3
U	3	1	2/3	−2
D	3	1	−1/3	−2
U^c	3^*	1	−2/3	−1
D^c	3^*	1	1/3	−1

served. Note also that the quadrilinear terms $QQQL$ and $u^c u^c d^c e^c$ (allowed in the MSSM) as well as $u^c d^c d^c N^c$ are forbidden by $U(1)_X$. Proton decay is thus strongly suppressed. It may proceed through the quintilinear term $QQQLS_1$ as the S_1 fields acquire VEVs, but this is a dimension-six term in the effective Lagrangian, which is suppressed by two powers of a very large mass, say the Planck mass, and may safely be allowed.

III. GAUGE SECTOR

The new Z_X gauge boson of this model becomes massive through $\langle S_{1,2,3}\rangle = u_{1,2,3}$, whereas $\langle \phi^0_{1,2}\rangle = v_{1,2}$ contribute to both Z and Z_X. The resulting 2×2 mass-squared matrix is given by [2]

$$\mathcal{M}^2_{Z,Z_X} = \begin{pmatrix} (1/2)g_Z^2(v_1^2 + v_2^2) & (3/2)g_Z g_X(v_2^2 - v_1^2) \\ (3/2)g_Z g_X(v_2^2 - v_1^2) & 2g_X^2[u_1^2 + 4u_2^2 + 9u_3^2 + (9/4)(v_1^2 + v_2^2)] \end{pmatrix}. \tag{3}$$

Since precision electroweak measurements require $Z - Z_X$ mixing to be very small [3], $v_1 = v_2$, i.e. $\tan\beta = 1$, is assumed from now on.

Consider the decay of Z_X to the usual quarks and leptons. Each fermionic partial width is given by

$$\Gamma(Z_X \to \bar{f}f) = \frac{g_X^2 M_{Z_X}}{24\pi}[c_L^2 + c_R^2], \tag{4}$$

where $c_{L,R}$ can be read off under $U(1)_X$ from Table I. Thus

$$\frac{\Gamma(Z_X \to \bar{t}t)}{\Gamma(Z_X \to \mu^+\mu^-)} = \frac{\Gamma(Z_X \to \bar{b}b)}{\Gamma(Z_X \to \mu^+\mu^-)} = \frac{27}{5}. \tag{5}$$

This will serve to distinguish it from other Z' models [4].

IV. EFFECTIVE TWO-HIGGS-DOUBLET STRUCTURE

In the MSSM, the scalar potential of the two Higgs doublets is given by

$$V = m_1^2\Phi_1^\dagger\Phi_1 + m_2^2\Phi_2^\dagger\Phi_2 + m_{12}^2(\Phi_1^\dagger\Phi_2 + \Phi_2^\dagger\Phi_1) + \tfrac{1}{2}\lambda_1(\Phi_1^\dagger\Phi_1)^2 + \tfrac{1}{2}\lambda_2(\Phi_2^\dagger\Phi_2)^2 + \lambda_3(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) + \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1), \tag{6}$$

where

$$\lambda_{1,2} = \tfrac{1}{4}(g_1^2 + g_2^2), \qquad \lambda_3 = \tfrac{1}{4}(-g_1^2 + g_2^2), \qquad \lambda_4 = -\tfrac{1}{2}g_2^2. \tag{7}$$

In the present model, there are extra Higgs singlets, but if they are heavier than the doublets by an order of magnitude and the soft A terms are of the elctroweak scale, they can be integrated out and the effective two-Higgs-doublet structure is given by [2,5–9]

$$\lambda_{1,2} = \frac{1}{4}(g_1^2 + g_2^2) + f^2 - \frac{f^4}{9g_X^2}, \tag{8}$$

$$\lambda_3 = \frac{1}{4}(-g_1^2 + g_2^2) + f^2 - \frac{f^4}{9g_X^2}, \tag{9}$$

$$\lambda_4 = -\tfrac{1}{2}g_2^2 + f^2, \tag{10}$$

where f is the Yukawa coupling of the trilinear term $S_3\phi_1\phi_2$, assuming that this one particular singlet dominates over all others in changing the scalar quartic couplings λ_i.

Since $\tan\beta = 1$ in this model, the lightest neutral Higgs boson has the upper bound [2]

$$(m_h^2)_{\max} = \epsilon + \frac{f^2}{\sqrt{2}G_F}\left[\frac{3}{2} - \frac{f^2}{9g_X^2}\right], \tag{11}$$

where

$$\epsilon = \frac{3g_2^2 m_t^4}{8\pi^2 M_W^2}\ln\left(1 + \frac{\tilde{m}^2}{m_t^2}\right) \tag{12}$$

is the well-known large radiative correction [10–14] due to the t quark and its supersymmetric scalar partners. The above upper bound may easily exceed that of the MSSM. For example, let $f = 3g_X = g_2$, then $(m_h)_{\max} = 2M_W^2 + \epsilon = 143$ GeV, assuming $\tilde{m} = 1$ TeV in Eq. (12). Contrast this with the upper bound in the MSSM, i.e. $M_Z^2\cos^2 2\beta + \epsilon < (126\ \text{GeV})^2$. If the experimental bound of $m_h > 114.4$ GeV is used, then $|\cos 2\beta| > 0.81$ is required. Here the prediction is that $\cos 2\beta = 0$ and yet m_h may be substantially greater than 114.4 GeV. This difference will have important implications for the Higgs search at the LHC. Another important difference is the mass of the charged Higgs boson:

$$m_{H^\pm}^2 = m_A^2 + M_W^2 - f v^2, \tag{13}$$

where m_A is the mass of the pseudoscalar Higgs boson, i.e. $\mathrm{Im}(\phi_1^0 - \phi_2^0)$. For example, if $f = g_2$, then $m_{H^\pm}^2 - m_A^2 = -M_W^2$, instead of $+M_W^2$ in the MSSM.

V. NEUTRINO MASSES

There are three copies each of the superfields L, e^c, N^c, and S_1. As such, a family symmetry in the lepton sector may be supported. For example, the discrete symmetry Z_4 may be used to realize the interesting proposal of Ref. [15] that the observed neutrino mass matrix has two zero subdeterminants. Under Z_4, the three copies of L, e^c, N^c, and S_1 separately all transform as 1, i, $-i$, with $\phi_{1,2}$ transforming as 1. From the $Le^c\phi_1$ and $LN^c\phi_2$ couplings, the charged-lepton and Dirac neutrino mass matrices are thus diagonal, whereas the $N^cN^cS_1$ couplings result in the Majorana mass matrix of the form [16,17]

$$\mathcal{M}_N = \begin{pmatrix} A & B & C \\ B & 0 & D \\ C & D & 0 \end{pmatrix}. \tag{14}$$

The resulting seesaw neutrino mass matrix is then given by [15,18]

$$\mathcal{M}_\nu = \begin{pmatrix} \alpha & \beta & \gamma \\ \beta & \alpha^{-1}\beta^2 & \delta \\ \gamma & \delta & \alpha^{-1}\gamma^2 \end{pmatrix}, \tag{15}$$

where the subdeterminants of the (1, 2) and (1, 3) blocks are clearly zero. Assuming $\gamma = \beta$, then $\theta_{23} = \pi/4$, $\theta_{13} = 0$, and

$$\Delta m^2_{21} \cos 2\theta_{12} = |\alpha^{-1}\beta^2 + \delta|^2 - |\alpha|^2, \tag{16}$$

$$\Delta m^2_{21} \sin 2\theta_{12} = 2\sqrt{2}|\alpha^*\beta + \beta^*(\alpha^{-1}\beta^2 + \delta)|, \tag{17}$$

$$\Delta m^2_{32} = |\alpha^{-1}\beta^2 - \delta|^2 - \tfrac{1}{2}|\alpha|^2 - \tfrac{1}{2}|\alpha^{-1}\beta^2 + \delta|^2 - 2|\beta|^2. \tag{18}$$

Since there can only be one nontrivial phase in the above, let α be real, $\epsilon = \alpha^{-1}\beta^2 + \delta$ real, and β complex. In the case β is also real, it has already been shown [15] that the resulting solution has the normal hierarchy of neutrino masses, i.e. $m_1 < m_2 < m_3$, such that the effective neutrino Majorana mass measured in neutrinoless double beta decay, i.e. the parameter α, is about 6×10^{-4} eV, much below the sensitivity of such experiments.

Consider now the case of a purely imaginary β, using $\beta = i\zeta$, where ζ is real, then a solution exists where α is several times larger, as shown below. Rewriting Eqs. (16)–(18),

$$\Delta m^2_{21} \cos 2\theta_{12} = (\epsilon - \alpha)(\epsilon + \alpha), \tag{19}$$

$$\Delta m^2_{21} \sin 2\theta_{12} = 2\sqrt{2}\zeta(\epsilon - \alpha), \tag{20}$$

$$\Delta m^2_{32} = \zeta^2(4\alpha^{-2}\zeta^2 + 4\alpha^{-1}\epsilon - 2) + \tfrac{1}{2}(\epsilon^2 - \alpha^2). \tag{21}$$

Let $\alpha = 0.0042$ eV, $\epsilon = 0.0067$ eV, and $\zeta = 0.0098$ eV, then $\Delta m^2_{21} = 7.45 \times 10^{-5}$ eV^2, $\tan^2\theta_{12} = 0.464$, and $\Delta m^2_{32} = 2.53 \times 10^{-3}$ eV^2, in good agreement with data [19].

VI. HEAVY NEUTRINO SINGLET AT THE LHC

In the canonical seesaw, even if the heavy neutrino singlet anchor has a mass of order TeV, it is very hard to produce at the LHC [20], because it couples only to leptons and the strength of that coupling is necessarily very weak, as required by the tiny neutrino mass [21]. Here, since N^c also couples to the leptoquarks (U, U^c) and (D, D^c), it can be produced at the LHC as the decay product of the latter, which are copiously produced themselves because they have strong interactions.

Since the scalar $\tilde{N}^c$ has odd R parity, it may also be a component of dark matter. In that case, the decays of the heavy leptoquark fermions (U, U^c) and (D, D^c) to quark jets and $\tilde{N}^c$ may lead to sizable missing-energy signals at the LHC, as recently discussed [22].

VII. CONCLUSION

The utilitarian supersymmetric $U(1)_X$ gauge extension of the standard model of particle interactions proposed 8 years ago [1] allows for two classes of anomaly-free models which have no μ term and conserve baryon number and lepton number automatically. A simple version with leptoquark superfields is discussed here with a number of interesting and verifiable properties.

The new Z_X gauge boson of this model has specified couplings to quarks and leptons which are distinct from other gauge extensions and may be tested at the LHC. The effective two-Higgs-doublet sector has $\tan\beta = 1$, and yet the mass of the lightest neutral Higgs boson may exceed the upper limit of 126 GeV predicted in the MSSM. A discrete Z_4 symmetry may be accommodated in the lepton sector so that the 3×3 neutrino mass matrix has two zero subdeterminants. The scalar partners of the heavy singlet neutrinos could be components of dark matter, and since they may be decay products of the leptoquark fermions, they are possible sources of missing energy at the LHC.

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