Z' boson from SU(6)×SU(2)$_h$ grand unified theory, CDF Wjj excess and Higgs boson mass bound

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A general electroweak scale Z' is applied in a supersymmetric $\mathrm{SU}(6)\times\mathrm{SU}(2)_h$ grand unification model, to have a $\mathbb{Z}_6$ for the hexality. We briefly show that there cannot exist any baryonic $\mathrm{U}(1)'_B$ in any subgroup of E_6. Any effect that requires sizable Z' couplings to quarks like the reported Wjj anomaly of CDF, if observed, implies a substantial Z' coupling to leptons or Higgs doublets. The kinetic mixing considered in a supersymmetric model from E_6 is restricted by the gauge coupling unification and neutrino mixing. The mass of Z' is strongly constrained by the electroweak ρ_0 parameter. We conclude that a Z' mass much above 10 TeV is favored by considering the neutrino mixing and proton decay constraint in supersymmetric models. In this sense, the CDF Wjj anomaly cannot be fitted to any electroweak model descending from E_6. Furthermore, if Z' is found at several hundred GeV, any grand unification group embedded in E_6 such as $\mathrm{SU}(6)\times\mathrm{SU}(2)$, SO(10), $\mathrm{SU}(5)\times\mathrm{U}(1)$, SU(5), $\mathrm{SU}(4)\times\mathrm{SU}(4)$, and $\mathrm{SU}(3)^3$, needs fine-tuned gauge couplings. We also discuss the $\mathrm{U}(1)'$ effect on the tree level mass of the lightest MSSM Higgs boson.

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In the standard model (SM), there is one electroweak scale neutral gauge boson Z [1]. Discovery of any new neutral gauge boson hints at a bigger gauge group beyond $\mathrm{SU}(3)\times\mathrm{SU}(2)\times\mathrm{U}(1)$. This possibility is widely discussed in view of the CDF reports such as Wjj final states having a bump in dijet invariant mass near 145 GeV [2].[1] The simplest extension is just assuming a new $\mathrm{U}(1)'$ beyond the SM gauge group, which does not fix the $\mathrm{U}(1)'$ quantum numbers except by the constraints from the anomaly freedom [5].

On the other hand, if the extension beyond the SM is achieved in (semi-)simple gauge groups, then the gauge quantum numbers are not arbitrary but fixed for given representations. The most analyzed grand unification (GUT) groups for $\mathrm{U}(1)'$ are SO(10) and E_6 GUTs [6]. However, there is the notorious gauge hierarchy problem in GUTs. The supersymmetry (SUSY) model was suggested to solve this problem. In addition, the doublet/triplet splitting problem is the most serious issue with the SUSY GUTs.

The doublet/triplet splitting problem in GUTs is surfaced as the μ-problem [7] in the minimal SUSY SM (MSSM). There are several ways to solve the μ-problem in some extension of the MSSM. In particular, the string solutions seem to be interesting because they touch upon all other plausible phenomenological aspects of the MSSM from the ultraviolet completed theory [8]. For instance, a SUSY electroweak group $\mathrm{SU}(3)_W\times\mathrm{U}(1)$ is exceptionally useful for obtaining one pair of Higgs doublets H_u and H_d in the MSSM, naturally solving the μ-problem [9].

In this paper, we present a $\mathrm{U}(1)'$ model from a SUSY $\mathrm{SU}(6)\times\mathrm{SU}(2)_h$ GUT. For the $\mathrm{U}(1)'$ phenomenology from E_6, the discussion from its subgroup $\mathrm{SU}(6)\times\mathrm{SU}(2)_h$ is as good as E_6 since their ranks are the same. From the chain of GUTs, $\mathrm{SU}(5)\subset\mathrm{SO}(10)\subset\mathrm{E}_6$, the discussion of SO(10) is included in the discussion of E_6. Note also that the extra Z' from SO(10) is in fact the Z' of B-L, which is not the one we try to introduce for the recent Wjj anomaly. Any $\mathrm{U}(1)'$ generator can be written as a linear combination of six E_6 Cartan generators or of six $\mathrm{SU}(6)\times\mathrm{SU}(2)_h$ Cartan generators. If the SU(6) is taken as a GUT, the electroweak part is $\mathrm{SU}(3)_W\times\mathrm{U}(1)$ [10]. Its SUSY extension was obtained from the F-theory compactification of string [11]. Note, however, that in our $\mathrm{U}(1)'$ discussion, the SU(6) GUT is not a necessity except from the proton stability condition. The representation of $\mathrm{SU}(6)\times\mathrm{SU}(2)_h$ can be shown as matrix elements on the plane without any attachment of U(1) quantum numbers. This is a nice feature to glimpse the Z' quantum numbers, just by looking at the representation on the plane. From these representations, we will notice that there are only two neutral SM singlets N and N' which are the heavy neutrinos needed for the seesaw mechanism. The chief motivation for the SUSY GUT group containing SU(6) is from the proton hexality condition that forbids proton decay operators of dimension-4 and dimension-5 terms [12]. The R-parity forbids the dimension-4 operator from the superpotential $u^cd^cd^c$ but allows the dimension-5 proton decay operator from the superpotential $qqq\ell$. The SU(5) GUT does not forbid this dimension-5 proton decay operator, and its coefficient is required to be as small as 0.995×10^{-8} considering the limit of the proton decay to

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[1]It must be also noticed that there is no such bump in the recent D0 result with 4.3 fb^{-1} [3] and the LHC result at the 1 fb^{-1} [4] integrated luminosity. With these results, the modeling of the SM background can be important as some of the papers in [2].

$K^+\bar{\nu}$ in [13]. The hexality of Ref. [12] is the product of R-parity and triality which forms a $\mathbb{Z}_6$. The operator $qqql$ is allowed by the R-parity but is forbidden by the triality since four triality nonsinglet fields (with q's having the same triality) are multiplied. The reason SU(5) cannot accommodate the hexality is that it does not have a discrete subgroup $\mathbb{Z}_6$. On the other hand, GUTs containing SU(6) can have the $\mathbb{Z}_6$ discrete subgroup since the center of SU(6) is $\mathbb{Z}_6$.

The SU(6) $\times$ SU$(2)_h$ SUSY model allows the following representations, e.g. for the first family,

$$\mathbf{15}_L \equiv (\mathbf{15},\mathbf{1}) = \begin{pmatrix} 0 & u^c & -u^c & u & d & D \\ -u^c & 0 & u^c & u & d & D \\ u^c & -u^c & 0 & u & d & D \\ -u & -u & -u & 0 & e^c & H_u^+ \\ -d & -d & -d & -e^c & 0 & H_u^0 \\ -D & -D & -D & -H_u^+ & -H_u^0 & 0 \end{pmatrix},$$

$$\bar{\mathbf{6}}_{\mathbf{2},1} \equiv (\bar{\mathbf{6}},\mathbf{2}^{\uparrow}) = \begin{pmatrix} d^c \\ d^c \\ d^c \\ -\nu_e \\ e \\ N \end{pmatrix}, \quad \bar{\mathbf{6}}_{\mathbf{2},2} \equiv (\bar{\mathbf{6}},\mathbf{2}^{\downarrow}) = \begin{pmatrix} D^c \\ D^c \\ D^c \\ -H_d^0 \\ H_d^- \\ N' \end{pmatrix}. \tag{1}$$

Note that we have not included any E_6 singlets.[2] The representations $\bar{\mathbf{6}}_{\mathbf{2},1}$ and $\bar{\mathbf{6}}_{\mathbf{2},2}$ form a doublet pair of the horizontal group SU$(2)_h$. Without the loss of generality, we choose $\bar{\mathbf{6}}_{\mathbf{2},1}$ as matter and $\bar{\mathbf{6}}_{\mathbf{2},2}$ as Higgs sextets. We need three families and at least a vector-like pair $n_{(\bar{\mathbf{6}},\mathbf{2})}$ and $n_{(\mathbf{6},\mathbf{2})}$, which is responsible for the breaking SU(6) down to SU(5). Therefore, to allow for three chiral families and SU(6) $\to$ SU(5) breaking, we assume $n_{(\bar{\mathbf{6}},\mathbf{2})} = 4$ and $n_{(\mathbf{6},\mathbf{2})} = 1$ [11]. By the vacuum expectation value (VEVs) of $\bar{\mathbf{6}}_{\mathbf{2}}$ and $\mathbf{6}_{\mathbf{2}}$, SU(6) $\times$ SU$(2)_h$ is broken down to SU(5) $\times$ U$(1)'$.

By the VEV of the adjoint representation **35** of SU(6) or by the hyper-flux in F-theory [15], our interest is focused on a rank 6 group SU$(3)_c \times$ SU$(3)_W \times$ U(1) $\times$ SU$(2)_h$.[3] To break SU$(3)_W \times$ U(1) $\times$ SU$(2)_h$ down to the one including the rank 4 SM group SU$(3)_c \times$ SU$(2)_W \times$ U$(1)_Y$, we assign a GUT scale VEVs to $(\bar{\mathbf{6}}, \mathbf{2})$ and $(\mathbf{6}, \mathbf{2})$, which reduce just one rank, and the low-energy gauge group is SU$(3)_c \times$ SU$(2)_W \times$ U$(1)_Y \times$ U$(1)'$. The tensor form of the representation $(\bar{\mathbf{6}}, \mathbf{2})$ is $\bar{\mathbf{6}}_i^{\alpha}$, where $\alpha = 1, 2, \cdots, 6$ and $i = 1, 2$. The fields of Eq. (1) couple as

$$f\mathbf{15}_{\alpha\beta}\bar{\mathbf{6}}_i^{\alpha}\bar{H}_j^{\beta}\epsilon^{ij}, \tag{2}$$

where we suppressed the family indices and used $\bar{H}$ as another $(\bar{\mathbf{6}}, \mathbf{2})$. The group SU$(2)_W \times$ U$(1)_Y \times$ U$(1)'$ contains three diagonal generators, $Q_{\rm em}$, Q_Z, and Y'. The SM $Q_{\rm em}$ and Q_Z are included in the GUT SU(6), representing the linear combinations of two SU(6) generators only in the vertical directions of $(\mathbf{6}, \mathbf{2})$: $T_3 = {\rm diag.}(000\frac{1}{2}\,\frac{-1}{2}0)$ and $Y = {\rm diag.}(\frac{-1}{3}\,\frac{-1}{3}\,\frac{-1}{3}\,\frac{1}{2}\,\frac{1}{2}0)$. The U$(1)'$ generator is a linear combination of two SU(6) $\times$ SU$(2)_h$ diagonal generators in the vertical and horizontal directions of $(\mathbf{6}, \mathbf{2})$: $Y_{\rm SU(6)}$ and X_3.

Let the gauge bosons corresponding to T_3, Y, and Y' be A_μ^3, B_μ, and C_μ (with coupling g''), respectively. Below, we will present the form of Y'. The mass eigenstates are defined as the photon A_μ, Z_μ-boson and Z'_μ-boson. In this extended weak interaction model,[4] we define a new weak mixing angle $\sin^2\varphi = g''^2/(g^2 + g'^2 + g''^2)$ in addition to $\sin^2\theta_W = g'^2/(g^2 + g'^2) \simeq 0.23$ of the SM. The diagonal gauge bosons of SU$(2)_W \times$ U$(1)_Y \times$ U$(1)'$ (A_μ^3, B_μ, C_μ) are related to the mass eigenstate gauge bosons (A_μ, Z_μ, Z'_μ) by an orthogonal matrix,

$$\begin{pmatrix} A_\mu^3 \\ B_\mu \\ C_\mu \end{pmatrix} = \begin{pmatrix} s_\theta & -c_\theta c_\varphi & c_\theta s_\varphi \\ c_\theta & s_\theta c_\varphi & -s_\theta s_\varphi \\ 0 & s_\varphi & c_\varphi \end{pmatrix} \begin{pmatrix} A_\mu \\ Z_\mu \\ Z'_\mu \end{pmatrix}, \tag{3}$$

where $s_\theta = \sin\theta_W$ and $s_\varphi = \sin\varphi$, and similarly for the cosines. The gauge boson masses depend on the Y' quantum numbers of Higgs fields, which will be discussed below.

Below, we prove the no-go theorem for a gauged U$(1)'_B$ from E_6 and its consequence on the Z boson and the lightest CP-even Higgs boson masses. Finally, we comment on the possibility of obtaining SU(6) $\times$ SU$(2)_h$ from the ultraviolet completed superstring.

On gauged U$(1)'_B$ *and leptophobic* U$(1)'$ *from* E_6: The chiral representation **27** of E_6 is split into $(\mathbf{15}, \mathbf{1})$ and $(\bar{\mathbf{6}}, \mathbf{2})$ of Eq. (1). Rank 6 E_6 has six diagonal generators: F_3 and F_8 of SU$(3)_c$, T_3 of SU$(2)_W$, Y of U$(1)_Y$, $Y_{\rm SU(6)} \equiv Y_6$, and X_3. In any subgroup of E_6, the diagonal generators are linear combinations of these. Therefore, without loss of generality, we consider the baryon number as a linear combination of Y, Y_6, and X_3. To have an R-parity, we include a global U$(1)_R$ symmetry and consider the following U$(1)'_B$

$$B = aY + bY_6 + cX_3 + dR, \tag{4}$$

where $Y_6 \equiv Y_{\rm SU(6)} = {\rm diag.}(\frac{-1}{6}\,\frac{-1}{6}\,\frac{-1}{6}\,\frac{-1}{6}\,\frac{-1}{6}\,\frac{5}{6})$ is for the representation $\mathbf{6}$, $X_3 \equiv {\rm diag.}(\frac{1}{2} - \frac{1}{2})$ is the third generator of SU$(2)_h$, and R is the U$(1)_R$ charge. The R-symmetry is broken at the high-energy scale, we set $d = 0$ and the resulting B would be a gauge group generator.

[2] According to the recent LHC result with the 1 fb^{-1} integrated luminosity, there are new strict lower bounds on the masses of colored exotics by the study of dijet resonances. The bounds are 2.91 TeV for the excited quarks, 3.21 TeV for axigluons, and 1.91 TeV for color octet scalars [14].

[3] Without confusion, we can use the GUT representations to simplify the notation at the scale where the broken group is effective.

[4] Even if only Z' survives down to the electroweak scale, we call it a new weak interaction model.

If B is a good symmetry, leptons and Higgs fields should carry vanishing B. In addition, u^c and d^c must carry the same B, which is opposite to that of the quark doublet $(u, d)_L$. The required conditions of leptons and Higgs fields are

$$\begin{aligned} e^c&:\ a - \tfrac{1}{3}b = 0, \\ (\nu, e)&:\ -\tfrac{1}{2}a + \tfrac{1}{6}b + \tfrac{1}{2}c = 0, \\ H_d&:\ -\tfrac{1}{2}a + \tfrac{1}{6}b - \tfrac{1}{2}c = 0, \\ H_u&:\ +\tfrac{1}{2}a + \tfrac{2}{3}b = 0, \end{aligned} \tag{5}$$

which cannot be satisfied unless $a = b = c = d = 0$. Therefore, it is not possible to have a gauged $U(1)'_B$ as a subgroup of E_6.[5]

But, not requiring a strict baryon number, it is possible to consider a useful nonbaryonic $U(1)'$ from E_6. It is the so-called "leptophobic" that leptons do not have the $U(1)'$ interaction, *i.e.* no $U(1)'$ charges for (ν, e) and e^c. Since H_d has the same charge as (ν, e) in the MSSM, H_d should not carry a $U(1)'$ charge. Therefore, by adopting the first three conditions only in Eq. (5), a solution $a = \frac{1}{3}b$ and $c = 0$ is obtained. Making H_u not the complex-conjugate of H_d, *i.e.* going beyond the SM, we note that this can be realized in the so-called two Higgs doublet model with H_u carrying a nonvanishing $U(1)'$ charge. In this case, the diagonal entries of the leptophobic $U(1)'$ charge is

$$Y'_{\mathrm{lp-phob}} = \frac{5}{6}\left(\frac{-1}{3}, \frac{-1}{3}, \frac{-1}{3}, 0, 0, 1\right). \tag{6}$$

There are two ways to realize this $Y'_{lp\text{-}phob}$. One is to introduce a VEV of the adjoint **78** of E_6, and the other is by considering the kinetic mixing between B_μ and C_μ in our model. Needing the adjoint representation **78** is very much involved in the orbifold construction [17], but is easily achievable in an F-theory construction.

Even without the VEV of **78**, the kinetic mixing between B_μ and C_μ has been considered with the branching $E_6 \to SO(10) \times U(1)_\psi \to SU(5) \times U(1)\chi \times U(1)_\psi$, including the running of the gauge couplings [18]. However, this case needs an extreme fine-tuning between masses of the split multiplet members to obtain such a large mixing. In addition, the leptophobia obtained by the kinetic mixing with coefficient $1/3$ is not achieved if one requires an anomaly-free model where the gauge couplings of the SM and $U(1)'$ are perturbative and unify at the GUT scale [19].

On the other hand, to give singlet neutrino masses, N of Eq. (1) should develop a VEV, implying that $Y'_{lp\text{-}phob}$ is broken at the heavy neutrino mass scale. Therefore, the exact leptophobic Z' from E_6 should be very heavy to induce the neutrino mixing, which cannot explain the recent Wjj anomaly of CDF. In this sense, our $U(1)'$ is introduced not to be leptophobic by assigning no charge to N so that the Z' charge $Y' = X_3 + \frac{3}{5}Y_6$.

[5] The $U(1)'_B$ model such as [16] is not originated from E_6.

Therefore, let us consider Z' from E_6, coupling to baryons, couples to leptons as well.

Z' and Higgs boson masses: If only the third family members have VEVs, without loss of generality, we can choose $\langle N\rangle = V_{\mathrm{heavy}}$ and $\langle N'\rangle = 0$. We also introduce at least one vector-like representations $(\mathbf{6}, \mathbf{2})$ and $(\bar{\mathbf{6}}, \mathbf{2})$ as in Ref. [11]. Since there is no parameter space where the leptonic $U(1)'$ currents are negligible, the high-precision NC experiments and the LEP II data for nonvanishing Z'- lepton coupling stringently restrict the Z' mass.

The neutral fields carrying nonvanishing $U(1)'$ charges are $H^0_{u,d}$, ν, and N'. For Z' to survive down to the electroweak scale, N' should not develop a superheavy VEV above the electroweak scale. However, for the neutrino oscillation, we also need them to be heavy [11] with mass larger than 10^{10} GeV. The N' Majorana mass can be generated by $\frac{1}{M_P}(\mathbf{6}, 2)(\mathbf{6}, 2)(\bar{\mathbf{6}}, 2)(\bar{\mathbf{6}}, 2)$ as in Ref. [11] by the VEV $\langle(\mathbf{6}, \mathbf{2})\rangle \to \langle\bar{H}^6_1\rangle = \langle N\rangle$. Using Eq. (2), the Dirac mass between ν and N' is generated by $\langle\bar{H}_u\rangle$. These lead to the seesaw mechanism for neutrino masses. The $H^0_{u,d}$ and ν_e fields at the electroweak scale carry nonvanishing $U(1)'$ charges. For the R-parity conservation, ν_e is not required to break $U(1)'$. In addition, three N' fields survive down to the electroweak scale because N' and $\bar{N}'$ fields in four $\bar{\mathbf{6}}$'s and one $\mathbf{6}$ remove only one heavy Dirac neutrino field, viz. by $\frac{1}{M_P}(\mathbf{6}, \mathbf{2})(\mathbf{6}, \mathbf{2})(\bar{\mathbf{6}}, \mathbf{2})(\bar{\mathbf{6}}, \mathbf{2})$,

$$\begin{array}{c} \\ N'_1 \\ N'_2 \\ N'_3 \\ N'_4 \\ \bar{N}' \end{array} \begin{array}{c} \begin{array}{ccccc} N'_1 & N'_2 & N'_3 & N'_4 & \bar{N}' \end{array} \\ \begin{pmatrix} 0 & 0 & 0 & 0 & M_1 \\ 0 & 0 & 0 & 0 & M_2 \\ 0 & 0 & 0 & 0 & M_3 \\ 0 & 0 & 0 & 0 & M_4 \\ M_1^* & M_2^* & M_3^* & M_4^* & M \end{pmatrix} \end{array}, \tag{7}$$

where the masses M and M_{1-4} are at the intermediate scale. The electroweak singlet neutrinos are called $N'_i(i=1,2,3)$ again. The VEVs of H_u, H_d, and $N'_i(i=1,2,3)$ break the $U(1)'$ symmetry at the electroweak scale. Since we look for the parameter space, where g'' is smaller than the $SU(2)_W$ coupling g, the contribution of the H_u and H_d VEVs to $M_{Z'}$ is smaller than their contribution to the W boson mass. Therefore, there must be a TeV scale VEV(s) of $N'_i(i = 1, 2, 3)$ to make Z' as heavy as 150 GeV, if the CDF Wjj rate is attributed to Z'. This additional VEVs are free parameters to tune the Z' mass. Since $Y_6 = \frac{2}{5}$ and $-\frac{2}{5}$ for H_u and H_d, respectively, the Z-Z' mass matrix becomes as following. Here s_φ, c_φ are defined in Eq. (3), $t_\varphi = s_\varphi/c_\varphi$, $\tan^2\gamma \equiv X^2/V^2$, $G \equiv \sqrt{g^2 + g'^2}$, $V^2 \equiv v_u^2 + v_d^2$ and X^2 is the contribution from the VEVs of N' fields.

$$M^2_{ZZ'} = \frac{G^2V^2}{4}\begin{pmatrix} c^2_\varphi + \frac{8}{5}s^2_\varphi + \left(\frac{16}{25} + 4\tan^2\gamma\right)t^2_\varphi s^2_\varphi, & -\frac{1}{5}c_\varphi s_\varphi + 4t_\varphi s^2_\varphi \tan^2\gamma - \frac{4}{25}t_\varphi s^2_\varphi \\ [0.5em] -\frac{1}{5}c_\varphi s_\varphi + 4t_\varphi s^2_\varphi\tan^2\gamma - \frac{4}{25}t_\varphi s^2_\varphi, & \left(\frac{16}{25} + 4\tan^2\gamma\right)s^2_\varphi - \frac{3}{5}s^2_\varphi \end{pmatrix}. \quad (8)$$

From Eq. (2), we note that the N' VEVs break the R-parity. If any four fields of N' in Eq. (7) does not develop a VEV, we can consider the limit $X^2 \ll v_d^2 \ll v_u^2$, *i.e.* $\tan\gamma \simeq 0$. Generally, this case leads to $M_{Z'}$ smaller than M_Z, and we are left with a large $\tan\gamma$ case. Not to be conflicted with the R-parity problem, a large $\tan\gamma$ must be provided by the heavy pair of N' and $\bar{N}'$. The VEV of the 4th N' combines the lepton doublets with the superheavy H_u. This case is not ruled out obviously. Even for this large $\tan\gamma$, the ρ parameter constrains the allowed mass of Z'. For this study, we satisfy the electroweak neutral current (NC) parameter $\rho_0 = 1.0004^{+0.0029}_{-0.0011}$ with the 2σ limit, which has no meaningful bound on the Higgs mass [13]. We show the allowed $\tan\varphi$ and M'_Z in the region $g''^2 < g'^2$ in Fig. 1(a), from which we note that the heavy Z' much above 10 TeV is favored in the region $M_{Z'} > M_Z$. Adding to this, there are more constraints such as Z boson decay width, $e^-e^+ \to W^-W^+$, etc. However, as seen in Fig. 1(a), the constraint on the ρ_0 parameter provides a strong enough conclusion to constrain the viable Z' parameters. Therefore, we concentrated only on the ρ_0 parameter in this paper.

The VEVs of H_u and H_d are related to the Higgs boson masses and can raise the upper bound of the lightest Higgs boson mass of the MSSM, even before including the radiative corrections [20]. If Z' is present and the Higgs doublets, H_u and H_d, carry nonzero U(1)$'$ charges, the lightest CP-even Higgs boson mass bound is changed. Our interest based on the ρ parameter constraint is Case $H2$ of Eq. (20) of Ref. [21]. In the limit $\tan\beta \equiv v_u/v_d \to \infty$, it can be shown succinctly. In the MSSM, we have $v_u^2 = 8\mu_u^2/G^2$ and hence $m_h^2 \simeq M_Z^2$ for $v_d \ll v_u$. With the Y_6 contribution in the D-term potential, $-\mu_u^2 H_u^\dagger H_u + \frac{1}{4}G^2(1 + \frac{16}{25}\tan^2\varphi)(H_u^\dagger H_u)^2 + (H_d\text{terms})$, we obtain $m_H^2 = 2\mu_u^2 + \cdots$ and $M_Z^2 = \frac{1}{4}G^2V^2$. Then, the Z boson mass has the same expression as in the MSSM, but the relation between the VEV v_u^2 and μ_u^2 is changed to $v_u^2 = 8\mu_u^2/G^2(1 + \frac{16}{25}\tan^2\varphi)$ if we neglect the Higgsino mixing μ term and obtain the tree level upper bound on the Higgs boson mass as, $m_H^2 = M_Z^2(1 + \frac{16}{25}\tan^2\varphi)$.

However, in this limit, the bound represents the heavier CP-even Higgs since the real part of H_d^0 is massless, if we neglected its mass parameter μ_d^2. If Z' is decoupled at high energy for $X \gg V$, still the MSSM Higgs boson masses are strongly affected. The reason is that the lightest MSSM Higgs boson mass encodes the quartic couplings or the gauge symmetry (whether it is broken or not). The lightest Higgs mass bound is not protected by the decoupling theorem since the dimensionless quartic couplings are renormalized only logarithmically.

In the limit $\tan\beta \equiv v_u/v_d \to \infty$, the pseudoscalar mass m_A goes to zero as commented above. So, we must consider a finite $\tan\beta$ case, *i.e.* for nonzero v_d and also for the Higgsino mixing term μ [7] to make the pseudoscalar heavy. In this case, we consider the following 2×2 CP even Higgs mass matrix

$$M^2_{\rm CPeven} = \begin{pmatrix} m_A^2 c_\beta^2 + M_Z^2 s_\beta^2 + \frac{8}{25}M_Z^2 t_\varphi^2(4s_\beta^2 - 1), & -\left(m_A^2 + M_Z^2\left[1 + \frac{16}{25}t_\varphi^2\right]\right)c_\beta s_\beta \\ -\left(m_A^2 + M_Z^2\left[1 + \frac{16}{25}t_\varphi^2\right]\right)c_\beta s_\beta, & m_A^2 s_\beta^2 + M_Z^2 c_\beta^2 + \frac{8}{25}M_Z^2 t_\varphi^2(4c_\beta^2 - 1) \end{pmatrix}, \quad (9)$$

where we parameterized the B_μ term as the pseudoscalar mass m_A^2, t_φ^2 is defined in Eq. (8), and $s_\beta(c_\beta)$ is $\sin\beta(\cos\beta)$.

Equation (9) leads to the following eigenvalues for the lighter and the heavier CP-even Higgs fields, for $m_A > M_Z$

$$m^2_{h,H} = \frac{1}{2}\left(m_A^2 + M_Z^2 + \frac{16}{25}t_\varphi^2\right) \mp \frac{1}{2}\left\{\left[m_A^2 + M_Z^2 + \frac{16}{25}t_\varphi^2\right]^2 - 4\cos^2 2\beta\left[m_A^2M_Z^2 - \frac{8}{25}M_Z^2t_\varphi^2\left(-3m_A^2 + M_Z^2 + \frac{24}{25}M_Z^2t_\varphi^2\right)\right]\right\}^{1/2}. \quad (10)$$

The condition to obtain positive m_h^2 is subtle because the term in the second line of Eq. (10) cannot be too large. It severely depends on the CP-odd Higgs mass m_A. The dependence of m_h on $\tan\varphi$ is depicted in Fig. 1(b) for a few values of m_A and $\tan\beta$. The $\tan\beta$ dependence converges in the large $\tan\beta$ region.

Comment related to F-theory: The above SU(6) $\times$ SU(2)$_h$ model can be obtained from the F-theory construction [22,23]. In this construction, we first obtain a visible six-dimensional (6D) GUT group, which is then broken to the four-dimensional (4D) SM group by fluxes. To obtain the 6D GUT group, one should consider the holonomy groups, the continuous and the discrete ones. The SU(3)$_\perp$ holonomy alone does not specify the 6D group completely. The additional information on the discrete holonomy is needed to choose one of the rank 6 groups among E$_6$, SU(6) $\times$ SU(2), and SU(3)$_3$. To guarantee the proton longevity, forbidding the dimension-5 operators, the hexality

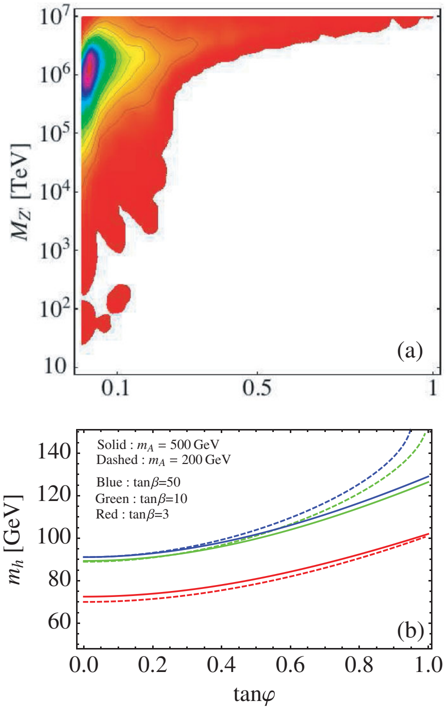

FIG. 1 (color). Masses of (a) Z' and (b) the lightest CP-even Higgs m_h as functions of $\tan\varphi$. The total number of data points is 14 235, and the region with no data points is white. Starting from the outmost contour (gray), the colors are separated by the density of points, in increments of ten for each step.

has been proposed. For a natural hexality, a visible sector should have $\mathbb{Z}_6$ [11,12,24]. The centers of SU(N) and E_6 are $\mathbb{Z}_N$ and $\mathbb{Z}_3$, respectively. So, we rule out 6D E_6 but require the 6D SU(6) part. Therefore, to introduce $\mathbb{Z}_6$ holonomy of SU(6) × SU(2), we look for $\mathbb{Z}_6$ holonomy of $SU(3)_\perp$. Since the center of $SU(3)_\perp$ is $\mathbb{Z}_3$, we need additional $\mathbb{Z}_2$, which is possible if the holonomy of the instanton is from the Belavin *et al.* type instanton [25]. Then, it is possible to have a 6D GUT group SU(6) × SU(2).

Conclusions: We analyzed the electroweak scale Z' in the context of a supersymmetric U(6) × SU(2)$_h$ grand unification model, which provides a $\mathbb{Z}_6$ for the hexality to make it safe from the dangerous proton decay. Motivated from the recent CDF result on the Wjj excess around 150 GeV, we analyzed the possibility of constructing a leptophobic Z' from our model. However, we briefly showed that there cannot exist any baryonic $U(1)'^B$ in any subgroup of E_6. Aside from $U(1)'_B$, the leptophobic Z' model from a supersymmetric E_6 is usually constructed through the kinetic mixing historically. Such mixing demands a large mixing coefficient $1/3$ which can arise from fine-tuned relations between the masses of the split multiplet members. It is also not achieved if one requires an anomaly-free model where the gauge couplings of the SM and $U(1)'$ are perturbative and unify at the GUT scale according to other research. Such a scenario is also constrained by the neutrino mixing.

Analyzing the electroweak ρ_0 parameter, the mass of our Z' is favored to be above 10 TeV by considering the neutrino mixing and proton decay constraint in supersymmetric models. In this sense, the CDF Wjj anomaly cannot be fitted to any electroweak model descending from E_6. We also discussed the $U(1)'$ effect on the tree-level mass of the lightest MSSM Higgs boson and the F-theory construction to obtain our supersymmetric SU(6) × SU(2)$_h$ model.

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