Some Advantages of SUSY SU(4) × SU(2)\textsubscript{L} × SU(2)\textsubscript{R} Model in String Derived SO(10) GUTs

Akihiro Murayama

Department of Physics, Faculty of Education, Shizuoka University
836 Ohya, Shizuoka 422, Japan

abstract

A D-parity violated SUSY SU(4) × SU(2)\textsubscript{L} × SU(2)\textsubscript{R} gauge model with the Higgs sector 2\{((4, 1, 2) + (\overline{4}, 1, 2)) + (1, 2, 2) + some (1, 1, 1)'s\} is shown to have the following advantages: (i) It is the simplest and almost unique solution that satisfies $M_X = M_{\text{string}} \approx 0.6 \times 10^{18}$ GeV and $M_{\text{INT}} \approx 5 \times 10^{11}$ GeV in superstring derived SUSY SO(10) GUTs. (ii) The proton is stable enough by the automatic ”doublet-triplet splitting” closely connected with the D-parity violation. (iii) The minimization of SUSY one-loop effective potential in a toy model suggests that the SO(10) gauge theory tends to break dynamically down to the SU(4) × SU(2)\textsubscript{L} × SU(2)\textsubscript{R} model.

* E-mail: murayama@ed.shizuoka.ac.jp
Supersymmetric (SUSY) SO(10) grand unified theory (GUT) is one of candidates of true GUT because:

(1) The group SO(10) contains the standard model (SM) gauge symmetry $SU(3)_C \times SU(2)_L \times U(1)_Y$.

(2) All members of each generation of light quarks and leptons belong to single irreducible representation $16$ of SO(10) together with a right-handed neutrino and the anomaly cancellation is automatic.

(3) It naturally provides us with an intermediate mass scale, $M_{INT}$, which might develop a new physics. The case of $M_{INT} \approx 10^{11-12}$ GeV is of special interest in connection with an invisible axion [1], massive neutrinos through the see-saw mechanism [2], an inflaton for generating the cosmological baryon asymmetry [3] and so on.

(4) It is a subgroup of $E_8$ and so the SO(10) GUT could be derived from the compactification of superstring theory (SST) such as $E_8 \times E'_8$ heterotic string theory.

(5) The grand unification scale $M_X$ can be raised from $O(10^{16}$ GeV) of minimal SUSY standard model (MSSM) to the string scale $M_{string} \sim 10^{18}$ GeV.

Suppose that the SUSY SO(10) GUT has been somehow derived from an underlying SST, e.g., the $E_8 \times E'_8$ heterotic string, by a compactification à la Witten [4] on some manifold $K$, and that the situation corresponds to level 1 Kač-Moody algebra representations. Then, the contents of the chiral scalar superfields at the GUT scale $M_X$ are given by [4,5]

$$n_g 16 + \delta (16 + \overline{16}) + \varepsilon 10 + \eta 1,$$

(1)

where $n_g (= 3 )$ is the number of generations of light quarks and leptons and $\delta, \varepsilon$ and $\eta$ denote the number of Higgs superfields of representations $16 + \overline{16}, 10$ and $1$, respectively. The values of $n_g, \delta, \varepsilon$ and $\eta$ depend, in principle, on the topology of $K$.

In passing, it should be remarked that the gauge group SO(10) cannot be obtained by the simple standard embedding of SU(3) holonomy of $K$ into $E_8$ but by an embedding of gauge connection of holomorphic SU(4) vector bundle over $K$ [4]. Consequently, the number of generations of quarks and leptons is not given by one half of Euler characteristics of $K$ but by more general Atiyah-Singer index theorem [4,6] which could develop a new potential of explaining $n_g = 3$.

There are several possible paths from the SUSY SO(10) GUT to MSSM. For simplicity, we will confine here ourselves to the cases of single intermediate scale, namely to the breaking pattern:

$$\text{SO}(10) \xrightarrow{M_X} G_{INT} \xrightarrow{M_{INT}} SU(3)_C \times SU(2)_L \times U(1)_Y,$$

(2)
and investigate the following cases of intermediate symmetries:

\[
G_{\text{INT}} = \begin{cases} 
(a) & \text{SU}(4) \times \text{SU}(2)_L \times \text{SU}(2)_R \times D, \\
(b) & \text{SU}(4) \times \text{SU}(2)_L \times \text{SU}(2)_R, \\
(c) & \text{SU}(4) \times \text{SU}(2)_L \times U(1)_R, \\
(d) & \text{SU}(3)_C \times \text{SU}(2)_L \times \text{SU}(2)_R \times U(1)_{B-L} \times D, \\
(e) & \text{SU}(3)_C \times \text{SU}(2)_L \times \text{SU}(2)_R \times U(1)_{B-L},
\end{cases}
\]

where \(D\) in (a) and (d) means that the spectrum of Higgs sector is symmetric under the exchange of \(L \leftrightarrow R\), i.e., \(D\)-parity \([7]\) is conserved, while in (b) and (e) the \(D\)-parity is violated. We discard the case \(G_{\text{INT}} = \text{SU}(5)\) or \(\text{SU}(5) \times U(1)\) because it requires rather big \(M_{\text{INT}}\) which does not deserve the name of "intermediate".

The purpose of this paper is to show that the path (b) with the absence of the component \((4, 2, 1) + (\bar{4}, 2, 1)\) in the Higgs multiplet \(16 + \overline{16}\) has exclusively remarkable virtues. We call the model with this breaking path a \(D\)-parity violated SUSY Pati-Salam model \([8]\). In the following, we present characteristic features of the model in three different aspects: (I) the renormalization group equation (RGE) analysis, (II) the "doublet-triplet splitting" by Witten mechanism and (III) the minimization of SUSY one-loop effective potential in a toy model.

(I) \textit{RGE analysis}

In this section, we analyze the evolution of gauge couplings for (2). In all the cases (a)\,\sim\,(e), we assume that the colored components in the Higgs multiplet \(10\) are at least as heavy as \(M_X\) and decoupled from the theory below the unification scale, so that the "doublet-triplet splitting" has been realized to avoid the danger of fast proton decay. It should be noticed here that the colored Higgses belong to the different representation of \(G_{\text{INT}}\) from the electroweak doublets. Therefore, the "doublet-triplet splitting" could be more easily implemented for (a)\,\sim\,(e) than for SU(5). This is another reason why we do not take SU(5) or SU(5) \(\times U(1)\) as candidates of \(G_{\text{INT}}\).

We find \([9-11]\) that in the \(D\)-parity violated SUSY Pati-Salam model, a simple choice of the Higgs sector, \(2\{(4, 1, 2) + (\bar{4}, 1, 2)\} + (1, 2, 2) + \text{some } (1, 1, 1)'\)s which corresponds to \(\delta = 2\) and \(\varepsilon = 1\) in (1), can attain \(M_X = M_{\text{string}} \approx 0.6 \times 10^{18}\) GeV and, at the same time, \(M_{\text{INT}} \approx 5 \times 10^{11}\) GeV \([9-11]\). The relation \(M_X = M_{\text{string}}\) indicates that the intermediate gauge symmetry \(G_{\text{INT}}\) might be directly realized from the string at \(M_{\text{string}}\) with the gauge couplings unified so that the "GUT" group SO(10) need not in fact be embodied explicitly. For such cases we will use the word "SO(10) GUT" in the sense that \(G_{\text{INT}} \subset \text{SO}(10)\).
This case of $G_{INT} = SU(4) \times SU(2)_L \times SU(2)_R$ with the Higgs sector $2\{(4, 1, 2) + (\bar{4}, 1, 2) \} + (1, 2, 2) + \text{some } (1, 1, 1)'s$ is actually the only possibility that both of the constraints (i) $M_X = M_{\text{string}}$ and (ii) $M_X / M_{INT} \approx 10^6$ are fulfilled in the SUSY SO(10) GUT with the single intermediate scale. Indeed, according to the RGE analysis of ref.[12], if we demand the constraints (i) and (ii), the relative changes in the beta functions of the MSSM above $M_{INT}$ due to the additional Higgs supermultiplets must satisfy the condition

$$\frac{2}{5}r \equiv \Delta b_2 - \Delta b_1 = 2, \quad q \equiv \Delta b_3 - \Delta b_2 = 1, \quad (3)$$

where $b_i = -2\pi \partial \alpha_i^{-1} / \partial \ln \mu$. The above case of $D$-parity violated SUSY Pati-Salam model just satisfies (3). In the other cases, we have

(a) $q + r = 18$, \quad (4)
(b) $r < 0$, \quad (5)
(c) $q + r = 0, 9$, \quad (6)
(d) $q + r = 9, 3$ or $q + r < 0$, \quad (7)

even with the help of exotics. Obviously, (4)$\sim$(7) cannot be compatible with the condition (3).

It is noteworthy that, although the cases (a) and (b) are both the SUSY Pati-Salam model, only the case (b) can satisfy (3), namely the $D$-parity violation is indispensable for the SUSY Pati-Salam model to realize $M_X = M_{\text{string}}$ and $M_{INT} \approx 5 \times 10^{11}$ GeV. In this context, it might not be accidental that, in a simple SST-derived SUSY Pati-Salam model [13,10] constructed by the fermionic formulation [14], the Higgs sector contains two copies of $(4, 1, 2) + (\bar{4}, 1, 2)$, but no $(4, 2, 1) + (\bar{4}, 2, 1)$, so that the $D$-parity is violated. In order to make the model more realistic, we must find a dynamical mechanism in which the gauge symmetry breaking $SU(4) \times SU(2)_L \times SU(2)_R \to SU(3)_C \times SU(2)_L \times U(1)_Y$ takes place actually at $M_{INT} \approx 5 \times 10^{11}$ GeV. An example of such a mechanism is found in ref.[11].

(II) "Doublet-triplet splitting"

Hereafter we assume that the compact manifold $K$ is multiply-connected. Then, all fields obey nontrivial boundary conditions and the gauge fields can develop vacuum expectation values (VEVs) on $K$. The gauge symmetry breaking is dynamically caused by Hosotani mechanism [15,16] through the modified Wilson loops $W$s[16]. Let $G$ be $\pi_1(K)$ and $\hat{G}$ the image of a homomorphim of $G$ into the gauge group SO(10). Then,
\( \hat{G} \) will be the symmetry of \( W \) and it is only \((G \oplus \hat{G})\)-invariant modes that remain massless after the breaking (Witten mechanism \([17]\)).

In the case of the path \((b)\), \( G_{\text{INT}} = \text{SU}(4) \times \text{SU}(2)_L \times \text{SU}(2)_R \) with the \( D \)-parity violated, the "doublet-triplet splitting" (in this case, the decoupling of the component \((6,1,1)\) in the Higgs multiplet \( 10 \)) by Witten mechanism is closely connected with the \( D \)-parity violation \([9]\). Namely, it is automatic in the sense that any two of the following three statements lead to the remaining one (for the proof, see ref.\([9]\)):

(A) The \( \text{SO}(10) \) breaks down to \( \text{SU}(4) \times \text{SU}(2)_L \times \text{SU}(2)_R \), i.e., the component \((6,2,2)\) in the gauge multiplet \( 45 \) becomes superheavy while the components \((15,1,1)\), \((1,3,1)\) and \((1,1,3)\) remain massless.

(B) The \( D \)-parity is violated, i.e., the component \((4,2,1)+(\bar{4},2,1)\) in the Higgs multiplet \( 16 + \bar{16} \) becomes superheavy while the component \((4,1,2)+(\bar{4},1,2)\) remains massless, or \textit{vice versa}.

(C) The "doublet-triplet splitting" is realized, i.e., the component \((6,1,1)\) in the Higgs multiplet \( 10 \) becomes superheavy while the component \((1,2,2)\) remains massless.

From (I) and (II), we understand that, in the \( D \)-parity violated SUSY Pati-Salam model, the achievement of \( M_X = M_{\text{string}} \approx 0.5 \times 10^{18} \) GeV and the appearance of \( M_{\text{INT}} \approx 5 \times 10^{11} \) GeV are closely connected with the stability of proton in terms of the \( D \)-parity violation. As to the potentially dangerous Higgs multiplets \((4,2,1)+(\bar{4},2,1)\)'s, which contain color triplet component \((\bar{3},1)_{2/3}+(3,1)_{-2/3}\) of \( \text{SU}(3)_C \times \text{U}(1)_Y \), their direct couplings with the quarks and leptons are forbidden by the symmetry of the model. This is due to the fact that \( 16 \times 16 \times 16 \) or \( 16 \times 16 \times \bar{16} \) do not contain \( 1 \) of \( \text{SO}(10) \) and this feature is inherited by \( \text{SU}(4) \times \text{SU}(2)_L \times \text{SU}(2)_R \). After the breaking of \( \text{SU}(4) \) to \( \text{SU}(3) \) at \( M_{\text{INT}} \), they effectively couple with the quarks and leptons through the process such as

\[
16 + 16 \rightarrow 10 \rightarrow 16^+ < 16 >,
\]

where \( < > \) indicates a VEV. However, the strength of induced effective coupling \( 16 \cdot 16 \cdot 16 \) is greatly suppressed by a factor \( M_{\text{INT}}/M_X \) and negligible.

(III) \textit{SUSY effective potential}

There is a hint that the path of SUSY Pati-Salam model is dynamically favorable. It is the embedding scheme of \( \hat{G} \) into \( \text{SO}(10) \) that characterizes the breaking direction \([17,18]\). However, we cannot arbitrarily choose the Wilson loop but should determine \( \hat{G} \) through the minimization of the effective potential for the SUSY \( \text{SO}(10) \) gauge theory. The minimization of SUSY one-loop effective potential in a toy model
suggests that the breaking direction $SO(10) \rightarrow SU(4) \times SU(2)_L \times SU(2)_R$ or $SU(4) \times SU(2)_L \times SU(2)_R \times D (\times D$ is henceforth abbreviated) is more probable than other directions. Indeed, in an $SO(10)$ gauge theory on $M_3 \times S^1$, where $M_3$ is the three-dimensional Minkowski spacetime, the SUSY one-loop effective potential $V_1$ can be estimated by the auxiliary field tadpole method [19]. The following derivation of $V_1$ ((16) below) is a revised version of ref.[20].

We assume that (i) the gauge fields have non-vanishing VEV $< A_y > \neq 0$ on $S^1$, $y$ indicating the coordinate of $S^1$, (ii) the VEVs of the scalar components of matter superfields are zero, (iii) the effective potential satisfies the SUSY boundary condition, $V_1(f = f^+ = d = 0) = 0$ where $f \equiv < F >$, $f^+ \equiv < F^+ >$ and $d \equiv < D >$ denote the VEVs of auxiliary fields and (iv) the gauge symmetry does not break spontaneously, which means that $f$ is gauge singlet and $d = 0$ so that we need not consider the $D$ tadpoles.

The action is given by [21]

$$S = \int dz \left[ \sum_i \phi_i^+ e^{gV} \phi_i \right] + \left[ \int d\sigma \left( \frac{1}{4} tr(W^a W_a) + P(\phi_i) \right) + h.c. \right],$$

where $\phi_i$ stands for chiral matter superfields in the irreducible representation of $SO(10)$ displayed in (1), $V \equiv V^a T^a$ defines the vector superfield with $T^a$ being the generators of $SO(10)$ in the adjoint representation [22], $W_a \equiv -\frac{1}{4} T^a \exp(-gV) D_a \exp(gV)$ denotes the SUSY field strength, $dz = d^3y dy d^2\theta d\theta \bar{\theta}$, $d\sigma = d^3y dy d^2\theta d\theta \bar{\theta}$ and $P(\phi_i) = \frac{1}{2} m_{ij} \phi_i \phi_j + \frac{1}{3!} \lambda_{ijk} \phi_i \phi_j \phi_k + \cdots$ is the superpotential.

By making a component field expansion for (8) in the Wess-Zumino gauge and by translating the fields as $F_i \rightarrow F_i + f_i$ and $A_\mu \rightarrow A_\mu + a \delta_\mu y$, where $A_\mu$ denotes the spin-1 component of $V$ and $a \equiv < A_y >$, we obtain propagators for spin-0 component $\Phi_i$ of $\phi_i$ which are relevant to the $F$ tadpole as follows:

$$\Phi_i \Phi_j : -\frac{\lambda_{ijk} f_k^+}{\Delta},$$

where

$$\Delta(a, f) = \det\{(p_\mu - ga \delta_{\mu y})^4 - \lambda_{ijk} \lambda_{ijk'} f_k^+ f_{k'} \}. \quad (10)$$

In terms of irreducible representations of $SO(10)$, $\Phi_i \Phi_j$ and $f_k^+ = f_{k'} \equiv f$ in (9) and (10) in fact belong respectively to $16 \times \overline{16}$ and $1$ due to the assumption (iv).

Then, using the SUSY boundary condition $V_1(f = 0) = 0$, we obtain

$$V_1 = \frac{1}{2} \int \frac{d^4p}{(2\pi)^4} [\ln \Delta(a, f) - \ln \Delta(a, 0)]$$
\[
\frac{1}{2} \text{tr} \int \frac{d^4 p}{(2\pi)^4} \left[ \ln \{(p_\mu - g a \delta_\mu y)^4 - \chi^2 f^2\} - \ln \{(p_\mu - g a \delta_\mu y)^4\}\right] \\
= \frac{\chi^2 f^2}{4\pi^2 L} \text{tr} \sum_{n=\infty}^{\infty} \int_0^\infty p^2 dp \ln[p^2 + (\omega_n - g a)^2] + O(f^3),
\]

where \(\omega_n = (2\pi n + \beta)/L\), with \(\beta\) being the phase which can enter the boundary conditions on \(S^1\) for the chiral superfields \(\phi_i\)'s and \(L\) being the periodicity of the coordinate \(y\) of \(S^1\). In arriving at the third line of (11), we have assumed \(|f| \ll L^{-1}\), that is, the SUSY breaking scale is much smaller than the compactification scale. Under the constraint that \(G_{INT}\) has the same rank with SO(10) and that SU(3)_C \times SU(2)_L is unbroken, the 16 \times 16 matrix \(a\) can be parametrized in terms of two real free parameters as [16, 18]

\[
a = (gL)^{-1}(\theta H^\theta + \psi H^\psi),
\]

where

\[
\begin{align*}
\theta H^\theta &= \text{diag}(\theta_1, \cdots, \theta_{16}) \\
&= \theta \text{diag}(-3, -3, 1, 1, 1, 1, 1, 1, 3, 3, -1, -1, -1, -1, -1, -1, -1), \\
\psi H^\psi &= \text{diag}(\psi_1, \cdots, \psi_{16}) \\
&= \psi \text{diag}(0, 0, 0, 0, 0, 0, 0, 1, -1, 1, 1, -1, -1, -1, -1).
\end{align*}
\]

The parametrization (12) corresponds to the Wilson loop

\[
W = \exp(\theta H^\theta + \psi H^\psi).
\]

The \(\theta\) and \(\psi\) dependence of \(V_1\) can easily be estimated by the method of ref.[23] to be

\[
V_1(\theta, \psi) = \frac{\chi^2 f^2}{\pi^2 L^4} \sum_{j=1}^{16} \sum_{n=-\infty}^{\infty} \frac{\cos n (\theta_j + \psi_j - \beta)}{n^4} + (\theta, \psi - \text{independent}),
\]

where \(O(f^3)\)–term has been neglected. It is evident that, if one takes \(\beta = 0\), which means that 16- or 16-dimensional chiral fields obey a periodic boundary condition, for simplicity, the potential \(V_1\) attains the minimum at \(\theta = \pi\) and \(\psi = 0\) (mod2\(\pi\)). Unfortunately, the corresponding Wilson loop \(W\) (15) can neither break SO(10) [18] nor determine \(\hat{G}\). Therefore, the minimization of \(V_1\) cannot discriminate \(G_{INT}\)'s.
However, it will be reasonable to assume $\hat{G} = Z_m \times Z_n$ and take an average $\bar{V}_1(m, n)$ of $V_1(\theta, \psi)$ such as

$$\bar{V}_1(m, n) = \frac{1}{mn-1} \sum_{k=0}^{m-1} \sum_{l=0}^{n-1} \{V_1(\frac{2k\pi}{m}, \frac{2l\pi}{n}) - V_1(0, 0)\}, \quad (17)$$

and minimize it in terms of $m, n$ except for a trivial case $m = 2, n = 1$. The result is remarkable. The $\bar{V}_1$ attains the minimum when $\hat{G} = Z_4$ ($m = 4, n = 1$), which gives $\text{SO}(10) \rightarrow G_{\text{INT}} = \text{SU}(4) \times \text{SU}(2)_L \times \text{SU}(2)_R$. In fact, we obtain

$$\bar{V}_1(4, 1) < \bar{V}_1(2, 2) < \bar{V}_1(5, 1) < \bar{V}_1(6, 1) < \cdots, \quad (18)$$

which implies that, in the breaking of $\text{SO}(10)$ to $G_{\text{INT}}$ via Wilson loops, $G_{\text{INT}} = \text{SU}(4) \times \text{SU}(2)_L \times \text{SU}(2)_R$ by $\hat{G} = Z_4$ or $\hat{G} = Z_2 \times Z_2$ \cite{18} is more probable than $G_{\text{INT}} = \text{SU}(3)_C \times \text{SU}(2)_L \times \text{SU}(2)_R \times U(1)_{B-L}$ as well as $G_{\text{INT}} = \text{SU}(4) \times \text{SU}(2)_L \times U(1)_R$ by $\hat{G} = Z_3, \hat{G} = Z_m, m \geq 5$ or other $\hat{G}$’s.

It is obvious that the above inference is valid also for $M_4 \times T^6$, where $M_4$ and $T^6$ are four-dimensional Minkowski spacetime and six-torus, respectively, which is derived from $E_8 \times E'_8$ heterotic string by torus compactification, provided that the VEVs of gauge fields are isotropic on $T^6$ in the sense that their 6 components are all identical independent of the way of their corresponding to $S^1$’s in $T^6$. It will not be so trivial to generalize the above toy model to more realistic cases. However, the characteristic feature of $V_1$ (16) is expected to hold as far as the manifold $K$ is multiply-connected and the VEV of the gauge field on $K$ can be parametrized as (12). It is one of next tasks to investigate if such argument is also applicable to orbifolds or more general multiply-connected compact manifolds of dimension six than $T^6$.

In summary, we have seen that the $D$-parity violated SUSY Pati-Salam model with the Higgs sector consisting of $2\{(4, 1, 2) + (\overline{4}, 1, 2)\} + \{1, 2, 2\} + \text{some } (1, 1, 1)$’s has the following advantages: (i) It is the simplest and almost unique solution that satisfies $M_X = M_{\text{string}} \approx 0.6 \times 10^{18}$ GeV and $M_{\text{INT}} \approx 5 \times 10^{11}$ GeV in SST-derived SUSY SO(10) GUTs. (ii) The stability of proton is guaranteed by the automatic “doublet-triplet splitting” and closely connected with the $D$-parity violation which is indispensable for the advantage (i) to be realized. (iii) The minimization of SUSY one-loop effective potential in a toy model suggests that the SUSY SO(10) gauge theory is expected to break dynamically down to the SUSY Pati-Salam model.
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