J. FAC. SCI., SHINSHU UNIVERSITY, Vol. 28, No. 1, pp.27-39 1993

A Variant Scenario in String Unification

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(Received June 8, 1993)

Abstract.

We report a new solution to the problem related to the gauge coupling unification in superstring theories. Our solution is based on a dynamical assumption and it is applied to 4-dimensional string models. It is shown that these models have phenomenologically interesting features (three families, $m_{3/2} \simeq O(1)$ TeV, the universal soft SUSY breaking terms,$\cdots$).

1 Introduction

The minimal supersymmetric standard model (MSSM) is the most attractive candidates for the realistic theory beyond standard model. The hierarchy problem is elegantly solved by the introduction of supersymmetry (SUSY)[1]. Furthermore the recent precision measurements at LEP [2] have given strong support to the supersymmetric grand unified theories (SUSY GUTs) [3], that is, if the renormalization group equations of MSSM are used, the three gauge coupling constants, g_3, g_2 and g_1 of $G_{ST}=SU(3)\times SU(2)\times U(1)$ meet at about 10^{16} GeV[4]. If superpartners of usual particles are found and the values of their masses are $O(1)$ TeV, MSSM is most likely established as the physics below the grand unification scale M_{GUT}. However, it is difficult to regard SUSY GUTs as the final theory because there are still several open questions not to answer within the framework of SUSY GUTs. First, SUSY GUTs do not include gravity, while the Planck scale M_{Pl} is around the corner. Second, there exists a great deal of arbitrariness on the model building, that is, a lot of freedom is left over on the choice of gauge group, matter multiplets and parameters. They do not explain the number of families (generations). Last, it is not yet known how to break SUSY and to get the desired low energy physics.

Superstring theories (SSTs)[5] are powerful candidates as the fundamental theory of nature and are expected to give definite answers for the above questions. For example, they probably describe quantum gravity in a consistent manner and some of them include a grand unified gauge group such as $SU(5)$, $SO(10)$and E_6,chiral matters and a hidden sector with sever constraints. The family number is supposed to be related to a sort of topological number in the extra compactified space. Moreover

SSTs are probably well described as effective $N=1$ supergravity coupled to supersymmetric Yang-Mills theory below the string scale $M_S=2/\sqrt{\alpha'}$ and it is known that SUSY can be softly broken in an observable sector by some scalar condensations [6] or gaugino condensation [7] in a hidden sector in supergravity theories.

Now it is natural that one tries to construct the unified models, which take over the advantages in SUSY GUTs, based on SSTs. In this attempt, one encounters a serious difficulty related to the unification of gauge coupling constants. In SSTs, the unification scale of all gauge coupling constants is believed to be M_S, that is, somewhat larger one $\sim 10^{18}$ GeV [8]. This fact apparently disagrees with the precise measurement on the Weinberg angle $sin^2\theta_w$ at LEP.

There have been, so far, proposed two general solutions to this problem. First, the gauge group at the string scale may be broken down to the GUT group such as SU (5) or SO (10) once. In this case, one needs Higgs scalars in the adjoint representations of SU (5) or SO (10) to break further the GUT groups down to the standard gauge group G_{ST} at M_{GUT}. It is possible to have adjoint Higgs scalars in string models if one uses the higher levels of Kac-Moody algebra ($k \geq 2$). However, realistic models with gauge groups at level greater than one have not been found yet [9].

A second possibility is to include string-loop threshold corrections [10] or to add extra matter multiplets [11] at the intermediate scale in order to shift the unification scale of three gauge coupling constants from 10^{16} GeV to 10^{18} GeV. Such a large shift may be possible, since an infinite number of massive states above the string scale can contribute to the threshold. However, one should consider in this case that the beautiful success of the minimal SUSY GUT is even accidental.

Recently, we have proposed a new solution [12]. Our proposal is founded on a new assumption that the GUT groups are broken down to G_{ST} dynamically by some effective, non-renormalizable interactions of the fundamental supermultiplets. The effective interactions may be remnant of compactification of extra space. As an example, we have applied our hypothesis to a simple 4-dimensional string model derived from the Z_7 orbifold compactification.

In this paper, we report previous results in detail with some additional ones. We introduce our hypothesis on the dynamical symmetry breakings and apply it to 4-dimensional string models. The features of these models are elucidated. As new subjects, we discuss the mass spectra after the dynamical symmetry breakings and the structure of soft SUSY breaking terms in our models. (Only the universality of scalar masses is commented on in our previous paper [12].)

The content of this paper is as follows. In section 2, we review the energy scales and parameters in SSTs. In section 3, we explain our dynamical symmetry breaking scenario in SSTs and apply it to two string models. Summary and discussions are given in section 4.

2 Parameters in SSTs

Let us first give a brief review on several energy scales and parameters in SSTs [5] [13]. SSTs have a distinctive feature that they include only one fundamental parameter α' which is called Regge slope parameter. This parameter is related to the string tension T as $T=(2\pi\alpha')^{-1}$. The string scale is defined as

$$M_S = \frac{2}{\sqrt{\alpha'}} \tag{1}$$

Other energy scales and parameters are expected to be generated dynamically.*

As is described in introduction, all gauge couplings are unified at the string tree level as [8],

$$\frac{4\pi}{\alpha'} G_N = k_i \cdot g_i (M_S)^2 \tag{2}$$

where G_N is the gravitational constant which is related to the Planck scale M_{Pl} as $M_{Pl} = 1/\sqrt{G_N}$ and k_i's are the Kac-Moody levels of gauge group whose gauge coupling constants (GCCs) are g_i. We consider only the level one Kac-Moody algebra for non-abelian gauge groups and hence $\frac{3}{5}k_1 = k_2 = k_3 = 1$.† The gauge coupling constants g_i at M_S and the size R of the extra compactified space are related to the vacuum expectation values (VEVs) of real parts of a dilaton field S and a moduli field T such as

$$< ReS > = \frac{1}{k_i \cdot g_i(M_S)^2} \tag{3}$$

and

$$< ReT > = R^2 \tag{4}$$

respectively. The effective potential that fixes $<ReS>$ and $<ReT>$ is not known. From eqs. (1), (2) and (3), the scale M_S is related also to the gravitational scale $M \equiv M_{Pl}/\sqrt{8\pi} = 2.4 \times 10^{18}$ GeV as

$$M_s = \sqrt{\frac{2}{< ReS >}} M \tag{5}$$

* Unfortunately, no parameter has been determined because the dynamics of SSTs are not fully understood.

†For $U(1)$ gauge group, the level k is not quantized. Here we choose $k_1 = \frac{5}{3}$ because it is consistent with the gauge coupling unification condition in GUTs at M_s.

‡The values of them are given in the unit of string scale M_s.

Since the quantity $<ReS>$ is anticipated to be order 1,‡ M_s is estimated as $\sim 10^{18}$ GeV. Here and hereafter a field variable represents its vacuum expectation value without a bracket. The VEV of some auxiliary field F is an order parameter of SUSY breaking and the breaking scale M_{SS} is defined as $M_{SS}{}^2 \equiv F$. In the case that the SUSY is broken by the gaugino condensation $<\lambda\lambda>$, the scale M_{SS} is given by [14],

$$M_{SS}{}^2 = \frac{|<\lambda\lambda>|}{M} \qquad (6)$$

The quantity $|<\lambda\lambda>|$ is estimated at order of Λ_C^3. Here Λ_C is the scale where the GCC blows up in the hidden gauge theory whose gaugino is λ. The mass of the gravitino is given by [14],

$$m_{3/2} = \frac{|<\lambda\lambda>|}{M^2} \sim \frac{\Lambda_C^3}{M^2} \qquad (7)$$

The masses of superpartner of usual particles (quarks, leptons and gauge bosons) are the same order of $m_{3/2}$, so the solution to hierarchy problem requires that $m_{3/2} \simeq O(1)$ TeV. Hence the favorite value of Λ_C is $\sim 10^{13}$ GeV.

3 Dynamical Breaking Scenario

We propose a new assumption that the GUT groups are broken down to G_{ST} dynamically by some effective, non-renormalizable interactions of the fundamental supermultiplets. Namely, the Higgs multiplets in the adjoint representations are bound states of the fundamental matter multiplets. The effective interactions may be remnant of compactification of extra space. Although it is not clear to us whether the dynamical breaking occurs or not in the string models, this working hypothesis opens a new window in the superstring phenomenology. In fact, the GUT with an exceptional group E_6 is a well-known example where all the symmetry breakings required phenomenologically are obtained from the fundamental fermion-fermion condensations [15]. As an example, we shall apply our hypothesis to two interesting 4-dimensional string models derived from the Z_7 orbifold compactification [16].

3.1 Two Z_7 Orbifold Models

We explain two 4-dimensional string models with gauge group $E_6 \times U(1)^2 \times E_6' \times U(1)'^2$ [17] obtained from the Z_7 orbifold compactification of the heterotic string with gauge group $E_8 \times E_8'$. (The prime (′) represents that they belong to the hidden sector. For a complete construction, see ref. [16].)

The first model is obtained by the choice of shift vectors V^I and v^t as follows,

$$V^I = (1, 2, -3, 0, \cdots, 0)/7(1, 2, -3, 0, \cdots, 0)/7 \qquad (8)$$

and

$$v^t = (0, 1, 2, -3)/7 \tag{9}$$

where $I = 1, 2, \cdots, 16$ and $t = 1, 2, 3, 4$. The shift V^I breaks the original gauge group $E_8 \times E_8'$ down to $E_6 \times U(1)^2 \times E_6' \times U(1)'^2$. Massless matter representations of the untwisted sector are $3(\mathbf{27}, \mathbf{1}') + 3(\mathbf{1}, \mathbf{27}') + 6(\mathbf{1}, \mathbf{1}')$ under $E_6 \times E_6'$. The ground state in the twisted sector consists of 147 singlets. When we choose simple roots of $E_6 \times E_6'$ as

$$\begin{aligned}
e^1 &= (0, 0, 0, 0, 0, 0, -1, 1)(0, \cdots, 0)\\
e^2 &= (0, 0, 0, 0, 0, -1, 1, 0)(0, \cdots, 0)\\
e^3 &= (0, 0, 0, 0, -1, 1, 0, 0)(0, \cdots, 0)\\
e^4 &= (0, 0, 0, -1, 1, 0, 0, 0)(0, \cdots, 0)\\
e^5 &= (1, 1, 1, 1, -1, -1, -1, -1)/2(0, \cdots, 0)\\
e^6 &= (0, 0, 0, 1, 1, 0, 0, 0)(0, \cdots, 0)\\
e'^1 &= (0, \cdots, 0)(0, 0, 0, 0, 0, 0, -1, 1)\\
e'^2 &= (0, \cdots, 0)(0, 0, 0, 0, 0, -1, 1, 0)\\
e'^3 &= (0, \cdots, 0)(0, 0, 0, 0, -1, 1, 0, 0)\\
e'^4 &= (0, \cdots, 0)(0, 0, 0, -1, 1, 0, 0, 0)\\
e'^5 &= (0, \cdots, 0)(1, 1, 1, 1, -1, -1, -1, -1)/2\\
e'^6 &= (0, \cdots, 0)(0, 0, 0, 1, 1, 0, 0, 0)
\end{aligned}$$

and $U(1)$ charges as

$$\begin{aligned}
U_1 &= (1, -1, 0, 0, 0, 0, 0, 0)(0, \cdots, 0)\\
U_2 &= (1, 1, -2, 0, 0, 0, 0, 0)(0, \cdots, 0)\\
U_1' &= (0, \cdots, 0)(1, -1, 0, 0, 0, 0, 0, 0)\\
U_2' &= (0, \cdots, 0)(1, 1, -2, 0, 0, 0, 0, 0)
\end{aligned}$$

the untwisted matters $3(\mathbf{27}, \mathbf{1}')$ and $3(\mathbf{1}, \mathbf{27}')$ have the following $U(1)$ charges

$$(\mathbf{27}, \mathbf{1}'\,;\ 1, 1, 0, 0) \quad \text{for } \textstyle\sum_{I=1}^{16} P^I V^I = \frac{1}{7} \tag{10}$$

$$(\mathbf{27}, \mathbf{1}'\,;\ -1, 1, 0, 0) \quad \text{for } \textstyle\sum_{I=1}^{16} P^I V^I = \frac{2}{7} \tag{11}$$

$$(\mathbf{27}, \mathbf{1}'\,;\ 0, -2, 0, 0) \quad \text{for } \textstyle\sum_{I=1}^{16} P^I V^I = \frac{4}{7} \tag{12}$$

and

$$(\mathbf{1}, \mathbf{27}'\,:\ 0, 0, 1, 1) \quad \text{for } \textstyle\sum_{I=1}^{16} P^I V^I = \frac{1}{7} \tag{13}$$

$$(\mathbf{1}, \mathbf{27}'\,;\ 0, 0, -1, 1) \quad \text{for } \textstyle\sum_{I=1}^{16} P^I V^I = \frac{2}{7} \tag{14}$$

$$(\mathbf{1}, \mathbf{27}'\,;\ 0, 0, 0, -2) \quad \text{for } \textstyle\sum_{I=1}^{16} P^I V^I = \frac{4}{7} \tag{15}$$

under $E_6 \times E_6' \times U(1)^2 \times U(1)'^2$. This model has three families. This is due to the fact that there exist three subsectors which correspond to three of six values of $\sum_{I=1}^{16} P^I V^I = \frac{1}{7}, \cdots, \frac{6}{7}$ in the untwisted sector where P^I's are quantized momenta which span the $E_8 \times E_8'$ root lattice.

The second model is obtained by taking shift vectors V^I, v^t and Wilson line a^I as follows,

$$V^I = (1, 2, -3, 0, \cdots, 0)/7(0, 0, 0, 0, \cdots, 0)/7, \tag{16}$$

$$v^t = (0, 1, 2, -3)/7 \tag{17}$$

and

$$a^I = (1, 2, -3, 0, \cdots, 0)/7(1, 2, -3, 0, \cdots, 0)/7. \tag{18}$$

The shift V^I breaks the original gauge group $E_8 \times E_8'$ down to $E_6 \times U(1)^2 \times E_8'$ and the Wilson line a^I breaks further it down to $E_6 \times U(1)^2 \times E_6' \times U(1)'^2$. There is no massless state in the untwisted sector because of the physical state condition $\sum_{I=1}^{16} P^I a^I \in Z$. Massless matter representations of the twisted sectors are $3(\mathbf{27}, \mathbf{1}') + 3(\mathbf{1}, \mathbf{27}')$ and some singlets under $E_6 \times E_6'$.If we choose the same simple roots and the same $U(1)$ charge assignments as those of the first model, the $U(1)$ charges of twisted matters 3 $(\mathbf{27}, \mathbf{1}')$ and $3(\mathbf{1}, \mathbf{27}')$ are as follows

$$(\mathbf{27}, \mathbf{1}'\,;\ 0, -2, 0, 0) \quad \text{for the first twisted sector} \tag{19}$$

$$(\mathbf{27}, \mathbf{1}'\,;\ 0, -2, 0, 0) \quad \text{for the second twisted sector} \tag{20}$$

$$(\mathbf{27}^*, \mathbf{1}'\,;\ 0, -4, 0, 0) \quad \text{for the third twisted sector} \tag{21}$$

and

$$(\mathbf{1}, \mathbf{27}',:\ 0, 0, 0, -2) \quad \text{for the first twisted sector} \tag{22}$$

$$(\mathbf{1}, \mathbf{27}'\,;\ 0, 0, 0, -2) \quad \text{for the second twisted sector} \tag{23}$$

$$(\mathbf{1}, \mathbf{27}^{*\prime}\,;\ 0, 0, 0, -4) \quad \text{for the third twisted sector} \tag{24}$$

under $E_6 \times E_6' \times U(1)^2 \times U(1)'^2$. This model has also three families. This is due to the fact that there exist three fixed points, which correspond to the origin, with respect to the first twist θ, the second one θ^2 and the third one θ^3 in the presence of the Wilson line.

3. 2 Dynamical Breaking in Observable Sector

We assume that one E_6 gauge group, which is interpreted as the observable one, is spontaneously broken down to G_{ST} with appropriate chiral multiplets at $M_{GUT} \sim 10^{16}$ GeV and that this symmetry breakings occur dynamically by the condensation of certain bound states whose constituents are fundamental **27**. This is a new phenomenological possibility in the building of the unified string models.

At first sight, we wonder why the only one E_6 gauge symmetry is spontaneously broken although our string models have the same structure of observable sector as that of hidden one. We shall give a possible solution to this question by taking a simple model § as an example. Consider the Lagrangian density for two real scalar fields ϕ and ϕ' with mass m,

$$\mathcal{L}=\frac{1}{2}(\partial_\mu\phi)^2-\frac{1}{2}m^2\phi^2-\frac{1}{4}g\phi^4+\frac{1}{2}(\partial_\mu\phi')^2-\frac{1}{2}m^2\phi'^2-\frac{1}{4}g\phi'^4-\frac{1}{2}\lambda\phi^2\phi'^2 \tag{25}$$

$$=\frac{1}{2}(\partial_\mu\phi)^2+\frac{1}{2}(\partial_\mu\phi')^2-V(\phi,\ \phi') \tag{26}$$

where g and λ are some coupling constants. The above Lagrangian density $\mathcal{L}$ is invariant under the exchange of one scalar field for another one. The vacuum solution is obtained by solving the following simultaneous equations,

$$\frac{\partial V}{\partial\phi}=m^2\phi+g\phi^3+\lambda\phi\phi'^2=0 \tag{27}$$

$$\frac{\partial V}{\partial\phi'}=m^2\phi'+g\phi'^3+\lambda\phi^2\phi'=0 \tag{28}$$

There exist various vacuum solutions for the different parameter regions of m,g and λ. For example,

1. $\phi=\phi'=0$
2. $\phi=\phi'=\pm\sqrt{\frac{-m^2}{g+\lambda}}$
3. $\phi=\pm\sqrt{\frac{-m^2}{g}}, \qquad \phi'=0$

   or

   $\phi=0, \qquad \phi'=\pm\sqrt{\frac{-m^2}{g}}$

and so on. The solution 1 and 2 are the symmetric solutions. On the other hand, the exchange symmetry is spontaneously broken in the solution 3. We postulate that a similar mechanism is applied to our models. That is, the Z_2 invariance under the exchange of the observable sector for the hidden one is supposed to be spontaneously broken on the presence of non-renormalizable interactions between them. Since the non-renormalizable interactions among fundamental chiral multiplets are suppressed by the factor $(\frac{1}{M_{Pl}})^n$, they seem not to be available. However, if renormalizable interactions among composite fields are induced effectively, the strength of the interactions can become order 1. Thus the exchange symmetry can be broken at the composite level.

§ This model is regarded as a special case of the $SU(2)_L\times SU(2)_R\times U(1)_{B-L}$ models of weak interactions whose Higgs potential induces the parity breaking [18].

Next we explain how E_6 is broken down to G_{ST} by the bound states of 27 dimensional chiral multiplet Φ_{27}. The mechanism is the almost same as the case of dynamical breakings by fundamental fermion-fermion condensations. According to Ref. [15], the breakings occur at the following three stages,

$$E_6 \xrightarrow{\mathbf{351}} SU(5)\times U(1)$$

$$SU(5) \xrightarrow{\mathbf{24}} G_{ST}$$

$$G_{ST} \xrightarrow{\mathbf{2}} SU(3)\times U(1)_{EM}$$

where the numbers over the arrows represent the dimensions of representation of the scalar objects that are needed to the corresponding breakings. In our models, there are chiral supermultiplets $(\mathbf{5}^* + \mathbf{5} + \mathbf{1})$ under $SU(5)$ subgroup of E_6 besides usual $(\mathbf{5}^* + \mathbf{10})$ in MSSM. The condensation of bound states made of $(\mathbf{5}^* + \mathbf{5} + \mathbf{1})$ induces the breaking $E_6 \rightarrow G_{ST}$ simultaneously, and the breaking $G_{ST} \rightarrow SU(3) \times U(1)_{EM}$ occurs by the condensation of bound states made of usual $(\mathbf{5}^* + \mathbf{10})$ similar to top-quark condensation [19].

At last, we discuss the mass spectra after the dynamical symmetry breakings. We can construct a dimension 4 operator O_4, which is invariant under $E_6 \times U(1)^2 \times E_6' \times U(1)'^2$ transformation and Z_7 transformation, as

$$O_4 \equiv \varepsilon_{ijk}\Phi_{27}^i\Phi_{27}^j\Phi_{27}^k \tag{29}$$

where Φ_{27}^i is the 27-dimensional chiral supermultiplet whose index i represents a family which it belongs to. The dimension 5 operator O_5 is constructed as

$$O_5 \equiv \bar{\Phi}_{27}^i\Phi_{27}^i\bar{\Phi}_{27}^j\Phi_{27}^j \tag{30}$$

where $\bar{\Phi}_{27}^i$ is the anti-chiral supermultiplets. We can write down the following $E_6 \times U(1)^2 \times E_6' \times U(1)'^2$ and Z_7 invariant interaction by using of O_4 and O_5,

$$S_{int} = \int d^4x d^2\theta f \Phi_{27}^1\Phi_{27}^2\Phi_{27}^3 + (h.\,c.) + \int d^4x d^2\theta d^2\bar{\theta} f_{ij}\bar{\Phi}_{27}^i\Phi_{27}^i\bar{\Phi}_{27}^j\Phi_{27}^j \tag{31}$$

where f is a Yukawa coupling constant and f_{ij} is a certain coupling constant suppressed by M_{Pl} as $f_{ij} = G_{ij}/M_{Pl}$. Here G_{ij} is some dimensionless parameter. The reduction of Φ_{27}^i is done as

$$\Phi_{27}^i = \Phi_{16}^i + \Phi_{10}^i + \Phi_1^i \tag{32}$$

under the subgroup $SO(10)$. Here the numbers in the lower index in the right-hand side represent the dimensions of representation under $SO(10)$ and Φ_{16}^i includes usual matters. We suppose that $\bar{\Phi}_{10}^i + \bar{\Phi}_1^i$ form bound states

$$\frac{G_{ij}}{2M_{Pl}}(\bar{\Phi}^i_{10}+\bar{\Phi}^i_1)\cdot(\bar{\Phi}^j_{10}+\bar{\Phi}^j_1)\sim\Psi^{ij}_{54}+\Psi^{ij}_{45}+\Psi^{ij}_{1(1)}+\Psi^{ij}_{10(1)}+\Psi^{ij}_{1(2)}+\Psi^{ij}_{10(2)} \quad (33)$$

and that all of them condense in the following form

$$<\Psi^{ij}_a> \sim \delta^{ij} M_{GUT} \quad (34)$$

Of course, the chiral symmetry is not broken without SUSY breaking in the framework of the system described by only S_{int} as interaction terms in the same way as the case of SUSY Nambu-Jona-Lasinio model [20]. More complex interactions are required to generate the above condensation. Here it is supposed that such effective interactions exist after the compactification and induce chiral symmetry breakings. And let us discuss the mass spectra by using only S_{int}. Extra matters $\Phi^i_{10} + \Phi^i_1$ acquire heavy masses $O(M_{GUT})$ by the above condensation if the values of $G_{ii}(i = 1, 2, 3)$is order 1. Note that the $U(1)^2$ are also broken by these condensations. Furthermore, we suppose that the composite field $\Psi^{ij}_{10(3)}$, which is made of $\bar{\Phi}^i_{16}$ such as

$$\frac{G_{ij}}{2M_{Pl}}\bar{\Phi}^i_{16}\cdot\bar{\Phi}^j_{16}\sim\Psi^{ij}_{126}+\Psi^{ij}_{120}+\Psi^{ij}_{10(3)} \quad (35)$$

condenses at weak scale M_W similar to top-condensation [19],

$$<\Psi^{ij}_{10(3)}>\sim\delta^{3i}\delta^{3j} M_W \quad (36)$$

In this way, $SU(2)_L \times U(1)_Y$ gauge symmetry can be broken dynamically and the mass of top quark can be order M_W.

There are many problems. The mechanism of mass generation in the first two families is not known. The origin of Kobayashi-Maskawa mixing angle [21] also is not known.

3. 3 Dynamical Breaking in Hidden Sector

In this subsection, we discuss the dynamical symmetry breaking in the hidden sector which triggers off SUSY breaking. The gauge coupling constant of the other E_6' gauge group becomes strong at some energy scale Λ_C ($\Lambda_C < M_{GUT}$). Then, its gaugino can condense and break SUSY. We shall examine by using renormalization group equations (RGEs) whether it is anticipated that the parameters M_S, $\alpha_U(\equiv g(M_S)^2/4\pi)$, Λ_C and $m_{3/2}$take phenomenologically reasonable values. The values of M_{GUT} and α_{GUT} ($\equiv g(M_{GUT})^2/4\pi$) are determined [4] by the use of RGEs in the usual SUSY GUT scheme as follows,¶

$$M_{GUT} = 10^{16\pm0.3} GeV$$

and

¶After E_6 gauge symmetry breaking, the extra supermultiplets ($5^* + 5 + 1$) acquire heavy masses $O(M_{GUT})$ and hence they don't contribute on the analysis of RGEs.

$$\alpha_{GUT}^{-1} = 25.7 \pm 1.7$$

Hereafter we use the values $M_{GUT} = 10^{16}$GeV and $\alpha_{GUT}^{-1} = 25.7$, for simplicity. The running of $\alpha(\mu)(\equiv g(\mu)^2/4\pi)$ in SSTs yields the following solution of RGEs at one loop level,

$$\alpha(\mu)^{-1} = \alpha_U^{-1} + \frac{b}{2\pi} \ln \frac{\Lambda_U}{\mu} + \Delta \tag{37}$$

where Δ represents string threshold effects [22] and Λ_U is the string unification scale in $\overline{MS}$ scheme. In our models, eq.(37) holds in the energy range from Λ_U to M_{GUT} for the observable E_6 and from Λ_U to Λ_C for the hidden E_6'. In the case of the overall modulus, b and Δ are given by,

$$b = -3C_2(G) + \sum_i T(R_i) \tag{38}$$

and

$$\Delta = \frac{b'}{4\pi} \ln(2ReT|\eta(T)|^4) \tag{39}$$

where $C_2(G)$ is the quadratic Casimir invariant of group G, $T(R_i)$ is the index of R_i-representation, $b' = 3C_2(G) - \sum_i (3 + 2n_i) T(R_i)$ (The number n_i is called 'modular weight' and $n_i = -1$ for untwisted matters.) and $\eta(T)$ is the Dedekind function. In our models, $b = -27$, and $\Delta = 0$ because Δ depends on the untwisted moduli which is absent in Z_3 and Z_7 orbifold models. Now when we estimate Λ_U at M_S, we can obtain values for Λ_U and α_U^{-1} by setting the scale μ at M_{GUT} in eq. (37),

$$\Lambda_U \sim 1.8 \times 10^{18} GeV \tag{40}$$

and

$$\alpha_U^{-1} \sim 48 \tag{41}$$

We now calculate the confining scale Λ_C of E_6' and estimate the gravitino mass $m_{3/2}$. The scale Λ_C is connected with the scale M_{GUT} and the structure constant α_{GUT} by RGEs as follows,

$$\Lambda_C = M_{GUT} \cdot exp(\frac{2\pi}{b} \alpha_{GUT}^{-1}) \tag{42}$$

We find $\Lambda_C \simeq 2.5 \times 10^{13}$ GeV and $m_{3/2}$ is estimated from eq.(7) as $m_{3/2} \simeq O(1)$ TeV which is nothing but what is assumed in the SUSY phenomenology.

The gaugino condensation and scalar condensations are tightly constrained by Konishi anomaly relation [23]. If there exists a gauge non-singlet chiral matter which does not appear in the superpotential, the gaugino does not condense. However, since all 27 dimensional chiral superfield Φ_{27}' have $E_6 \times U(1)^2 \times E_6' \times U(1)'^2$ invariant

Yukawa couplings in our models, the E_6' gaugino can condense.

3. 4 Soft SUSY Breaking Terms

In this subsection, we examine whether soft SUSY breaking terms have universal structures or not in our models. The soft SUSY breaking terms mean terms that break SUSY without the introduction of quadratic divergences. For example, scalar mass terms, gaugino mass terms and trilinear scalar coupling terms. They are generated in the observable sector through the spontaneous SUSY breaking in the hidden sector.

First, the masses of scalars ϕ_i are given by [24],

$$m_i^2 = m_{3/2}^2 + n_i\ m_0^2 + V_0 \tag{43}$$

where m_0^2 is some mass parameter which depends on the model and V_0 is the cosmological constant. Here the modular weight n_i takes a different value between untwisted matters and twisted matters. In our first model, all matter multiplets **27**'s belong to the untwisted sector and hence the soft SUSY breaking mass terms of **27**'s have a universal structure which seems needed for the sufficient suppression of flavour changing neutral currents [25]. In our second model, they also have a universal structure since the modular weights have a universal value $n_i = -2$.

Second, the masses of gaugino λ_a are given by [24]

$$M_a = \alpha_a m_{3/2}\,(Ck_a + b'_a C_a) \tag{44}$$

where C and C_a are some constant factors and α_a are the structure constants. There is no threshold correction in our models, so $C_a = 0$ and M_a have a universal structure, ||

$$M_1 : M_2 : M_3 = \alpha_1 : \alpha_2 : \alpha_3 \tag{45}$$

Last, the trilinear scalar coupling terms are as follows [24],

$$L_{tril} = -\ m_{3/2} A\hat{h}\varepsilon_{ijk}\phi_{27}^i\phi_{27}^j\phi_{27}^k + (h.\ c.) \tag{46}$$

$$\hat{h} \equiv \hat{C} + \sum_{i=1}^{3} n_i\,\hat{D} \tag{47}$$

where A,$\hat{C}$and $\hat{D}$are some constant factors.

4 Summary

We have proposed a new approach in the superstring phenomenology provided by the dynamical symmetry breaking. And we have applied it to two Z_7 orbifold models and shown that these models have phenomenologically interesting features (three families, $m_{3/2} \simeq O(1)$ TeV, the universal soft SUSY breaking terms, $\cdots$). It is no

||Here the Kac-Moody levels k_a are chosen as $\frac{3}{5}k_1 = k_2 = k_3 = 1$ as is described in section 2.

wonder that non-renormalizable interactions are generated by the compactification of extra space. However, it is well known [20] that the chiral symmetry is not broken without SUSY breaking in SUSY Nambu-Jona-Lasinio models. Therefore, more complex interactions are required to generate the dynamical breaking at the GUT scale. It is not clear to us whether such effective interactions are indeed induced in the string compactification. There are of course many remaining problems. None of them is not fully analysed, since it requires more detailed non-perturbative dynamics of SSTs. So it is an important subject that we investigate the dynamical aspects of SSTs.

Acknowledgement

The author would like to give special thanks to Prof. T. Yanagida, who has collaborated with the author in the search of 'Topics of String Unification', for many useful advises and discussions. The authors also would like to thank H. Inoue for useful discussions.

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