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Orbifold unification for the gauge and Higgs fields and their couplings

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Abstract

We present an orbifold GUT model in which the Higgs trilinear couplings are unified with the three Standard Model gauge couplings. The model is constructed as an $N = 2$ supersymmetric $SU(8)$ gauge theory in six dimensions, which is reduced to a supersymmetric Standard Model with three singlets and extra $U(1)$ factors upon compactification. Such an unification is in good agreement with experiments. The predicted upper limit for the lightest CP-even neutral Higgs boson is somewhat larger than in the MSSM, and can be tested in the upcoming large hadron collider.

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1. Introduction

The minimal supersymmetric Standard Model (MSSM) is the most natural extension of the Standard Model (SM). It elegantly solves the gauge hierarchy problem, contains neutralino as the cold dark matter candidate, and naturally accommodates the gauge coupling unification [1,2]. Depending on the supersymmetry breaking mechanism, it also has distinct

predictions for the sparticles' (supersymmetric partners of the SM particles) spectra which can be tested at the upcoming colliders such as the large hadron collider (LHC) and the future international linear collider (ILC). Despite all these successes, there are several unanswered questions within the MSSM. Why the bilinear supersymmetric Higgs mass μ (in the superpotential) involving the up and down Higgs superfields, $\mu H_u H_d$, is at the TeV scale but not at the Planck scale? This is known as the μ problem. Also the predicted upper bound for the mass of the lightest CP-even neutral Higgs boson h^0 is around 130 GeV [3],

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which is not much higher than the current experimental lower limit, 114 GeV. Moreover, the prediction for the proton decay rate through dimension-5 operators is uncomfortably close to the current experimental bounds.

These problems in the MSSM have prompted many to consider the possible extensions to the next to the minimal supersymmetric Standard Model (NMSSM) [4] by the addition of one or more Higgs singlets to the usual two doublets present in the MSSM. Such extensions can be used to solve the μ problem, and extend the upper mass limit for the Higgs mass. Additionally, in an orbifold GUT model (such as the one being proposed here), the doublet–triplet splitting problem is naturally solved, thus avoiding the possible problem with the proton decay rate [5]. With the addition of one singlet Higgs field, we can have the additional trilinear interaction terms $\lambda SH_u H_d$ and $\kappa S^3/3$ in the superpotential. After minimization of the scalar Higgs potential, S can obtain a vacuum expectation value (VEV) around the supersymmetry-breaking scale (M_{SUSY}), and generate an effective μ term with $\mu = \lambda \langle S \rangle$. Thus, the μ problem is solved. However, such NMSSM lacks definite predictions because the Higgs couplings λ and κ are completely arbitrary. Is there a theoretical framework in which the values of λ and κ get determined?

In this work, we present a supersymmetric Standard Model with three SM singlets and following superpotential

$$W = \lambda H_u H_d S - \kappa S S_1 S_2, \quad (1)$$

where the Yukawa couplings λ and κ get determined in terms of the gauge couplings, and thus making this model predictive and testable by experiment. The idea is simple, and very attractive, and has led to the unification of gauge and Yukawa couplings [6]. We use the framework of extra dimensions with supersymmetry. The two Higgs doublets, as well as the singlets are all part of the gauge multiplet in higher dimensions, and the non-minimal interactions involving the λ and κ are just part of the gauge interactions in higher dimensions. We present the realization of this idea below.

2. Formalism and the model

We consider a theory with a gauge symmetry G in six dimensions (6D) with $N = 2$ supersymmetry.

(The two extra dimensions will be compactified on a suitable orbifold such that the gauge symmetry is broken down to the SM with possibly some extra $U(1)$ factors, and the supersymmetry is broken down to $N = 1$.) The $N = 2$ supersymmetry in 6D corresponds to $N = 4$ supersymmetry in 4D, and thus only the gauge multiplet can be introduced in the bulk. In terms of 4D $N = 1$ language, the six-dimensional gauge multiplet contains a vector multiplet, V , and three chiral multiplets Σ_1 , Σ_2 , and Σ_3 in the adjoint representation of the gauge group G . The bulk action [7], written in 4D $N = 1$ language and in the Wess–Zumino gauge, contains the following trilinear term of the chiral multiplets

$$S = \int d^6x \int d^2\theta \frac{1}{kg^2} \text{Tr}(-\sqrt{2}\Sigma_1[\Sigma_2, \Sigma_3]) + \text{H.C.}, \quad (2)$$

where k is the normalization factor for the group generators. If the SM singlet Higgs field, and the up- and down-type Higgs doublets are contained in the zero modes of the chiral multiplets Σ_1 , Σ_2 , and Σ_3 , the gauge interaction term, Eq. (2), includes the trilinear Higgs interaction term $\lambda SH_u H_d$ with the coupling λ determined in terms of the gauge coupling g . In this construction, the singlet Higgs field, the two Higgs doublets, and the gauge fields are all unified in a single multiplet of the gauge symmetry group G in higher dimensions. In the NMSSM, we also need a cubic term, $\kappa S^3/3$ for the singlet field S to develop a VEV. We can see from Eq. (1) that we need three SM singlet Higgs fields to be present in the zero modes of Σ_1 , Σ_2 , and Σ_3 leading to a trilinear term $\kappa S S_1 S_2$.

We now address what bulk gauge symmetry we need to unify both λ and κ with the gauge couplings. To obtain both the Higgs doublets and the singlets as zero modes in 4D from the extra-dimensional components of the higher-dimensional gauge multiplet, and also to break the supersymmetry to $N = 1$, the minimal bulk gauge symmetry needed is $SU(4)_W$. In this case, $SU(4)_W$ is broken down to $SU(2)_L \times U(1)_Y \times U(1)'$ upon compactification, and the adjoint **15**-dimensional representation will have two doublets, H_u and H_d , and a singlet S as the zero modes. In this case, we can obtain only the trilinear interaction $\lambda SH_u H_d$ from the bulk gauge interaction with $\lambda = g_2$, where g_2 is the weak gauge coupling.

Table 1

The zero modes of the chiral multiplets Σ_1 , Σ_2 and Σ_3 , and their quantum numbers under the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_\alpha \times U(1)_\beta \times U(1)_\gamma$ gauge symmetry. The anti-symmetric subscripts Q_{ij} ($Q_{ij} = -Q_{ji}$) are the charges under the $U(1)_Y \times U(1)_\alpha \times U(1)_\beta \times U(1)_\gamma$ gauge symmetry: $Q_{12} = (-5/6, 0, 0, 0)$, $Q_{13} = (-1/3, 8/15, -1/2, -1/2)$, $Q_{14} = (-1/3, 8/15, 1/2, -1/2)$, $Q_{15} = (-1/3, 8/15, 0, 1)$, $Q_{23} = (1/2, 8/15, -1/2, -1/2)$, $Q_{24} = (1/2, 8/15, 1/2, -1/2)$, $Q_{25} = (1/2, 8/15, 0, 1)$, $Q_{34} = (0, 0, 1, 0)$, $Q_{35} = (0, 0, 1/2, 3/2)$, $Q_{45} = (0, 0, -1/2, 3/2)$

Zero modes for Σ_i (in the i th row)

$Q_X: (\mathbf{3}, \mathbf{2})_{Q_{12}}$	$H'_u: (\mathbf{1}, \mathbf{2})_{Q_{23}}$	$S: (\mathbf{1}, \mathbf{1})_{Q_{45}}$	$\bar{D}_\delta: (\bar{\mathbf{3}}, \mathbf{1})_{Q_{51}}$
$D_\delta: (\mathbf{3}, \mathbf{1})_{Q_{13}}$	$S_2: (\mathbf{1}, \mathbf{1})_{Q_{34}}$	$\bar{D}_X: (\bar{\mathbf{3}}, \mathbf{1})_{Q_{41}}$	$H_d: (\mathbf{1}, \mathbf{2})_{Q_{52}}$
$H_u: (\mathbf{1}, \mathbf{2})_{Q_{24}}$	$S_1: (\mathbf{1}, \mathbf{1})_{Q_{53}}$	$S'_1: (\mathbf{1}, \mathbf{1})_{Q_{35}}$	$H'_d: (\mathbf{1}, \mathbf{2})_{Q_{42}}$

The minimal gauge symmetry in the bulk to include both the λ and κ terms from the zero mode bulk interaction is $SU(5)_W$. The $SU(5)_W$ gauge symmetry in the bulk, upon compactification to 4D, is broken down to $SU(2)_L \times U(1)_Y \times U(1)' \times U(1)''$. The $SU(5)_W$ adjoint representation, $\mathbf{24}$, decomposed under the $SU(2)_L \times U(1)_Y \times U(1)' \times U(1)''$ contains the two Higgs doublets, H_u and H_d , as well as three singlets S , S_1 and S_2 as zero modes. The bulk gauge interaction contains the $\lambda S H_u H_d$, as well as $\kappa S S_1 S_2$ terms, giving rise to $\lambda = \kappa = g_2$ at the compactification scale. With $SU(3)_C \times SU(5)_W$ as the gauge symmetry in the bulk, we can include color interaction, but this does not unify the three SM gauge couplings. Thus we are naturally lead to an $SU(8)$ gauge symmetry in the bulk to unify all three SM gauge couplings with λ and κ .

The model we propose for the gauge and Higgs trilinear coupling unification is in six dimensions with $N = 2$ supersymmetry, and $SU(8)$ gauge symmetry. The two extra dimensions x_5 and x_6 are compactified on a T^2/Z_6 orbifold, which is obtained from torus T^2 by moduloing the Z_6 equivalent class: $z \sim \omega z$, where z is the complex coordinate of the extra dimensions and $\omega = e^{i\pi/3}$. The transformation property for the vector multiplet, V is

$$V(x^\mu, \omega z, \omega^{-1}\bar{z}) = R V(x^\mu, z, \bar{z}) R^{-1}, \quad (3)$$

where R is an 8×8 matrix and $R^6 = I$. The transformation rules for the three chiral multiplets Σ_1 , Σ_2 , and Σ_3 are obtained by multiplying the right-hand side of the Eq. (3) by the additional factors ω^{-1} , ω^{-1-m} , and ω^{2+m} , respectively, where m is an integer. These transformations keep the bulk action invariant and non-

trivial R breaks the bulk gauge symmetry G at the 4D fixed point [8]. We choose the matrix R to be

$$R = \text{diag}(+1, +1, +1, \omega^{n_1}, \omega^{n_1}, \omega^{n_2}, \omega^{n_3}, \omega^{n_4}). \quad (4)$$

Then, for unequal values of the integers n_1 , n_2 , n_3 and n_4 , upon compactification to 4D, the $SU(8)$ gauge symmetry breaks to $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_\alpha \times U(1)_\beta \times U(1)_\gamma$, and the $N = 4$ supersymmetry in 4D is broken down to $N = 1$ by an appropriate choice of m . For the choice of $m = 1$, and $n_1 = 5$, $n_2 = 4$, $n_3 = 2$ and $n_4 = 1$, the zero modes of the $\mathbf{63}$ -dimensional vector multiplet are the gauge bosons (and gauginos) corresponding to the unbroken gauge symmetry, while the zero modes of the three $\mathbf{63}$ -dimensional chiral multiplets Σ_1 , Σ_2 and Σ_3 , and their quantum numbers are given in Table 1. These zero modes include the Higgs bosons of above model in the compactified 4D theory. From the bulk action in Eq. (2), we obtain the non-minimal Higgs interactions for the zero modes of the kinetic-normalized chiral multiplets

$$\begin{aligned} S = \int d^6x \left[\int d^2\theta g_6 (S H_u H_d - S S_1 S_2 \right. \\ \left. - Q_X \bar{D}_X H_u + S_2 H'_u H'_d - S'_1 H'_u H_d \right. \\ \left. + S'_1 D_\delta \bar{D}_\delta) + \text{H.C.} \right], \quad (5) \end{aligned}$$

where g_6 is the 6D gauge coupling whose mass dimension is -1 .

The extra $U(1)$ gauge symmetries, $U(1)_\alpha \times U(1)_\beta \times U(1)_\gamma$ can be broken at the compactification scale via Higgs mechanism, and thus the exotic quarks Q_X , \bar{D}_X , D_δ , \bar{D}_δ , and the exotic doublets H'_u and H'_d can acquire superheavy masses at this scale after these extra $U(1)$ gauge symmetries are broken. This can be achieved on the 3-brane at the Z_6 fixed point, for example, $z = 0$, by introducing two exotic quarks \bar{Q}'_X and D'_δ with quantum numbers $(\bar{\mathbf{3}}, \mathbf{2})_{(5/6, 0, -1, 0)}$ and $(\mathbf{3}, \mathbf{1})_{(-1/3, 8/15, -1, 1)}$ respectively under the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_\alpha \times U(1)_\beta \times U(1)_\gamma$ gauge symmetry. We also introduce a SM singlet Higgs field \tilde{S}_2 which has the same quantum number as that of S_2 and is localized on the 3-brane at $z = 0$. After \tilde{S}_2 gets a VEV, the exotic quarks and Higgs doublets can obtain the vector-like masses through the following brane-

localized superpotential,

$$W = \tilde{S}_2 H'_u H'_d + \tilde{S}_2 D_\delta \bar{D}_X + \tilde{S}_2 Q_X \bar{Q}'_X + \tilde{S}_2 D'_\delta \bar{D}_\delta. \tag{6}$$

Similarly, S'_1 can be made superheavy. Furthermore, the extra $U(1)$ symmetries can have gauge anomalies in 4D since some of vector-like pairs of the zero modes are projected out by orbifolding. Thus, we need to add brane-localized fields which has extra $U(1)$ charges to cancel the gauge anomalies, and this can be an origin of the extra fields such as \tilde{S}_2 and breaking of extra $U(1)$ symmetries.

After the $U(1)_\alpha \times U(1)_\beta \times U(1)_\gamma$ gauge symmetry is broken, we have the relevant superpotential

$$S = \int d^6x \left[\int d^2\theta g_6 (S H_u H_d - S S_1 S_2) + \text{H.C.} \right]. \tag{7}$$

Integrating out the two extra dimensions, we obtain, at the GUT scale,

$$g_3 = g_2 = g_1 = \lambda = \kappa = g_6 / \sqrt{V}, \tag{8}$$

where V is the volume of extra dimensions, and g_3, g_2 and $g_1 \equiv \sqrt{5/3} g_Y$ are the 4D gauge couplings for the SM gauge symmetry $SU(3)_C, SU(2)_L$ and $U(1)_Y$, respectively. The true unification scale of the couplings is the cutoff scale (M_*) in the orbifold models, but for simplicity, we here assume that the compactification scale is the GUT scale ($\sim 2 \times 10^{16}$ GeV) so that the Higgs trilinear couplings λ and κ can be predicted. We also neglect the brane-localized gauge kinetic terms which are assumed to be suppressed by $\sqrt{V} M_*$ compared to the bulk kinetic term.

In our model, we introduce three families of the SM fermions on the 3-brane at the Z_6 fixed point $z = 0$. Moreover, we emphasize that the hypercharge interaction in the SM can be one linear combination of the $U(1)_\alpha$ and $U(1)_\gamma$ in above $SU(8)$ model, and then the hypercharge normalization may not be determined in this model as in the usual orbifold GUT models when all the SM fermions are brane-localized fields. However, if we identify \bar{D}_δ as a right-handed down-type quark field in the presented choice of Z_6 charge assignment, the hypercharge normalization in the SM can be the same as usual $SU(5)$ normalization.

These additional zero modes such as exotic quarks, extra doublets and singlets can also be eliminated from the zero modes of the compactified 4D spectrum by considering a 7-dimensional theory with $N =$

1 supersymmetry and $SU(8)$ bulk gauge symmetry, and compactifying the three extra dimensions on a $T^2/Z_6 \times S^1/Z_2$ orbifold. Due to the orbifold projections, the bulk $SU(8)$ gauge symmetry is broken directly down to the SM-like gauge symmetry, and there are only one pair of Higgs doublets and three SM singlets in the Higgs sector arising from the zero modes of bulk vector multiplet. In this case, however, the hypercharge normalization is not determined completely. When we consider larger gauge group, the quark and lepton fields can also be unified with the bulk gauge multiplet and then the hypercharge normalization will be fixed naturally [9].

The superpotential in Eq. (7) contains five SM neutral complex scalar fields and (in the general case) three phase symmetries in the scalar potential. One of these is the $U(1)$ gauge symmetry related to the Z boson, implying two unwanted global symmetries. These will generally be spontaneously broken, implying two massless Goldstone bosons. One of these has large H'_d and H'_u components and is clearly excluded by the known experiment. The second consists mainly of the S, S_1 and S_2 fields, and is most likely also excluded, although a detailed investigation is beyond the scope of this Letter. Let us give one possible solution to this problem. To break the $U(1)_\beta \times U(1)_\gamma$ gauge symmetry at GUT scale by Higgs mechanism, we introduce the SM singlet vector-like fields $(N_1, \bar{N}_1), (N_2, \bar{N}_2)$ and (N_3, \bar{N}_3) . And the quantum numbers for N_1, N_2 and N_3 under the $U(1)_Y \times U(1)_\alpha \times U(1)_\beta \times U(1)_\gamma$ gauge symmetry are $(0, 0, 7/2, -21/2), (0, 0, -9, 0)$ and $(0, 0, 4, 12)$, respectively. So, we can have the following non-renormalizable terms in the superpotential

$$W' = h_1 \frac{S^7 N_1}{M_*^5} + h_2 \frac{S_2^9 N_2}{M_*^7} + h_3 \frac{S_1^8 N_3}{M_*^6}, \tag{9}$$

where h_1, h_2, h_3 are Yukawa coupling constants. After N_1, N_2 and N_3 get VEVs, we obtain the effective superpotential at low scale

$$W' = h'_1 \frac{S^7}{M_*^4} + h'_2 \frac{S_1^9}{M_*^6} + h'_3 \frac{S_2^8}{M_*^5}, \tag{10}$$

where $h'_i = h_i \langle N_i \rangle / M_*$. As shown in Ref. [10], these non-renormalizable terms in the above superpotential do not generate the dangerous quadratically divergent tadpoles for the SM singlet fields S and S_i [11]. Also, we do not have global symmetries in our model, so,

there are no massless Goldstone bosons and there is no domain wall problem after the Higgs doublets and SM singlets obtain VEVs.

3. Phenomenology

We now briefly discuss the phenomenological implication of our model, in particular, the implication of the effective superpotential given by Eq. (7), and the unification of Higgs trilinear couplings with the SM gauge couplings, Eq. (8). The unification prediction can be tested by using the appropriate renormalization group equations (RGEs) for these couplings. For numerical calculations, we consider two-loop RGE runnings for the SM gauge couplings and top quark Yukawa coupling y_t , and one-loop RGE runnings for the Higgs trilinear couplings (λ and κ) [12], with conversion from $\overline{\text{MS}}$ scheme to dimensional reduction ($\overline{\text{DR}}$) scheme. We also include the standard supersymmetric threshold corrections at low energy by choosing a single scale $M_{\text{SUSY}} = M_Z$ where M_Z is the Z-boson mass [13]. The relevant RGEs are

$$\frac{d\alpha_\lambda}{dt} = \frac{\alpha_\lambda}{2\pi} \left(\alpha_\kappa + 4\alpha_\lambda + 3\alpha_t - \frac{3}{5}\alpha_1 - 3\alpha_2 \right), \quad (11)$$

$$\frac{d\alpha_\kappa}{dt} = \frac{\alpha_\kappa}{2\pi} (3\alpha_\kappa + 2\alpha_\lambda), \quad (12)$$

$$\frac{d\alpha_t}{dt} = \left[\frac{d\alpha_t}{dt} \right]_{\text{MSSM}} + \frac{\alpha_t}{2\pi} \left(\alpha_\lambda - \frac{1}{4\pi} \alpha_\lambda (3\alpha_t + 3\alpha_\lambda + \alpha_\kappa) \right), \quad (13)$$

$$\frac{d\alpha_2}{dt} = \left[\frac{d\alpha_2}{dt} \right]_{\text{MSSM}} + \frac{\alpha_2^2}{8\pi^2} (-2\alpha_\lambda), \quad (14)$$

$$\frac{d\alpha_1}{dt} = \left[\frac{d\alpha_1}{dt} \right]_{\text{MSSM}} + \frac{\alpha_1^2}{8\pi^2} \left(-\frac{6}{5}\alpha_\lambda \right), \quad (15)$$

where t is the log of renormalization scale, $\alpha_i = g_i^2/(4\pi)$, $\alpha_\lambda = \lambda^2/(4\pi)$, $\alpha_\kappa = \kappa^2/(4\pi)$, $\alpha_t = y_t^2/(4\pi)$ and the bracket $[]$ denotes the corresponding two-loop RGEs in the MSSM. We use the values of SM gauge couplings at M_Z in Ref. [14] and the top quark mass to be 178 GeV.

Our results for the unification of the gauge and Higgs trilinear couplings using the RGEs, Eqs. (11)–(15), are shown in Fig. 1 for $\tan\beta = 5$ where $\tan\beta \equiv \langle H_u^0 \rangle / \langle H_d^0 \rangle$. Also the predictions for the couplings λ

and κ at the weak scale, M_Z , for various values of $\tan\beta$ are shown in Fig. 2.

Since the values of λ and κ are predicted in our model, we can calculate the upper bound on the mass of the lightest CP-even neutral Higgs boson h^0 by using the full one-loop and the leading logarithmic two-loop corrections [3]. In Fig. 3, we plot the upper bounds for the h^0 mass in the MSSM, the NMSSM and our model versus $\tan\beta$. Note that for a large range of $\tan\beta$, the mass bound in our model is larger than in the MSSM, but less than in the NMSSM in which the values of λ and κ are arbitrary. For this calculation we use the approximation that the mass of CP-odd Higgs (M_A) is order of the square root of the arithmetic average of the stop squared-mass eigenvalues (M). The

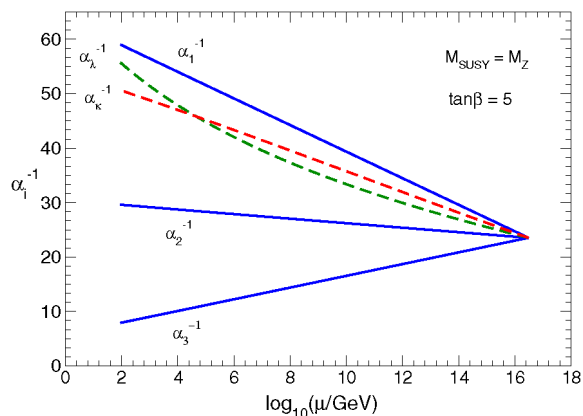


Fig. 1. For $\tan\beta = 5$, the unification of the SM gauge couplings ($\alpha_1, \alpha_2, \alpha_3$) and Higgs trilinear couplings (α_λ and α_κ).

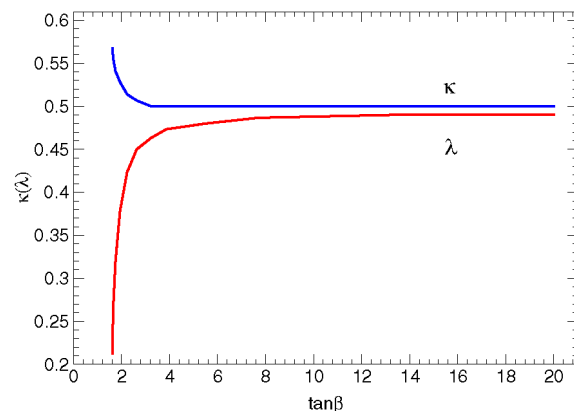


Fig. 2. The Higgs trilinear couplings λ and κ at the weak scale versus $\tan\beta$.

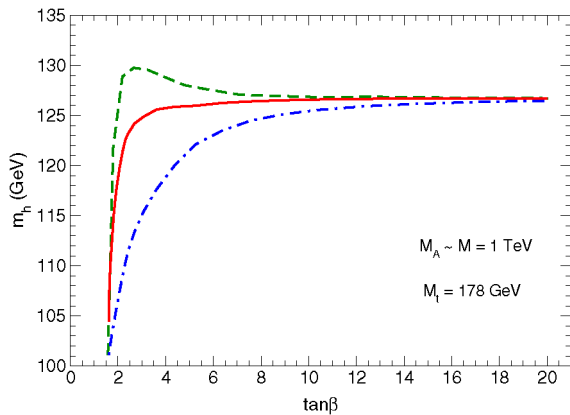


Fig. 3. The upper bounds on the lightest CP-even neutral Higgs h^0 mass in the MSSM, the NMSSM and our model versus $\tan\beta$: the blue dash-dot line, green dash line, and red solid line correspond to the MSSM, the NMSSM, and our model, respectively. (For interpretation of the references to colour in this figure legend, the reader is referred to the web version of this Letter.)

validity of our prediction can be tested in the upcoming LHC.

4. Conclusions

We have presented a supersymmetric Standard Model in which the Higgs trilinear couplings λ and κ are unified with the three SM gauge couplings at the unification scale. This is an orbifold GUT model in 6D with $N = 2$ supersymmetry. The symmetry is broken down to the SM gauge symmetry in four dimensions via orbifold compactification as well as via Higgs mechanism. The unification prediction is in good agreement with experiments. The predicted upper bound for the lightest CP-even Higgs mass is somewhat larger than in the MSSM, and can be tested at the LHC. The detail model buildings which also include the possible unification of the third-family Yukawa couplings, and their phenomenological consequences will be presented elsewhere.

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References

- [1] S. Dimopoulos, H. Georgi, Nucl. Phys. B 193 (1981) 150; S. Dimopoulos, S. Raby, F. Wilczek, Phys. Rev. D 24 (1981) 1681; N. Sakai, Z. Phys. C 11 (1981) 153; L.E. Ibanez, G.G. Ross, Phys. Lett. B 105 (1981) 439; M.B. Einhorn, D.R.T. Jones, Nucl. Phys. B 196 (1982) 475; W.J. Marciano, G. Senjanovic, Phys. Rev. D 25 (1982) 3092.
- [2] U. Amaldi, W. de Boer, H. Furstenau, Phys. Lett. B 260 (1991) 447; J.R. Ellis, S. Kelley, D.V. Nanopoulos, Phys. Lett. B 249 (1990) 441; P. Langacker, M.X. Luo, Phys. Rev. D 44 (1991) 817; C. Giunti, C.W. Kim, U.W. Lee, Mod. Phys. Lett. A 6 (1991) 1745.
- [3] Y. Okada, M. Yamaguchi, T. Yanagida, Prog. Theor. Phys. 85 (1991) 1; J.R. Ellis, G. Ridolfi, F. Zwirner, Phys. Lett. B 257 (1991) 83; H.E. Haber, R. Hempfling, Phys. Rev. Lett. 66 (1991) 1815; M. Carena, J.R. Espinosa, M. Quiros, C.E.M. Wagner, Phys. Lett. B 355 (1995) 209.
- [4] P. Fayet, Nucl. Phys. B 90 (1975) 104; H.P. Nilles, M. Srednicki, D. Wyler, Phys. Lett. B 120 (1983) 346; J.M. Frere, D.R.T. Jones, S. Raby, Nucl. Phys. B 222 (1983) 11; J.P. Derendinger, C.A. Savoy, Nucl. Phys. B 237 (1984) 307.
- [5] Y. Kawamura, Prog. Theor. Phys. 103 (2000) 613; Y. Kawamura, Prog. Theor. Phys. 105 (2001) 999; Y. Kawamura, Prog. Theor. Phys. 105 (2001) 691; G. Altarelli, F. Feruglio, Phys. Lett. B 511 (2001) 257; A.B. Kobakhidze, Phys. Lett. B 514 (2001) 131; L. Hall, Y. Nomura, Phys. Rev. D 64 (2001) 055003; A. Hebecker, J. March-Russell, Nucl. Phys. B 613 (2001) 3; T. Li, Phys. Lett. B 520 (2001) 377; T. Li, Nucl. Phys. B 619 (2001) 75.
- [6] I. Gogoladze, Y. Mimura, S. Nandi, Phys. Lett. B 562 (2003) 307; I. Gogoladze, Y. Mimura, S. Nandi, Phys. Rev. Lett. 91 (2003) 141801; I. Gogoladze, Y. Mimura, S. Nandi, Phys. Rev. D 69 (2004) 075006; I. Gogoladze, Y. Mimura, S. Nandi, K. Tobe, Phys. Lett. B 575 (2003) 66; T. Kobayashi, S. Raby, R.J. Zhang, Nucl. Phys. B 704 (2005) 3.
- [7] N. Marcus, A. Sagnotti, W. Siegel, Nucl. Phys. B 224 (1983) 159;

- N. Arkani-Hamed, T. Gregoire, J. Wacker, JHEP 0203 (2002) 055.
- [8] T. Li, JHEP 0403 (2004) 040.
- [9] I. Gogoladze, T. Li, Y. Mimura, S. Nandi, hep-ph/0504082.
- [10] S.A. Abel, S. Sarkar, P.L. White, Nucl. Phys. B 454 (1995) 663;
S.A. Abel, Nucl. Phys. B 480 (1996) 55.
- [11] H.P. Nilles, M. Srednicki, D. Wyler, Phys. Lett. B 124 (1983) 337;
A.B. Lahanas, Phys. Lett. B 124 (1983) 341;
U. Ellwanger, Phys. Lett. B 133 (1983) 187;
J. Bagger, E. Poppitz, Phys. Rev. Lett. 71 (1993) 2380;
- J. Bagger, E. Poppitz, L. Randall, Nucl. Phys. B 455 (1995) 59;
V. Jain, Phys. Lett. B 351 (1995) 481;
C. Panagiotakopoulos, K. Tamvakis, Phys. Lett. B 446 (1999) 224.
- [12] S.P. Martin, M.T. Vaughn, Phys. Rev. D 50 (1994) 2282;
S.F. King, P.L. White, Phys. Rev. D 52 (1995) 4183.
- [13] P. Langacker, N. Polonsky, Phys. Rev. D 47 (1993) 4028;
M. Carena, S. Pokorski, C.E.M. Wagner, Nucl. Phys. B 406 (1993) 59.
- [14] S. Eidelman, et al., Particle Data Group Collaboration, Phys. Lett. B 592 (2004) 1.