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Fat branes, orbifolds and doublet-triplet splitting

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Abstract. A simple higher dimensional mechanism of the doublet-triplet splitting is presented in a five dimensional supersymmetric SU(5) GUT on S^1/Z_2. The splitting of multiplets is realized by a mass term of the Higgs hypermultiplet which explicitly breaks SU(5) gauge symmetry. Depending on the sign of the mass, zero mode Higgs doublets and triplets are localized on either side of the fixed points. The mass splitting is realized due to the difference of magnitudes of the overlap with a brane localized or a bulk singlet field. No unnatural fine-tuning of the parameters is needed. The proton stability is ensured by locality – without symmetries.

The doublet-triplet splitting problem is one of the notorious problems in grand unified theories (GUTs) [1]. Many proposals for this problem have been made in four dimensional models [2] or recently in higher dimensional models [3]; see also for string-derived models [4]. In our previous papers [5, 6], we have presented a doublet-triplet splitting mechanism in the context of the fat brane scenario [7]. In this approach, a conventional mass splitting, namely the weak-GUT scale splitting, can be realized by an overlap of zero mode wave functions between the doublet and the triplet Higgs fields without unnatural fine-tuning of the parameters [5]. Interestingly, it has also been shown that an alternative mass splitting, namely the weak-TeV scale splitting, can be realized by the same machanism [6]. Proton stability is guranteed by the strong suppression of the coupling of matter fields to the triplet Higgs due to a small overlap of wave functions. This TeV scale triplet Higgs scenario deserves attention as an alternative signature of GUT instead of the proton decay.

The setup itself has crucial problems although these scenarios are very attractive. In the fat brane scenario, extra dimensions are considered to be non-compact; otherwise the effective four dimensional theory becomes vector-like. This non-compactness leads to the circumstance that the gravity and gauge fields cannot propagate in the bulk and have to be localized on a fat brane (domain wall) to obtain finite coupling constants. As is well known, it is highly nontrivial to realize the localization of the gravity and gauge fields on a domain wall in infinite extra dimensions. Even if the localization is realized by the "quasi-localization" mechanism [8, 9], zero mode wave functions have to be almost flat on a domain wall because of the charge universality constraints. It seems to be very difficult to obtain such flat wave functions.

To avoid this situation, we consider a theory on an orbifold.[1] By construction, the bulk gravity and gauge fields propagate in a finite extra space and it is easy to obtain flat zero mode wave functions in flat extra dimensions. A first explicit realization of the fat brane scenario on an orbifold has been given in [10]. By developing extra coordinate dependent vacuum expectation values (VEVs) of the parity odd scalar field, the kink solution (domain wall) is constructed. The chiral fermion zero modes which couples to the scalar field generating domain wall are localized on either fixed point depending on the sign of the coupling constant.

In this letter, we apply this mechanism to the doublet-triplet splitting in a five dimensional supersymmetric (SUSY) SU(5) GUT compactified on S^1/Z_2.[2] There are fixed points at $y = 0, \pi R$, where y denotes the fifth dimensional coordinate, and R is the compactification radius.

We shall focus on the Higgs sector. Two hypermultiplets are introduced as Higgs multiplets:

$$H_1 = (H_1(\mathbf{5}), H_1^c(\mathbf{5}^*)), \quad H_2 = (H_2(\mathbf{5}^*), H_2^c(\mathbf{5})), \quad (1)$$

where the representations under SU(5) are specified. We assign their Z_2 parities as follows:

$$H_1(-y) = +H_1(y), \quad H_1^c(-y) = -H_1^c(y), \quad (2)$$

$$H_2(-y) = +H_2(y), \quad H_2^c(-y) = -H_2^c(y). \quad (3)$$

[1] In the smooth compact dimensional case, S^1 for instance, the anti-domain wall as well as the domain wall appears. This case cannot only yield chiral fermions but also cause the instability of the system.

[2] For other applications, see [11, 12].

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The action we consider is given by

$$S_5 = \int \mathrm{d}^5x \Big\{ \Big[H_i^\dagger \mathrm{e}^{-V} H_i + H_i^c \mathrm{e}^{V} H_i^{c\dagger} \Big]_{\theta^2\bar\theta^2} + \Big[H_i^c \left(\partial_y - \frac{1}{\sqrt{2}} \Phi \right) H_i - \frac{1}{2} M(y) H_i^c H_i + \delta(y - \pi R) \frac{S}{M_\mathrm{P}} H_1 H_2 \Big]_{\theta^2} + \mathrm{h.c.} \Big\} \quad (i = 1, 2), \tag{4}$$

where the action S_5 is written in terms of the 4D, $\mathcal{N} = 1$ superspace formalism [13], S is a singlet chiral superfield, which is assumed to be localized on the brane at $y = \pi R$.[3] M_5, M_P are the Planck scales in five and four dimensions, which are related by $M_5^3 \pi R = M_\mathrm{P}^2$. Here, the vector superfield V and the chiral superfield Φ in the adjoint representation are explicitly given by

$$V = -\theta\sigma^\mu\bar\theta A_\mu + \mathrm{i}\bar\theta^2\theta\lambda_1 - \mathrm{i}\theta^2\bar\theta\bar\lambda_1 + \frac{1}{2}\theta^2\bar\theta^2 D, \tag{5}$$

$$\Phi = \frac{1}{\sqrt{2}}(\Sigma + \mathrm{i}A_5) + \sqrt{2}\theta\lambda_2 + \theta^2 F, \tag{6}$$

where A_μ ($\mu = 0, 1, 2, 3$) is a gauge field in four dimensions, $\lambda_{1,2}$ are gauginos, F, D are auxiliary fields, A_5 is an extra component of the gauge field, and Σ is a real scalar field in the adjoint representation. Note that an extra dimensional coordinate dependent mass term is introduced in the second line of (4). The mass parameter $M(y)$ is given by

$$M(y) = M_5 \,\mathrm{diag}(2, 2, 2, -3, -3)\varepsilon(y), \tag{7}$$

where $\varepsilon(y)$ is the sign function with respect to y. Thus, SU(5) is explicitly broken to the standard model (SM) gauge group at the cutoff scale M_5.

Throughout this paper, we consider the following SUSY vacuum at M_5[4]:

$$\langle \Sigma \rangle = \langle H_i \rangle = \langle H_i^c \rangle = \langle S \rangle = 0. \tag{8}$$

Let us concentrate on the fermionic components of the action to study the zero modes in the background (7) and (8),

$$S_5 \supset \psi_i^c \mathrm{i}\sigma^\mu \partial_\mu \bar\psi_i^c + \bar\psi_i \mathrm{i}\bar\sigma^\mu \partial_\mu \psi_i - \psi_i^c \partial_y \psi_i + \bar\psi_i \partial_y \bar\psi_i^c + \frac{1}{2} M(y)(\psi_i^c \psi_i + \bar\psi_i^c \bar\psi_i) - \delta(y - \pi R) \left(\frac{S}{M_\mathrm{P}} \psi_1 \psi_2 + \frac{S^*}{M_\mathrm{P}} \bar\psi_1 \bar\psi_2 \right) \quad (i = 1, 2), \tag{9}$$

where (ψ_i, ψ_i^c) obviously denote the fermionic components of (H_i, H_i^c). By expanding in modes as follows:

$$\psi_i(x, y) = \sum_n \psi_i^{(n)}(x) f^{(n)}(y),$$

$$\psi_i^c(x, y) = \sum_n \psi_i^{c(n)}(x) f^{c(n)}(y), \tag{10}$$

we obtain the mode equations

$$0 = m_n \bar\psi_1^{(n)} f_1^{c(n)} + \bar\psi_1^{(n)} \partial_y \bar f_1^{(n)} - \frac{1}{2} M(y) \bar\psi_1^{(n)} \bar f_1^{(n)}, \tag{11}$$

$$0 = m_n \bar\psi_2^{(n)} f_2^{c(n)} + \bar\psi_2^{(n)} \partial_y \bar f_2^{(n)} - \frac{1}{2} M(y) \bar\psi_2^{(n)} \bar f_2^{(n)}, \tag{12}$$

$$0 = m_n \bar\psi_1^{c(n)} f_1^{(n)} - \bar\psi_1^{c(n)} \partial_y \bar f_1^{c(n)} - \frac{1}{2} M(y) \bar\psi_1^{c(n)} \bar f_1^{c(n)} + \delta(y - \pi R) \frac{S^*}{M_\mathrm{P}} \bar\psi_2^{(n)} \bar f_2^{(n)}, \tag{13}$$

$$0 = m_n \bar\psi_2^{c(n)} f_2^{(n)} - \bar\psi_2^{c(n)} \partial_y \bar f_2^{c(n)} - \frac{1}{2} M(y) \bar\psi_2^{c(n)} \bar f_2^{c(n)} + \delta(y - \pi R) \frac{S^*}{M_\mathrm{P}} \bar\psi_1^{(n)} \bar f_1^{(n)}, \tag{14}$$

where the mass in four dimensions is defined as

$$-\mathrm{i}\sigma^\mu \partial_\mu \bar\psi_i^{c(n)} = m_n \psi_i^{(n)}, \quad -\mathrm{i}\bar\sigma^\mu \partial_\mu \psi_i^{c(n)} = m_n \bar\psi_i^{(n)}. \tag{15}$$

It is easy to find solutions of the above equations (11)–(14) in the bulk:

$$f_{1,2}^{(n)}(y) = N_n \exp\left[\frac{1}{2} \int_0^y \mathrm{d}x_5 M(x_5) \right] \cos(m_n y), \tag{16}$$

$$f_{1,2}^{c(n)}(y) = N_n^c \exp\left[-\frac{1}{2} \int_0^y \mathrm{d}x_5 M(x_5) \right] \sin(m_n y), \tag{17}$$

where $N_n^{(c)}$ are normalization constants.

On the other hand, the following boundary conditions at $y = \pi R$ should be satisfied from (13) and (14):

$$\tan(m_n \pi R) = -\frac{S}{2M_\mathrm{P}} \frac{\psi_2}{\psi_1^c} \exp\left[\int_0^{\pi R} \mathrm{d}y M(y) \right], \tag{18}$$

$$\tan(m_n \pi R) = -\frac{S}{2M_\mathrm{P}} \frac{\psi_1}{\psi_2^c} \exp\left[\int_0^{\pi R} \mathrm{d}y M(y) \right]. \tag{19}$$

Mass eigenvalues can be obtained by eliminating $\psi_{1,2}$ in (18) and (19) as follows:

$$m_n = \frac{1}{R} \left(n + \frac{1}{\pi} \arctan\left[\frac{S}{2M_\mathrm{P}} \exp\left(\int_0^{\pi R} \mathrm{d}y M(y) \right) \right] \right) \quad (n = 0, 1, 2, \ldots). \tag{20}$$

For $m_0 = 0$, zero mode wave functions take the form

$$f_{1,2}^{(0)}(y) \simeq \exp\left[\frac{1}{2} \int_0^y \mathrm{d}x_5 M(x_5) \right] \tag{21}$$

$$= \begin{cases} \sqrt{\frac{2M_5}{\mathrm{e}^{2M_5\pi R} - 1}} \exp[M_5 y] \ \text{(for triplets)}, \\ \sqrt{\frac{3M_5}{1 - \mathrm{e}^{-3M_5\pi R}}} \exp\left[-\frac{3}{2} M_5 y \right] \ \text{(for doublets)}. \end{cases} \tag{22}$$

[3] The cases of localization at $y = 0$ and being constant in the bulk will be discussed later.

[4] Later we will consider the VEV of S. We assume that S will take a VEV below the energy scale of M_5 by, for an example, an inverted hierarchy scenario.

It turns out that the triplet Higgs zero mode is peaked at $y = \pi R$ and the doublet Higgs zero mode at $y = 0$.

The doublet-triplet splitting is realized by the coupling of the Higgs doublets to S,

$$\delta(y - \pi R)\left[\frac{S}{M_\mathrm{P}} H_1 H_2\right]_{\theta^2} \Rightarrow \tag{23}$$

$$\begin{cases} m_3 = \frac{\langle S\rangle}{M_\mathrm{P}} \frac{2M_5}{\mathrm{e}^{2M_5\pi R}-1}\mathrm{e}^{2M_5\pi R} \simeq \frac{\langle S\rangle}{M_\mathrm{P}} 2M_5, \\ m_2 = \frac{\langle S\rangle}{M_\mathrm{P}} \frac{3M_5}{1-\mathrm{e}^{-3M_5\pi R}}\mathrm{e}^{-3M_5\pi R} \simeq \frac{\langle S\rangle}{M_\mathrm{P}} 3M_5 \mathrm{e}^{-3M_5\pi R} \simeq m_W, \end{cases}$$

where the weak scale m_W is $m_W \simeq 100\,\mathrm{GeV}$, and $m_{2,3}$ are the doublet, triplet Higgs masses respectively. Remarkably, the triplet mass is unsuppressed since the overlap with a singlet is large, while the doublet mass is *exponentially* suppressed since the overlap with a singlet is very small.

In order to constrain the parameters in our model further, let us study proton decay. We assume that the matter fields are localized on the brane at $y = 0$. Then the coupling of matter fields to the triplet Higgs H_3 is given by

$$\delta(y)\left[\frac{Y}{\sqrt{M_\mathrm{P}}} H_3 QQ\right]_{\theta^2}, \tag{24}$$

where Y is the Yukawa coupling, and the Q mean the SM matter superfields. Integrating out with respect to the fifth coordinate y and substituting the value of zero mode wave functions of Higgs triplets at $y = 0$ yield a 4D effective Yukawa coupling Y_eff, which is determined by the normalization constant,

$$Y_\mathrm{eff} \simeq \sqrt{\frac{2M_5}{M_\mathrm{P}}}\mathrm{e}^{-M_5\pi R}. \tag{25}$$

Let us first discuss the proton decay constraints from a dimension six operator. There are two sources for a dimension six operator; one is X, Y gauge boson exchange, the other is the triplet Higgs scalar exchange at tree level. In our scenario, the amplitude of the X, Y gauge boson exchange is $1/m_{X,Y}^2 \simeq 1/M_5^2$, where $m_{X,Y}$ are X, Y gauge boson masses of order 5D Planck scale M_5, since the GUT symmetry is broken by a mass of a hypermultiplet (7). Namely, the constraint is simply $M_5 > 10^{16}\,\mathrm{GeV}$. On the other hand, the amplitude of the proton decay process from the triplet Higgs boson exchange is

$$\frac{(Y_\mathrm{eff})^2}{m_3^2} \simeq \left(\frac{M_\mathrm{P}}{2M_5\langle S\rangle}\right)^2 \left(\frac{2M_5}{M_\mathrm{P}}\mathrm{e}^{-2M_5\pi R}\right) = \frac{M_\mathrm{P}}{2M_5\langle S\rangle^2}\mathrm{e}^{-2M_5\pi R}. \tag{26}$$

Therefore, the triplet Higgs boson exchange constraints to be satisfied lead to

$$\sqrt{\frac{2M_5}{M_\mathrm{P}}}\langle S\rangle \mathrm{e}^{M_5\pi R} > 10^{16}\,\mathrm{GeV}. \tag{27}$$

Using the second relation in (23), the upper bound on the compactification scale is obtained as

$$R^{-1} < 5.5 \times 10^{17}\,\mathrm{GeV}. \tag{28}$$

Next we turn to the constraints from dimension five operators. The proton decay amplitude coming from the dimension five operator is evaluated to

$$\frac{g^2}{16\pi^2}\frac{Y_\mathrm{eff}^2}{m_3 M_\lambda} \simeq \frac{g^2}{16\pi^2}\frac{y_1 y_2\left(\frac{2M_5}{M_\mathrm{P}}\mathrm{e}^{-2M_5\pi R}\right)}{m_3 M_\lambda}, \tag{29}$$

where g is an SU(2) gauge coupling, $y_{1,2}$ are Yukawa couplings of the first and the second generations in five dimensions, and M_λ is the gaugino mass. Comparing the 4D GUT case, the following condition can be obtained:

$$\frac{1}{\langle S\rangle} \times 10^{16}\mathrm{e}^{-2M_5\pi R} < 1 \Leftrightarrow \frac{3M_5}{M_\mathrm{P}} \times 10^{14}\mathrm{e}^{-5M_5\pi R} < 1; \tag{30}$$

the second expression can be obtained by using the second relation in (23). Summarizing the proton decay constraints, the X, Y gauge boson exchange amplitudes proportional to M_5^{-2} give a lower bound for M_5, namely an upper bound for the compactification scale R^{-1}. On the other hand, the amplitudes from Higgs triplet scalar exchange and the dimension five operator are roughly proportional to M_5/M_P for $M_5/M_\mathrm{P} \ll 1$; therefore an upper bound for M_5 is obtained, namely a lower bound for the compactification scale is obtained. Searching for the allowed parameter region satisfying the constraints of X, Y gauge boson exchange, (28) and (30) with $M_5^3\pi R = M_\mathrm{P}^2$, the following results are obtained:

$$1.8 \times 10^{17}\,\mathrm{GeV} < R^{-1} < 5.5 \times 10^{17}\,\mathrm{GeV}, \tag{31}$$

$$2.6 \times 10^{9}\,\mathrm{GeV} < \langle S\rangle < 6.7 \times 10^{17}\,\mathrm{GeV}, \tag{32}$$

$$2.1 \times 10^{9}\,\mathrm{GeV} < m_3 < 3.9 \times 10^{17}\,\mathrm{GeV}, \tag{33}$$

where the upper (lower) bounds of (31) ((32) and (33)) come from the dimension six proton decay constraints, while the other bounds come from the cutoff scale $\langle S\rangle$ ($\langle S\rangle < M_\mathrm{P}$). Note that the upper bound for $\langle S\rangle$ becomes more stringent since $\langle S\rangle$ depends on M_5 or R^{-1} through the doublet Higgs mass in (23). A large $\langle S\rangle$ corresponds to a small M_5, namely a small R^{-1} and vice versa. It is very interesting in that an intermediate scale triplet Higgs can be consistent with the proton stability. It should be stressed that the proton stability can be ensured not by symmetries but by locality only, namely by strong suppression of the coupling constants coming from overlap of the wave functions between the triplet Higgs and the SM matter.

So far, we have discussed the doublet-triplet splitting and the proton stability in the case with a localized singlet at $y = \pi R$. It is also interesting to consider other situations. What is to happen if S is localized at $y = 0$? In this case, the splitting is realized by

$$\delta(y)\left[\frac{S}{M_\mathrm{P}} H_1 H_2\right]_{\theta^2} \Rightarrow$$

$$\begin{cases} m_3 = \frac{\langle S\rangle}{M_\mathrm{P}} \frac{2M_5}{\mathrm{e}^{2M_5\pi R}-1} \simeq 2M_5 \frac{\langle S\rangle}{M_\mathrm{P}}\mathrm{e}^{-2M_5\pi R}, \\ m_2 = \frac{\langle S\rangle}{M_\mathrm{P}} \frac{3M_5}{1-\mathrm{e}^{-3M_5\pi R}} \simeq 3M_5 \frac{\langle S\rangle}{M_\mathrm{P}} \simeq m_W \;(\simeq 100\,\mathrm{GeV}); \end{cases} \tag{34}$$

then the ratio between the doublet and the triplet Higgs masses is modified as follows:

$$\frac{m_3}{m_2} \simeq \frac{2}{3} \mathrm{e}^{-2M_5 \pi R} \ll 1. \tag{35}$$

The triplet Higgs mass is extremely smaller than the doublet Higgs mass since an overlap between the triplet Higgs localized on $y = \pi R$ and S is very small. The following argument leads us to see that this case is inconsistent with the proton stability. Consider the dimension six operator constraints from the triplet Higgs boson exchange (27). The result is

$$\sqrt{\frac{2M_5}{M_\mathrm{P}}} \langle S \rangle \mathrm{e}^{-M_5 \pi R} > 10^{16}\,\mathrm{GeV}. \tag{36}$$

One can easily check that (36) has no solution. Thus, this case is excluded.

Finally, we consider the case that S is a bulk field and has a constant profile. In this case, the mass splitting is realized by both (23) and (34), which means

$$m_3 \simeq 2M_5 \frac{\langle S \rangle}{M_\mathrm{P}}, \quad m_2 \simeq 3M_5 \frac{\langle S \rangle}{M_\mathrm{P}} \simeq m_W. \tag{37}$$

Remarkably, the triplet Higgs and the doublet Higgs masses are comparable since exponentially suppressed contributions are negligible. Furthermore, the triplet Higgs mass is predicted[5] to become

$$m_3 \simeq \frac{2}{3} m_2. \tag{38}$$

Taking into account the lower bound on the triplet Higgs mass coming from the collider experiments, $m_3 \sim 200\,\mathrm{GeV}$ [14, 15],[6] the proton decay constraints are obtained:

$$R^{-1} < 9.3 \times 10^{16}\,\mathrm{GeV}\ (\mathrm{dim}\ 5), \tag{39}$$

$$R^{-1} < 3.2 \times 10^{16}\,\mathrm{GeV}\ (\text{triplet Higgs exchange}), \tag{40}$$

$$R^{-1} > 5.5 \times 10^{11}\,\mathrm{GeV}\ (X, Y\ \text{gauge boson exchange}). \tag{41}$$

The allowed parameters are obtained as

$$5.5 \times 10^{11}\,\mathrm{GeV} < R^{-1} < 3.2 \times 10^{16}\,\mathrm{GeV}, \tag{42}$$

$$393\,\mathrm{GeV} < \langle S \rangle < 13.5\,\mathrm{TeV}, \tag{43}$$

where, as in the case with S localized at $y = \pi R$, a large $\langle S \rangle$ coressponds to a small M_5, namely a small R^{-1} and vice versa. Thus, the bulk constant S case is very interesting, in that the extremely light Higgs triplets is consistent with the proton decay constraints and might provide an alternative signature of GUT in future collider experiments, as discussed in [6].

[5] Taking into account SUSY breaking, triplet Higgsinos obtain an additional SUSY breaking mass.

[6] We thank G.C. Cho for discussing this point.

Before summarizing this paper, we comment on the gauge coupling unification. In the case of the S brane localized at $y = \pi R$, the gauge coupling unification is achieved by taking the suitable triplet Higgs as $m_3 \geq 10^{16}$ GeV in the range of (33). On the other hand, to recover the gauge coupling unification in the bulk S case, we introduce extra fields $\mathbf{5}' + \mathbf{5}'^*$, another singlet S' localized at $y = \pi R$ and an interaction such as (23)[7]. We have

$$\delta(y - \pi R) \left[\frac{S'}{M_\mathrm{P}} \mathbf{5}' \mathbf{5}'^* \right]_{\theta^2} + \mathrm{h.c.} \tag{44}$$

Then the doublet and triplet Higgs masses are split as (23); thus,

$$m_{3'} \simeq \frac{\langle S' \rangle}{M_\mathrm{P}} 2M_5, \tag{45}$$

$$m_{2'} \simeq \frac{\langle S' \rangle}{M_\mathrm{P}} 3M_5 \mathrm{e}^{-3M_5 \pi R}. \tag{46}$$

Gauge coupling unification requires

$$m_{2'} = m_3 \sim 200\,\mathrm{GeV}, \quad m_{3'} \sim M_5. \tag{47}$$

We find that these requirements are satisfied if we take

$$\langle S' \rangle \sim M_\mathrm{P}, \quad R^{-1} \sim 1.3 \times 10^{17}\,\mathrm{GeV}, \tag{48}$$

which is inconsistent with the above allowed parameter region (42). Thus, the proton stability and the gauge coupling unification are not compatible in the bulk singlet case.

In summary, we have presented a simple higher dimensional mechanism of the doublet-triplet splitting in a five dimensional SUSY SU(5) GUT on S^1/Z_2. The splitting of the multiplets is realized by a mass term of the Higgs hypermultiplet which explicitly breaks SU(5) gauge symmetry. Depending on the sign of the mass, zero mode Higgs doublet and triplet are localized on the opposite side of the fixed points. In a phenomenologically viable setup, the mass splitting is realized due to the difference of magnitudes of the overlap with a brane localized singlet field at $y = \pi R$ not due to the boundary conditions. An unnatural fine-tuning of parameters is not necessary. Gauge coupling unification is easily achieved.

We would like to stress that proton stability is ensured by locality without symmetries! In the recent orbifold GUT literature [16], dimension five baryon- and lepton-number violating operators are forbidden by a $U(1)_R$ symmetry. However, dimension five and six baryon- and lepton-number violating operators in our model are strongly suppressed by a small overlap of the wave functions despite an orbifold setup being adopted.

[7] We need further to forbid the superpotentials $S\mathbf{5}'\mathbf{5}'^*$ and $S' H_1 H_2$ to recover the gauge coupling unification. This can easily be realized by introducing a global Z_2 symmetry, for example, $S, (H_1 H_2)$: even, $S', (\mathbf{5}'\mathbf{5}'^*)$: odd.

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