Light Higgs Triplets in Extra Dimensions

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\textbf{Abstract}

We discuss the possibility that the Higgs triplet can be light (1 TeV in the most interesting case) without contradictions with the proton stability in the context of higher dimensional theory. The proton stability is ensured by the suppression of Yukawa coupling of the Higgs triplet to the matter through its small overlap of wave functions in extra dimensions. The light Higgs triplets might be detected in future collider experiments as an alternative signature of GUT instead of the proton decay. The gauge coupling unification can be preserved by introducing extra bulk matter fields.
One of the serious problems in Grand Unified Theories (GUTs) \[1\] is the doublet-triplet splitting problem. We have to explain in a natural way why the Higgs doublet mass is the weak scale and the Higgs triplet mass is at least the GUT scale after the GUT symmetry breaking. \[1\] In usual, this mass splitting is explained by an unnatural fine-tuning of parameters in the theory. The Higgs triplets must be superheavy (at least the GUT scale \(\simeq 10^{16} \text{GeV.}\)) otherwise the rapid proton decay is caused by the dimension five operators \[17\]. \[2\] If Yukawa coupling constants of the Higgs triplet to the matter are extremely small, the proton decay can be suppressed and the Higgs triplets need not to be superheavy. The question is whether such a situation is naturally possible or not. The answer is yes. If we consider the extra dimensions and the Higgs triplets and the matter are localized at different points in extra dimensions, the effective Yukawa couplings in four dimensions are highly suppressed due to the small overlap of wave functions \[18\].

In this paper, we apply this mechanism to the doublet-triplet splitting and discuss the possibility that the Higgs triplets can be light without contradictions with the proton stability in the extra dimensional framework. Dvali \[3\] has also proposed a similar scenario in four dimensional theory, where Yukawa couplings of the Higgs triplet to the matter are suppressed for group theoretical reasons. However, the complicated superpotential in the Higgs sector is required to obtain the desired vacuum expectation value (VEV) of the adjoint Higgs field. In our scenario, these couplings are suppressed by the dynamics in extra dimensions (see also \[9\].) and our model is very simple.

Let us consider a supersymmetric (SUSY) SU(5) GUT in five dimensions, for concreteness. The model is based on Ref. \[15\]. The action of the Higgs sector is

\[
S = \int d^4 x d y \left[ \int d^2 \theta (H^c e^{-V} H + H^c e^V H^c + \bar{H}^c e^V \bar{H} + \bar{H}^c e^{-V} \bar{H}^c) + \left\{ \int d^2 \theta \left( H^c (\partial_y + X(y) + M) H + \bar{H}^c (\partial_y + \bar{X}(y) + \bar{M}) \bar{H} \right) + \delta(y) \int d^2 \theta \left( \lambda_1 \text{tr}(X^2 \Sigma) + \lambda_2 \text{tr}(\bar{X}^2 \Sigma) + \lambda_3 \text{tr}(X \Sigma^2) + \lambda_4 \text{tr}(\bar{X} \Sigma^2) + \frac{1}{2} m_0 \text{tr}(\Sigma^2) \right) + \text{h.c.} \right\} \right],
\]

where \(H(\bar{H}), H^c(\bar{H}^c)\) are left-handed (charge conjugated right-handed) \(\mathcal{N} = 1\) SUSY in four dimensional chiral superfield components of the single \(\mathcal{N} = 1\) SUSY in five dimensional chiral superfield \(H(\bar{5}) = (H, \bar{H}^c)\) and \(\bar{H}(\bar{5}) = (\bar{H}, H^c)\). \(5, \bar{5}\) are the corresponding representations of SU(5). \(X(y), \bar{X}(y)\) are the bulk fields in the 24 dimensional representation under SU(5). \(\Sigma\) is an usual SU(5) GUT adjoint Higgs field, which is assumed to

\[1\] There has been many proposals for this problem \[2\]-\[16\].

\[2\] Throughout this paper, R-parity is assumed.
be localized on the brane at $y = 0$. We assume that $X(y), \bar{X}(y)$ depends on $y$, and $M, \bar{M}$ do not. $\lambda_{1-4}$ are dimensionless constants and $m_0$ is a mass parameter. This formulation of the action Eq. (I) is useful because it is written by using the $N = 1$ superfield formalism and $N = 1$ SUSY in four dimensions is manifest [19, 20]. From F-flatness conditions $\partial W/\partial X = 0$ and $\partial W/\partial \bar{X} = 0$, one obtains

$$H^c H - \frac{1}{5} \text{tr}(H^c H) = 0, \quad \bar{H}^c \bar{H} - \frac{1}{5} \text{tr} (\bar{H}^c \bar{H}) = 0,$$

$$2\lambda_1 X(0) + \lambda_3 \Sigma = 0, \quad 2\lambda_2 \bar{X}(0) + \lambda_3 \Sigma = 0. \quad (2)$$

It is remarkable that Eqs. (3) connect $\langle X(0) \rangle$ and $\langle \bar{X}(0) \rangle$ in the bulk with $\langle \Sigma \rangle$ on the brane at $y = 0$. Using Eqs. (3), the last term in Eq. (I) reproduces the Higgs superpotential in the minimal $SU(5)$ GUT. Expanding $H, H^c, \bar{H}$ and $\bar{H}^c$ by the mode functions as

$$H(x, y) = \sum_n \phi_n(y) H_n(x), \quad H^c(x, y) = \sum_n \phi^c_n(y) H^c_n(x), \quad (4)$$

$$\bar{H}(x, y) = \sum_n \tilde{\phi}_n(y) \bar{H}_n(x), \quad \bar{H}^c(x, y) = \sum_n \tilde{\phi}^c_n(y) \bar{H}^c_n(x), \quad (5)$$

where $x$ denotes the coordinates of the four dimensional space-time, the equations of motions for the zero mode wave functions of Higgs fields are obtained

$$(\partial_y + X(y) + M) \phi_0(y) = 0, \quad (6)$$

$$(-\partial_y + X(y) + M) \phi^c_0(y) = 0, \quad (7)$$

$$(\partial_y + \bar{X}(y) + \bar{M}) \tilde{\phi}_0(y) = 0, \quad (8)$$

$$(-\partial_y + \bar{X}(y) + \bar{M}) \tilde{\phi}^c_0(y) = 0. \quad (9)$$

Assuming that $X(y) = X(0) + a^2 y, \bar{X}(y) = \bar{X}(0) + a^2 y$ in a small region of the point crossing zero, where $a$ is a dimensionful constant, we obtain two Gaussian normalizable zero mode wave functions

$$\phi_0(y) \sim \exp \left\{ -\frac{a^2}{2} \left( y - \frac{X(0) + M}{a^2} \right)^2 \right\}, \quad (10)$$

$$\tilde{\phi}_0(y) \sim \exp \left\{ -\frac{a^2}{2} \left( y - \frac{\bar{X}(0) + \bar{M}}{a^2} \right)^2 \right\}. \quad (11)$$

3 One might think that the GUT breaking VEV of the bulk field can be directly obtained from the minimization of the potential. But it is impossible because $N = 1$ SUSY in five dimensions highly constrains the form of the superpotential.

4 The other zero modes should be vanished since they are not normalizable.
Before discussing the doublet-triplet splitting problem in detail, we comment on various scales in our model. There are three typical mass scales, the Planck scale in five dimensions $M_\ast$, the wall thickness scale $L^{-1}$, which should be considered as the compactification scale and the inverse width of Gaussian zero modes $a^{-1}$. As explained in Ref. [18], for the description to make sense, the wall thickness $L$ should be larger than the inverse width of Gaussian zero modes $a^{-1}$ so that the wall has enough width to trap matter and Higgs modes. Furthermore, $a^{-1}$ should be smaller than or equal to the five dimensional Planck length $M_\ast^{-1}$,

$$L^{-1} < a \leq M_\ast.$$  

(12)

We take $L^{-1}$ to be $M_{GUT}$ in order to preserve the gauge coupling unification. The five dimensional Planck scale $M_\ast$ can be taken to be about $10^{17}$ GeV or $10^{18}$ GeV. Throughout this paper, $M_\ast$ is taken to be $10^{18}$ GeV and $a \simeq M_\ast$ for simplicity.

Now, we propose two scenarios of the doublet-triplet splitting problem. The first one which realizes the doublet-triplet splitting is based on the shining mechanism [19]. We introduce a singlet superfield and consider the overlap between the Higgs fields and the singlet field. As explained in Ref. [15], the simplest case without a singlet is not realistic because the Higgs doublets are too apart from each other to yield the hierarchy between the top and the bottom Yukawa couplings naturally.

The action of the singlet sector is

$$S = \int d^4xdy \left[ \int d^4\theta(S^\dagger S + S^{c\dagger}S^c) + \left\{ \int d^2\theta S^c(\partial_y + m_s)S - \delta(y) \int d^2\theta JS^c + h.c. \right\} \right],$$

(13)

where $S$ is a bulk SU(5) singlet superfield, $S^c$ is its conjugated superfield, $J$ is a constant and $m_s$ is a mass parameter. F-flatness conditions lead to

$$S = \theta(y)Je^{-m_s y}, \quad S^c = 0,$$

(14)

where $\theta(y)$ is a step function of $y$. The doublet-triplet splitting is achieved by the coupling

$$\frac{1}{\sqrt{M_\ast}} \int d^4xdy \left\{ \int d^2\theta S(x, y)H(x, y)\bar{H}(x, y) + h.c. \right\}$$

(15)

On this wall, the matter and Higgs fields are localized.

*Assuming the top Yukawa coupling constants in five and four dimensional theories and the bottom Yukawa coupling constant in five dimensional theory to be of order unity, the effective bottom Yukawa coupling constant in four dimensional theory becomes smaller than $O(10^{-21})$. In order to explain the observed bottom mass, we have to assume an unnatural huge bottom Yukawa coupling constant in five dimensional theory.
\[ \simeq M_s \exp \left[ -\frac{1}{2M_s^2} \left\{ (X(0) + M)^2 + (\bar{X}(0) + \bar{M})^2 \right\} + \frac{(X(0) + M + \bar{X}(0) + \bar{M} - m_s)^2}{4M_s^2} \right] \times \int d^4x \sigma(x) \right] \]

\[ \times \delta^4(x) H_0(x) H_0(x) + h.c., \tag{16} \]

where \( J \simeq M_s^{3/2} \) are assumed. As mentioned in Ref. \[15\], an R-symmetry for instance has to be imposed to forbid the bulk Higgs mass term. In order for Higgs doublets \( H_2, \bar{H}_2 \) to be the weak scale,

\[ M_2 \simeq M_s \exp \left[ -\frac{1}{2} \left\{ (3x + m)^2 + (\bar{x} + \bar{m})^2 \right\} + \frac{1}{4} (-s - 3x + m - 3\bar{x} + \bar{m})^2 \right] \simeq 10^2 \text{GeV} \tag{17} \]

should be satisfied, where \( m_s \equiv sM_s, M \equiv mM_s, \bar{M} \equiv \bar{m}M_s, X(0) = x \text{diag}(2, 2, 2, -3, -3)M_s, \bar{X}(0) = \bar{x} \text{diag}(2, 2, 2, -3, -3)M_s \) where \( s, x, \bar{x}, m \) and \( \bar{m} \) are dimensionless constants.

Unlike Ref. \[15\], we does not impose here that the mass of Higgs triplets should be above the GUT scale because this is not necessarily required to ensure the proton stability in our framework. If we consider the case that the Higgs doublet and anti-doublets (triplet and anti-triplet) are localized at the same point for simplicity, then the condition (17) is written as

\[ \exp\left[ -\frac{1}{2} s(-6x + 2m - \frac{1}{2}s) \right] \simeq 10^{-16}. \tag{18} \]

On the other hand, the mass of Higgs triplets \( H_3, \bar{H}_3 \) is

\[ M_3 \simeq M_s \exp \left[ -(2x + m)^2 + \frac{1}{4} (-s + 4x + 2m)^2 \right] \simeq 10^2 \exp(-5xs), \tag{19} \]

where Eq. \[18\] was used to obtain the last expression. Imposing \( M_3 \geq 1 \text{TeV} \) leads to

\[ -5xs \geq \ln 10. \tag{20} \]

One can easily check that there is a parameter region allowed by Eqs. \[18\] and \[20\] if \( m^2 \geq \frac{7}{5} \ln 10 \simeq 35.46 \).

Since the Higgs doublets are localized at \( y = (3x + m)M_s^{-1} \) and the Higgs triplets are localized at \( y = (2x + m)M_s^{-1} \), the relative distance between them is \( 5|x|M_s^{-1} \). By adjusting \( x \) appropriately, the baryon number violating dimension 5 operators are suppressed enough even if the Higgs triplets are light because Yukawa couplings of the Higgs triplet and the matter field localized around the Higgs doublets are small enough. \[7\]

\[ \text{For} \]

\[ \text{The matter fields must be localized around the Higgs doublet to reproduce fermion masses and mixings.} \]
example, we consider the dimension five operators $QQQL$. Suppose that the relative distance between $Q$ and the Higgs doublets is $q$ in units of $M_*^{-1}$, and the relative distance between $L$ and the Higgs doublets is $l$ in units of $M_*^{-1}$. Assuming the Gaussian zero mode wave functions for the matter fields, one can write down the suppression factor of $QQQL$ as follows.

$$
\frac{1}{M_3} \exp \left[ -\frac{1}{3} \left\{ 3(5|x| - q)^2 + (q - l)^2 + (5|x| - l)^2 \right\} \right] < 10^{-16}.
$$

(21)

The inequality is required to be consistent with the experimental data. Consider the case $M_3 \simeq \text{TeV}$ and using the the typical solution in Ref. [15], we obtain

$$|x| > 1.684.$$

(22)

As for the dimension five operators $UUDE$, the similar estimation tells us that $|x| > 1.653$ is enough for avoiding the rapid proton decay. Therefore, the proton decay via the dimension five operators are suppressed if we adjust the parameter $x$ satisfying the above conditions. We comment on the suppression of the dimension six baryon number violating operators. First, the dimension six operators by the X, Y gauge boson exchange are trivially suppressed since the masses of the X, Y gauge boson are the order of $10^{16}$ GeV. Second, the constraint for the dimension six operators by the Higgs scalar triplet exchange can be written as

$$
\frac{1}{M_3^2} \exp \left[ -\frac{1}{3} \left\{ 3(5|x| - q)^2 + (q - l)^2 + (5|x| - l)^2 \right\} \right] < (10^{-16})^2.
$$

(23)

One can easily see that the bound (22) is enough to satisfy the above constraint because the upper bound of the exponential factor is $\exp \left[ -25|x|^2 \right] \simeq 10^{-30.8}$. In the light of this fact, the Higgs triplets with mass of order TeV is very interesting because we might be able to detect the Higgs triplets in collider experiments as an alternative signature of GUT even if the proton decay cannot be observed.

The second scenario is that the doublet-triplet splitting through the bulk Higgs mass term is acheived as a result of supersymmetry breaking, namely Giudice-Masiero mechanism [21]. Naively, GUT and Giudice-Masiero mechanism are incompatible because the light Higgs triplets necessarily appear and lead to the rapid proton decay. In our case,

\footnote{Here we consider the case that $Q$ and $L$ are localized at the same side close to the Higgs triplets. The coefficients of $QQQL$ is much more suppressed in the case that either or both of $Q$ and $L$ are localized at the opposite side with respect to the Higgs doublets.}
however, this is not true. Even if the Higgs triplets are light, the proton decay is suppressed enough by naturally small Yukawa couplings of the Higgs triplets to the matter. Let us assume that a singlet superfield $Z$ with nonvanishing F-term is localized on the brane at $y = 0$. We consider the following Kähler potential

$$\frac{1}{M_*} \int dyd^2\theta d^2\bar{\theta} \delta(y) (Z^\dagger(x)H(x,y)\bar{H}(x,y) + h.c.) = \frac{F_Z}{M_*} \phi_0(y = 0)\bar{\phi}_0(y = 0) \int d^2\theta H_0(x)\bar{H}_0(x) + h.c..$$  

Substituting $\phi_0(y = 0)$ and $\bar{\phi}_0(y = 0)$ in (10) and (11), the masses of the Higgs triplets and doublets are obtained as follows,

$$M_3 \simeq \frac{F_Z}{M_*} \exp[-(2x + m)^2],$$  

$$M_2 \simeq \frac{F_Z}{M_*} \exp[-(-3x + m)^2].$$  

Here we assumed that the Higgs triplet (doublet) and anti-triplet (anti-doublet) localize at the same point, for simplicity. Requiring $M_2 \simeq M_W$, $\exp[-(-3x + m)^2] \simeq M_* M_W / F_Z$ is obtained. This means $(-3x + m)^2 \simeq 2\ln10$ for $\sqrt{F_Z} \simeq 10^{11}$ GeV. In this case, the masses of the Higgs triplets become

$$M_3 \simeq 10^4 \exp[-(5x \pm \sqrt{2\ln10})^2] \text{ GeV.}$$

In order to be $M_3 \simeq \text{TeV}$, $x \simeq -3.66, -0.628$. As is clear from the earlier discussion, it turns out that the dimension five and six operators are suppressed enough for $x \simeq -3.66$.

Finally, let us discuss the gauge coupling unification. In our scenario with light Higgs triplets, the gauge coupling unification is lost. We can improve this point by simply introducing extra bulk matter fields $5', \bar{5}'$ and by giving the GUT scale mass to the triplet components (denoted by $3'$ and $\bar{3}'$) and the same mass as the Higgs triplets to the doublet components (denoted by $2'$ and $\bar{2}'$). Yukawa couplings between these extra fields and ordinary chiral matter fields can be suppressed by the overlap of wave functions. The gauge coupling unification is preserved since $2'$ and $\bar{2}'$ form a complete SU(5) multiplets with Higgs triplets.

In our scenario, the spectrum of a kind of the gaugino mediation is expected. Gauginos receive volume suppressed masses at the tree level since the wave functions of gaugino zero modes are flat in an extra dimension. Sfermions receive exponentially suppressed masses at the tree level since the wave functions of the matter zero modes are Gaussian and are localized on our wall apart from the brane at $y = 0$. 

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We discuss how the above statement is realized. The action of the extra matter fields is similar to Eq. (1) and the mass splitting is achieved by the bulk mass term

$$\int d^4x dy \int d^2\theta M_5'(x,y)\bar{5}'(x,y).$$

After expanding in mode functions, one obtains masses of the $3', \bar{3}'$ and $2', \bar{2}'$

$$M_{3'} \simeq M_\ast \exp\left[-\{(x' - \bar{x}') + \frac{1}{2}(m' - \bar{m}')(\ln 10)^{1/2}\}\right] \simeq M_{\text{GUT}} \simeq 10^{16}\text{GeV},$$

$$M_{2'} \simeq M_\ast \exp\left[-\left\{-\frac{3}{2}(x' - \bar{x}') + \frac{1}{2}(m' - \bar{m}')\right\}^2\right] = M_3,$$

where $X' = x'\text{diag}(2,2,2,-3,-3)M_\ast$, $\bar{X}' = \bar{x}'\text{diag}(2,2,2,-3,-3)M_\ast$, $M' \equiv m'M_\ast$ and $\bar{M}' \equiv \bar{m}'M_\ast$. $x'$, $\bar{x}'$, $m'$ and $\bar{m}'$ are dimensionless constants. These conditions are written as follows,

$$-\sqrt{2\ln10} \leq (x' - \bar{x}') + \frac{1}{2}(m' - \bar{m}') \leq \sqrt{2\ln10},$$

$$-3(x' - \bar{x}') + (m' - \bar{m}') \simeq \pm \sqrt{s(2x + m - s/4)}.$$ 

One can easily see that these two conditions are satisfied ($x' - \bar{x}' = 3, m' - \bar{m}' = -3$ for example.).

In summary, we have discussed the possibility that the Higgs triplet can be light without contradictions with the proton stability in the context of higher dimensional theory. The proton stability is ensured by the suppression of Yukawa coupling of the Higgs triplet to the matter through its small overlap of wave functions in extra dimensions. Phenomenologically interesting is that Higgs triplets with mass of order TeV might be detected in future collider experiments as an alternative signature of GUT instead of the proton decay. The gauge coupling unification can be preserved by introducing extra bulk matter fields and causing the mass splitting so that the doublet components form a complete SU(5) multiplets with the Higgs triplets and the triplet components become superheavy.

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