Light GUT Triplets and Yukawa Splitting

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Triplet-mediated proton decay in Grand Unified Theories (GUTs) is usually suppressed by arranging a large triplet mass. Here we explore instead a mechanism for suppressing the couplings of the triplets to the first and second generations compared to the Yukawa couplings, so that the triplets can be light. This mechanism is based on a “triplet symmetry” in the context of product-group GUTs. We study two possibilities. The first possibility, which requires the top Yukawa to arise from a non-renormalizable operator at the GUT scale, is that all triplet couplings to matter are negligible, so that the triplets can be at the weak scale, giving new evidence for grand unification. The second possibility is that some triplet couplings, and in particular $Tlb$ and $Tu$, are equal to the corresponding Yukawa couplings. This would give a distinct signature of grand unification if the triplets were sufficiently light. However, we derive a model-independent bound on the triplet mass in this case, which is at least $10^{14}\text{GeV}$. Finally, we construct an explicit viable GUT model based on Yukawa splitting, with the triplets at $10^{14}\text{GeV}$, as required for coupling unification to work. This model requires no additional thresholds below the GUT scale.

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I. INTRODUCTION.

One of the main challenges of supersymmetric Grand Unified Theories (GUTs), is the rate of proton decay mediated by the GUT partners of the Standard-Model Higgs doublets. These color-triplet partners couple to Standard Model fermions violating both baryon and lepton number, through

$$W = y_{ij}^{Tqq}Tq_{i}q_{j} + y_{ij}^{Ta\bar{e}}T\bar{u}_{i}\bar{e}_{j} + y_{ij}^{Tql}Tq_{i}l_{j} + y_{ij}^{T\bar{u}d}T\bar{u}_{i}d_{j}.$$  \hspace{1cm} (1)

Here $T$ and $\bar{T}$ are the color-triplets, barred fields are SU(2) singlets and unbarred fields are SU(2) doublets, and $i, j$ are generation indices. Typically, the triplet couplings of $\mathbf{1}$ and the Yukawa couplings $y_{ij}^{U}H_{U}q_{i}\bar{u}_{j} + y_{ij}^{D}H_{D}q_{i}\bar{d}_{j}$, originate from the same GUT superpotential terms, so that $y_{ij}^{Tqq} = y_{ij}^{Ta\bar{e}} = y_{ij}^{U}$ and $y_{ij}^{Tql} = y_{ij}^{T\bar{u}d} = y_{ij}^{D}$. These couplings mediate proton decay at the level of dimension-five operators. To suppress this contribution, one typically tries to arrange a GUT-scale triplet mass, while keeping the Standard Model Higgses at the weak scale. In fact, even models with triplets at the GUT scale are in conflict with current experimental bounds on the proton lifetime. For example, in minimal SU(5), the lower bound on the triplet mass is about $10^{17}\text{GeV}$.

Here we explore instead the possibility that triplet couplings to matter are smaller than the Yukawas, namely, $y_{ij}^{Tqq}, y_{ij}^{Ta\bar{e}} \ll y_{ij}^{U}$ and $y_{ij}^{Tql}, y_{ij}^{T\bar{u}d} \ll y_{ij}^{D}$. This relaxes the proton decay bound on the triplet mass, allowing the triplets to be lighter than the GUT scale.

Thinking about triplets below the GUT scale is motivated by two reasons. The observation of proton decay, taken together with coupling unification, would be a strong indication for GUTs but still far from conclusive evidence. It is therefore intriguing to see whether the triplets can be sufficiently light that they can give a more direct experimental signature of the GUT. We will exhibit a model in which all triplet couplings to matter are tiny, so that the triplets can be around the weak scale. To allow for coupling unification, we will need to arrange an extra pair of light doublets. Thus, at low energies, there would be new fields in complete SU(5) representations, providing additional evidence for grand unification.

An even more spectacular GUT signature would potentially come from models in which triplet couplings to first and second generation fields are suppressed, so that the proton does not decay too fast, but some triplet couplings to third generation fields are order one. This is naturally the case in many of our models. However, we will show that just the presence of the couplings $y_{33}^{Tqq}$ and $y_{33}^{T\bar{u}d}$ requires the triplets to be heavier than at least $10^{6}\text{GeV}$, in order to satisfy the bounds on proton decay.

As mentioned above, current bounds on the proton lifetime imply a lower bound of $10^{17}\text{GeV}$ on the triplet mass in minimal SU(5). On the other hand, for coupling unification to occur in minimal SU(5), the triplet mass should be close to $10^{14}\text{GeV}$. This mismatch is the second motivation for triplets below the unification scale. We will present a simple model with triplets at $10^{14}\text{GeV}$. In this model, all $T$ couplings to matter are very small, so that the dimension-five contribution to proton decay is suppressed, and the dominant contribution is from $X$ and $Y$ gauge boson exchange. Many more variants with

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these properties can be constructed.

The possibility that triplet couplings to matter are small was also discussed in Ref. [3]. In Ref. [4] fermion masses were assumed to originate from two different Higgs fields, such that the two triplet contributions to proton decay cancel. Ref. [5] considered an SO(10) GUT, and argued that the triplet couplings are forbidden by a symmetry. It is therefore closest in spirit to our current work. It was assumed, however, that the top Yukawa originates from a higher dimension operator suppressed by $M_{\text{GUT}}$ only, and not by a higher scale. Thus, there is no energy region in which there exists a sensible effective theory with all couplings being of the same order of magnitude. The mechanism of Ref. [5] involves an extra dimension.

Our models are all based on a “triplet symmetry”, $U_T$, that distinguishes between the triplets and the doublets $\mathbf{10}$. We take $U_T$ to be either a $U(1)$ or a $Z_N$. Since it does not commute with the GUT group, $U_T$ must arise from the combination of some global symmetry and a subgroup of the GUT. Moreover, in order for some coupling involving the triplet to have a different $U_T$ charge from the corresponding doublet coupling, the GUT group must be semi-simple $\mathbf{10}$. For concreteness, we will take the GUT to be $SU(5) \times SU(5)$. Such a setup was used in Ref. [10] to generate a doublet-triplet mass hierarchy. As was pointed out in Ref. [10], the symmetry may also lead to a suppression of dimension-five operators contributing to proton decay, through the suppression of triplet coupling to fermions. In Ref. [13], explicit models that have these properties were constructed. However, the possibility of exploiting “Yukawa splitting” in order to lower the triplet mass was not explored.

As we will see, a triplet-matter coupling can be different from the corresponding Yukawa coupling whenever the relevant Higgs field and matter fields transform under different $SU(5)$ factors. Then, the Yukawa coupling must arise from a non-renormalizable term in the GUT superpotential, which involves some combination of GUT-breaking fields. The key point is that this operator does not lead to a triplet coupling: the triplet coupling can only come from a different GUT term, involving a different combination of GUT-breaking fields. If some Yukawa coupling has its origin in a non-renormalizable GUT term, it involves some power of $M_{\text{GUT}}/M_{\text{Planck}}$. Thus, if the Planck scale is $10^{18}$ GeV, the top Yukawa must arise from a renormalizable GUT term, and $y_{\bar{T}qq} \sim y_{\bar{T}q\bar{u}} \sim 1$. We therefore distinguish between two classes of models.

1. The top Yukawa coupling is renormalizable at the GUT scale. Then, $y_{\bar{T}qq} \sim y_{\bar{T}q\bar{u}} \sim 1$.

2. The top Yukawa originates from a non-renormalizable term at the GUT scale. In this case, all triplet couplings to matter can be small. This is only possible if the Planck scale is around $10^{17}$ GeV, as would be the case with a small extra dimension. Then, the top Yukawa can be of the order of $O(10^{-1})$ at the GUT scale, with running effects driving it to order one at the weak scale.

II. MODELS: BASIC STRUCTURE.

We now turn to the basic structure of our models, following Ref. [13]. The models have the symmetry $SU(5)_1 \times SU(5)_2 \times U_T^0$, where $U_T^0$ is a global symmetry which we will take to be either a $U(1)$ or a $Z_N$. The symmetry is broken by two sets of bifundamental fields $\Phi_3 \sim (5, 5, 1)$, $\Phi_3 \sim (5, 5, -1)$, $\Phi_2 \sim (5, 5, q)$, and $\Phi_2 \sim (5, 5, -q)$, where the third entry corresponds to the $U_T^0$ charge. For $U_T^0 = Z_N$ ($U_T^0 = U(1)$) we take $q = N/2$ ($q = -N/2$).

For the VEVs

$$
\langle \Phi_3 \rangle = \langle \bar{\Phi}_3 \rangle = \text{diag}(v_3, v_3, v_3, 0, 0),
$$
$$
\langle \Phi_2 \rangle = \langle \bar{\Phi}_2 \rangle = \text{diag}(0, 0, 0, v_2, v_2),
$$

with $v_2 \sim v_3 \sim 10^{16}$ GeV, the symmetry is broken to $[SU(3) \times SU(2) \times U(1)]_{\text{SM}} \times U_T$, where $U_T$ is a combination of $U_T^0$ and a discrete hypercharge subgroup of $SU(5)_1$. The standard-model gauge group lies in the diagonal $SU(5)_3$, so that the standard-model gauge couplings all start from the diagonal $SU(5)$ coupling, and hypercharge is quantized as in minimal $SU(5)$.

As was shown in Ref. [13], one can add three $SU(5)_1$ adjoints and a singlet such that the direction $\mathbf{2}$ is flat, all uneaten GUT-breaking fields get heavy, and the ratio $v_3/v_2$ is naturally of order one.

There are now different possible choices for the MSSM matter fields, since they can transform under either of the two $SU(5)$ gauge groups. This, and the charge assignments under the triplet symmetry, will define the different models we consider.

III. NON-RENORMALIZABLE TOP YUKAWA.

As mentioned above, if the Planck scale is near $10^{17}$ GeV, the top Yukawa may originate from a non-renormalizable term at the GUT scale. It is then easy to construct models in which the triplets can be as light as the weak scale. Take the Standard Model Higgs fields to come from $h \sim (5, 1, 0)$ and $\tilde{h} \sim (5, 1, 0)$ under $SU(5)_1 \times SU(5)_2 \times U_T^0$, and the matter fields to be three copies of $(1, 10, 0)$ and $(1, 5, -q)$. We also add the fields $N \sim (5, 1, 2)$, $\bar{N} \sim (1, 5, -1)$, $M \sim (5, 1, 1)$, $\bar{M} \sim (1, 5, 2)$. $N$ and $\bar{N}$ are needed to restore coupling unification, and $M$ and $\bar{M}$ are required for anomaly cancellation.

The Yukawa couplings now come from

$$
\frac{1}{M_{\text{Planck}}} \Phi_2 h(1, 10)(1, 10) + \frac{1}{M_{\text{Planck}}} \bar{\Phi}_2 \tilde{h}(1, 10)(1, 5),
$$

where we suppress generation indices. All Yukawa couplings are suppressed by $M_{\text{GUT}}/M_{\text{Planck}} \sim 10^{-1}$ at the

\footnote{In the models of Ref. [13], the Dirac mass for the triplets was suppressed, resulting in smaller dimension-five operators.}
GUT scale. Since running to the weak scale enhances the top Yukawa coupling by roughly 3, these models are viable if the order-one coefficient multiplying the top coupling at the GUT scale is around 3.

As long as the triplet symmetry is unbroken, the analogous terms with $\Phi_3$ and $\bar{\Phi}_3$ are forbidden, and the triplets have no couplings to matter. In addition, the triplet symmetry forbids both a triplet mass and a $\mu$ term. Thus, it should ultimately be broken in order to generate a $\mu$ term. The most attractive possibility is that this breaking is related to supersymmetry breaking. Then some triplet-matter couplings would typically be generated, suppressed by powers of the weak scale over the Planck scale.

Since the triplet Higgses are light, coupling unification is lost and we must split another GUT multiplet to restore it. It is easy to do so by allowing the superpotential term $N\bar{\Phi}_3\bar{N}$. The triplets in $N$ and $\bar{N}$ get a $M_{\text{GUT}}$ mass, but their doublet partners remain light.

As mentioned above, the fields $M$ and $\bar{M}$ are only needed for anomaly cancellation. The triplets and doublets of these fields cannot both get mass at the GUT scale because of the triplet symmetry. The simplest possibility is then to forbid their masses altogether in the limit of unbroken supersymmetry. Indeed, with the $U_T$ charges specified above, no masses are allowed for $M$ and $\bar{M}$. Once the triplet symmetry is broken, the triplets and doublets in $M$ and $\bar{M}$, as well as the doublets in $N$ and $\bar{N}$, will get mass around the supersymmetry-breaking scale.

At low energies, we then have two extra 5’s and two extra 5’s: One pair coming from Higgs triplets and the doublets of $N$ and $\bar{N}$, and the other from $M$ and $\bar{M}$. This would provide an additional hint for a grand-unified structure. However, because the light triplets have no triplet-matter couplings would typically be generated, and in particular, non-zero mixing of the third generation with the first two. Still, this unnatural scenario will allow us to obtain a useful bound for phenomenological purposes, because even in this most favorable case, the resulting bound is very strong.

We assume then that the only nonzero $T$ couplings are $y_{33}^{T qq}$ and $y_{33}^{T u\bar{d}}$. Upon rotating to the mass basis, tree level $T$ couplings to first generation fields will typically be generated. These are however model dependent. For example, if the up mass matrix is diagonal, such terms do not appear.

One can still obtain a model-independent bound from loop diagrams, such as the two-loop diagram of Fig. 1. Assuming that, in a basis with well-defined $U_T$ charges, the up matrix is diagonal, the $W$-loop is proportional to

$$\frac{\alpha_w}{4\pi} V_{td}^* \sum_i V_{ti}^* V_{ui} \frac{m_i^2}{m_t - m_W} \ln \left( \frac{m_i^2}{m_W^2} \right) \sim 10^{-10},$$

where $V$ is the CKM matrix. The charged Higgs loop scales as

$$\frac{\alpha_w}{4\pi} V_{td}^* \frac{m_t}{m_W} \sim \left\{ \begin{array}{ll} 10^{-6} & l = \tau \\ 10^{-10} & l = e \end{array} \right..$$

Thus we see that single $T$ exchange, even in the presence of triplet couplings to third-generation quarks only, implies a lower bound on the triplet mass between $10^6$ and $10^9$ GeV, depending on whether the lepton belonging to the same 10 as the top is the electron or the tau. It is therefore unlikely that the triplets can be detected in any foreseeable experiment.

In fact, as was mentioned above, in any natural model one would expect to have some other nonzero triplet couplings besides the ones we considered here, so that proton decay would occur already at tree level. If the model is based on a symmetry which is broken by some small parameter, it is reasonable to expect the tree contribution all other $T$ and $\bar{T}$ couplings vanish at the tree-level. We emphasize that such an assumption is unnatural. Terms that vanish at tree level would appear at one loop as we demonstrate shortly. Moreover, it is probably impossible to forbid all $T$ couplings in [1] for $i, j \neq 3$, while generating acceptable fermion masses, and in particular, non-zero mixing of the third generation with the first two. Still, this unnatural scenario will allow us to obtain a useful bound for phenomenological purposes, because even in this most favorable case, the resulting bound is very strong.

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and the loop contribution to involve the same parametric suppression. Thus, in any natural model, the bound on the triplet mass would be about one loop factor, or two orders of magnitude, stronger than what we found above.

Indeed, the most promising models from the point of view of getting small triplet couplings to the first generations, are models with the Higgses in the first SU(5), one 10 in the first SU(5) (so that the top Yukawa is renormalizable), and all other matter fields in the second SU(5), but we have been unable to find any model of this type for which proton decay would allow triplets below, roughly, $10^{11}\text{GeV}$.

V. TRIPLETS AT $10^{14}\text{GeV}$.

We will now use “Yukawa splitting” to obtain a model which reconciles the conflicting requirements from proton decay and coupling unification on the triplet mass. Many models of this type can be constructed, but we will present one simple example.

We take the Higgses to be $h = (5, 1, 0)$ and $\tilde{h} = (1, 5, 0)$, and the matter fields to be three copies of $(10, 1, 0)$ and $(5, 1, 0)$ with $q = \frac{\sqrt{5}}{2} (-3/2)$ for $Z_N (U(1))^4$. To cancel anomalies, we also add the fields $M \sim (5, 1, 1)$ and $M \sim (1, 5, 1)$.

As in the usual case, the up-type Yukawas originate from the superpotential terms $\tilde{h}(10, 1)(10, 1)$, so that $\langle y \rangle T q = y T e = y U$. On the other hand, down-type Yukawas must come from non-renormalizable terms involving a bifundamental,

$$y_{ij}^D \frac{1}{M_{\text{Planck}}} \tilde{\Phi}_2 \tilde{h}(10, 1)_i (5, 1)_j ,$$

so that all down Yukawas are uniformly suppressed by $\langle \tilde{\Phi}_2 \rangle / M_{\text{Planck}} \sim 10^{-2}$. The superpotential does not give rise to a $T$-matter coupling, since that requires a $\tilde{\Phi}_3$ instead of $\tilde{\Phi}_2$ in eqn. (6).

The dominant triplet contribution to proton decay now comes from single $T$ exchange. This is very similar to the usual $X$ and $Y$ gauge-boson mediated decay, but is much smaller, since the couplings involved are the Yukawas rather than gauge couplings. We then find that the bound on the triplet mass implied by proton decay is roughly $10^{12}\text{GeV}$.

The triplet symmetry so far forbids a triplet mass as well as a $\mu$ term. In order to generate an acceptable triplet mass, we assume that the symmetry is broken by some small parameter $\eta$ (which arises, say, as $\langle \tilde{\Phi}_2 \rangle / M_{\text{Planck}}$ with $S$ a fundamental or composite gauge singlet). For $\eta$ of charge +1, with $\eta \sim 10^{-2}$, the superpotential term

$$\eta h \tilde{\Phi}_3 \tilde{h} ,$$

is allowed and gives a $10^{14}\text{GeV}$ triplet mass. Coupling unification is then recovered, while proton decay is below experimental bounds.

The remaining aspect of “doublet-triplet” splitting, the $\mu$ problem, is readily solved as well. Suppose first that the triplet symmetry is a $U(1)$. Then, once we allow the coupling $\eta \sim 10^{-2}$, the $\mu$ term is automatically forbidden by the triplet symmetry and holomorphy. One can alternatively generate an acceptable $\mu$ term by taking the triplet symmetry to be discrete $U(1)$. For example, with $U_T = Z_{17}$ we have $M_T \sim 10^{14} \text{GeV}$ as before and $\mu \sim \eta M_{\text{GUT}} \sim 10^2 \text{GeV}$.

Since the triplet symmetry is now broken, $T$ couplings to matter may in principle be generated from terms such as $\eta^p \tilde{\Phi}_3 h(10, 1)(5, 1)$. One can check that, in the examples we consider, these terms are either forbidden, or are very suppressed, so that the dominant triplet contribution to proton decay is still from single $T$ exchange as described above.

One could also worry that, once the triplet symmetry is broken, the pattern of VEVs is no longer protected, and some triplet couplings are re-introduced from terms already present in the superpotential, such as $\tilde{\Phi}_3$. However, the pattern can be guaranteed using additional symmetries. One possibility is an $R$ symmetry under which the bifundamentals and $\eta$ have zero charges, and so cannot appear alone in the superpotential. This $R$ symmetry was present anyway in the models of section II in order to ensure a flat potential for the bifundamentals.

Finally, we comment on the fields $M$ and $M$. Just as the analogous fields $M$ and $M$ in the model of section II their mass is protected by the triplet symmetry and supersymmetry. Alternatively, their masses can be forbidden by the $R$-symmetry discussed above. In any case, once supersymmetry is broken, these fields will get mass near the supersymmetry-breaking scale.

To summarize, this model avoids the mismatch of coupling unification by having the triplets near $10^{14} \text{GeV}$, and gives proton decay below current experimental bounds. At low energies, there are new fields in complete $SU(5)$ representations, giving additional evidence for grand unification.

VI. CONCLUDING REMARKS.

In the examples we discussed, all matter fields transform under a single $SU(5)$. More generally, if either the 10s or the 5s are split between the different $SU(5)$ factors, the triplet symmetry will necessarily behave as a horizontal symmetry: Some matter fields will have generation-dependent $U_T$ charges. Clearly, these models give interesting patterns of fermion masses, which can significantly differ from the usual GUT relations.

We have assumed that there are no flavor-violating contributions from the superpartner sector. This holds for example for universal scalar masses.

Finally, we stress that, even in minimal GUT models,
the triplet contribution to proton decay is highly sensitive to the details of the Yukawa couplings. For example, if the up-quark mass vanishes, so that, in the up-mass basis, \( y_4^{11} = 0 \), there is no dimension-five triplet contribution to proton decay.

To summarize, we have shown that triplets below the GUT scale can be compatible with proton stability. If all triplet-matter couplings are suppressed, the triplets can be near the weak scale, and thus directly detectable. This would provide additional evidence for grand unification, which would otherwise be very hard to establish conclusively, even if proton decay is observed one day.

If the triplets only have large couplings to the top, their mass must be at least \( 10^6 \text{GeV} \). It would therefore be impossible to detect them. However, it is still interesting to consider triplets near \( 10^{14} \text{GeV} \), since then they supply just the right contributions to the running so that MSSM gauge couplings do unify. We present simple models that realize this possibility, with no other thresholds below the GUT scale.

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