Suppression of Proton Decay in the Missing-Partner Model for Supersymmetric SU(5) GUT

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Abstract

We construct a missing-partner model for supersymmetric SU(5) GUT assuming the Peccei-Quinn symmetry, in which the SU(5) gauge coupling constant remains in the perturbative regime below the gravitational scale \( \sim 2.4 \times 10^{18}\text{GeV} \). The Peccei-Quinn symmetry suppresses the dangerous dimension-five operators for the nucleon decay much below the limit from the present proton-decay experiments. We also stress that due to this suppression mechanism our model can accommodate even the large \( \tan \beta_H \) (\( \sim 60 \)) scenario which has been recently suggested to explain the observed value of the \( m_b/m_\tau \) ratio.

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Grand unified theories (GUT’s) of strong and electroweak interactions are based on the assumption of a large hierarchy between two mass scales, $M_{GUT} \sim 10^{16}$GeV and $M_{weak} \sim 10^2$GeV. In constructing a realistic GUT model, a serious problem arises from a phenomenological requirement that the SU(2)$_L$-doublet Higgs $H_f$ must have a mass of order of the electroweak scale to cause the breaking of SU(2)$_L \times U(1)_Y$ symmetry, while the mass of the color-triplet Higgs $H_c$ should be at the GUT scale in order to guarantee the observed stability of proton. This problem is not easily solved even in the supersymmetric (SUSY) extension of GUT’s, since an extremely precise adjustment of parameters in the superpotential is required to achieve such a large mass splitting of the doublet- and triplet-Higgs multiplets. Although such a tree-level hierarchy survives quantum corrections, thanks to the cancellation in the SUSY theories, the required fine-tuning of parameters seems very much unlikely.

The missing-partner model for the SUSY SU(5) GUT is a well-known example in which the above doublet-triplet splitting is naturally obtained without fine-tuning of parameters. However, this model becomes strongly interacting below the gravitational scale $M = M_{pl}/\sqrt{8\pi} \simeq 2.4 \times 10^{18}$GeV, since it contains somewhat high-rank representations of SU(5) for Higgs multiplets, 75, 50, and 500. Therefore, the perturbative description of GUT’s is broken down before reaching supergravity. A solution to this problem may be given if one puts the mass of 50 + 500 at the gravitational scale $M$. In this case, however, the color-triplet Higgses $H_c$ and $\overline{H}_c$ have a relatively smaller mass $M_{H_c} \sim M_{GUT}^2/M \sim 10^{(14-15)}$GeV which has been already excluded by the present proton-decay experiments.

In this letter we show that introduction of the Peccei-Quinn symmetry solves all problems mentioned above in the missing-partner model. That is, the Peccei-Quinn symmetry suppresses the $H_c$-mediated dimension-five ($D=5$) operators for the nucleon decay much well below the present experimental limit even for the relatively light color-triplet Higgses ($M_{H_c} \sim 10^{(14-15)}$GeV). We stress that this...
suppression also allows us to have a large \( \tan \beta_H \sim 60 \) which has been suggested as one of the parameter regions giving the correct \( m_b/m_\tau \) ratio \[1\].

The original missing-partner model \[4\] in the SUSY-SU(5) GUT consists of the following chiral supermultiplets;

\[
\psi_i(10), \quad \phi_i(\overline{5}), \quad H(5), \quad \overline{H}(\overline{5}), \quad \theta(50), \quad \theta(\overline{50}), \quad \Sigma(75),
\]

where \( i(=1–3) \) represents family index. In this model we incorporate a Peccei-Quinn symmetry \( U(1)_{PQ} \) under which the chiral multiplets in Eq. (1) transform as

\[
\begin{align*}
\psi_i(10) &\rightarrow e^{i\alpha/2} \psi_i(10), \\
\phi_i(\overline{5}) &\rightarrow e^{i\beta/2} \phi_i(\overline{5}), \\
H(5) &\rightarrow e^{-i\alpha} H(5), \\
\overline{H}(\overline{5}) &\rightarrow e^{-i\alpha + \beta/2} \overline{H}(\overline{5}), \\
\theta(50) &\rightarrow e^{i\alpha} \theta(50), \\
\theta(\overline{50}) &\rightarrow e^{i\alpha + \beta/2} \theta(\overline{50}), \\
\Sigma(75) &\rightarrow \Sigma(75),
\end{align*}
\]

with \( 3\alpha + \beta \neq 0 \). These \( U(1)_{PQ} \) charges are chosen such that the following superpotential is allowed,

\[
W = \frac{1}{4} h_{ij} \psi_i^{(AB)} \psi_j^{(CD)} H^E \epsilon_{ABCDDE} + \sqrt{2} f_{ij} \psi_i^{(AB)} \phi_A \overline{\phi}_B \\
+ G_H H^A \Sigma_{(FG)}^{(BC)} \theta^{(DE)(FG)} \epsilon_{ABCDDE} + G_{ABCDDE} \theta^{(DE)(FG)} \epsilon_{ABCDDE} \\
+ M_{75} \Sigma^{(AB)}_{(CD)} \Sigma^{(CD)}_{(AB)} - \frac{1}{3} \lambda_{75} \Sigma^{(AB)}_{(EF)} \Sigma^{(CD)}_{(AB)} \Sigma^{(CD)}_{(EF)},
\]

and that the Peccei-Quinn symmetry is not broken by \( \langle \Sigma \rangle \neq 0 \). Here, indices \( A, B, C \ldots \) are the SU(5) indices which run from 1 to 5, \( \epsilon_{ABCDDE} \) and \( \epsilon^{ABCDDE} \) are the fifth-antisymmetric tensors, and the indices in \( (AB) \) are antisymmetric.

The above model is still incomplete, since \( \theta(50) \) and \( \theta(\overline{50}) \) can not have an invariant mass term. To give large masses to them we double the \( H \) and \( \theta \) sector

3This is required for the Peccei-Quinn mechanism to work. On the other hand, the baryon-number violating \( D = 5 \) operators, \( \phi_i \psi_j \psi_k \psi_l \), are forbidden under this condition, irrespective of their origins.
introducing a new set of chiral multiplets, \( H'(5), \overline{H}'(\overline{5}), \theta'(50), \) and \( \theta'(\overline{50}) \), which have opposite \( U(1)_{PQ} \) charges of the corresponding original fields, \( H, \overline{H}, \theta, \) and \( \overline{\theta} \). We add a new superpotential to Eq. (3),

\[
W' = G'_H H'^A \Sigma_{(BC)}^{(FG)} \theta'^{(DE)(FG)} \epsilon_{ABCDE} + G_{\overline{H}} \overline{H}' A \Sigma_{(BC)}^{(FG)} \overline{\theta}'^{(DE)(FG)} \epsilon_{ABCDE} + M_1 \overline{\theta}(AB)(CD) \theta^{(AB)(CD)} + M_2 \overline{\theta}'(AB)(CD). \tag{4}
\]

To avoid that the \( SU(5) \) gauge coupling constant blows up below the gravitational scale \( M \), we assume

\[
M_1, M_2 \gtrsim 10^{18}\text{GeV}. \tag{5}
\]

In this letter, we take \( M_1 = M_2 = M(\equiv 2.4 \times 10^{18}\text{GeV}) \) for simplicity. Then, we have four Higgses, \( H, \overline{H}, H', \overline{H}' \), and one Higgs \( \Sigma \) much below the gravitational scale \( M \).

The 75-dimension Higgs \( \Sigma \) has the following vacuum-expectation value that causes the breaking \( SU(5) \to SU(3)_C \times SU(2)_L \times U(1)_Y \),

\[
\langle \Sigma \rangle^{(a\beta)}_{(\gamma\delta)} = \frac{1}{2} \left( \delta^a_\gamma \delta^\beta_\delta - \delta^a_\delta \delta^\beta_\gamma \right) V_\Sigma,
\]

\[
\langle \Sigma \rangle^{(ab)}_{(cd)} = \frac{3}{2} \left( \delta^a_c \delta^b_d - \delta^a_d \delta^b_c \right) V_\Sigma, \tag{6}
\]

\[
\langle \Sigma \rangle^{(aa)}_{(bb)} = -\frac{1}{2} \left( \delta^a_b \delta^b_a \right) V_\Sigma,
\]

where

\[
V_\Sigma = \frac{3}{2} \frac{M_{75}}{\lambda_{75}} \tag{7}
\]

obtained from the superpotential Eq. (3). Here, \( \alpha, \beta \ldots \) are the \( SU(3)_C \) indices and \( a, b \ldots \) the \( SU(2)_L \) indices. This vacuum-expectation value generates masses for the color-triplet Higgses as (after integrating out the heavy fields, \( \theta, \overline{\theta} \) and \( \theta', \overline{\theta}' \)),

\[
M_{H_c} H'^c \overline{H}'_{\alpha \alpha} + M_{\overline{H}_c} H'^{\alpha a} \overline{H}_{ca}, \tag{8}
\]

with

\[
M_{H_c} \approx 48 G_H G'_H V_\Sigma^2 M, \quad M_{\overline{H}_c} \approx 48 G_H G'_H V_\Sigma^2 M. \tag{9}
\]
The four SU(2)$_L$-doublet Higgses, $H_f$, $H'_f$, and $H''_f$, remain massless.

In order to break the Peccei-Quinn symmetry, we introduce a pair of SU(5)-
singlet chiral multiplets $P$ and $Q$ whose U(1)$_{PQ}$ charges are chosen as $P \rightarrow e^{-i\frac{1}{2}(3\alpha+\beta)}P$ and $Q \rightarrow e^{i\frac{1}{2}(3\alpha+\beta)}Q$ so that the following superpotential is allowed

$$W'' = \frac{f}{M} P^3 Q + g_P H'_A H'^A P.$$  \hfill (10)

We have a very flat scalar potential for $P$ and $Q$ as

$$V(P, Q) = \frac{f^2}{M^2} |P|^6 + \frac{f^2}{M^2} |3P^2Q|^2.$$ \hfill (11)

As pointed out in Ref. [12], the introduction of negative soft-SUSY breaking mass $\sim -m^2$ for $P$ induces very naturally the Peccei-Quinn symmetry breaking at the intermediate scale,

$$\langle P \rangle \simeq \langle Q \rangle \simeq \sqrt{\frac{Mm}{f}} \sim 10^{11}\text{GeV},$$ \hfill (13)

provided $m \sim 1\text{TeV}$ and $f \sim 1$. This Peccei-Quinn symmetry breaking produces an intermediate-scale mass for a pair of light SU(2)$_L$-doublet Higgses, $H'_f$ and $H''_f$,

$$M_{H'_f} = g_P \langle P \rangle,$$ \hfill (15)

through the Yukawa interaction in Eq. (10).

$^4$This negative soft SUSY breaking mass may be induced by radiative corrections from the $j_{ij} N_i N_j P$ interactions given in Eq. (16). See Ref. [12] for details.

$^5$To generate non-vanishing vacuum-expectation value for $Q$, we have assumed the other soft SUSY-breaking term $\sim (m/M)P^3Q$ (see Ref. [12] in detail). Here, $Q$ denotes the scalar component of the chiral multiplet $Q$.

$^6$The breaking scale of the Peccei-Quinn symmetry is bounded by astrophysics and cosmology as

$$10^{10}\text{GeV} \leq \langle P \rangle, \langle Q \rangle \leq 10^{13}\text{GeV}.$$ \hfill (12)

$^7$With this charge assignment for $P$ and $Q$, the Higgses $H_f$ and $H''_f$ receive a mass only from the following superpotential,

$$\frac{P^2 Q}{M^2} H.$$ \hfill (14)

This gives a small invariant mass $\mu$ for $H_f$ and $H''_f$ as $\mu \sim 1\text{GeV}$ for $\langle P \rangle \sim \langle Q \rangle \sim 10^{12}\text{GeV}$, which is not excluded for $\tan \beta_H < \sqrt{2}$ [12]. However, if one takes $f \sim 10^{-4}$, then $\langle P \rangle$ and $\langle Q \rangle$ are $\sim 10^{13}\text{GeV}$. In this case the invariant mass $\mu$ becomes $O(1)\text{TeV}$. 


So far we have two independent charges $\alpha$ and $\beta$ defined in Eq. (2) and hence there are two global $U(1)$'s. To eliminate one of them, we introduce right-handed neutrino multiplets $N_i$ (1) [15]. In fact, with two possible Yukawa couplings

$$W'' = k_{ij}N_i\phi_jH + j_{ij}N_iN_jP,$$

we have only one $U(1)_{PQ}$ and the charge $\alpha$ is fixed as $\alpha = 3\beta$. The Yukawa couplings $j_{ij}N_iN_jP$ in Eq. (16) induce Majorana masses for the right-handed neutrino multiplets $N_i$, with $\langle P \rangle \neq 0$. Interesting is that the Majorana masses for the right-handed neutrinos are expected to be $O(10^{11})$ GeV, which naturally induce very small masses of neutrinos through the cerebrated see-saw mechanism [16] in a range of the MSW solution [17] to the solar neutrino problem.

The color-triplet Higgses have an off-diagonal element in their mass matrix as

$$
\left( \begin{array}{c}
H_c, \bar{H}_c \end{array} \right) \left( \begin{array}{cc}
M_{H_c} & 0 \\
g_P \langle P \rangle & M_{H_c} \end{array} \right) \left( \begin{array}{c}
H'_c \\
H_c \end{array} \right).
$$

The baryon-number violating $D = 5$ operators [10] mediated by the color-triplet Higgses are given by in the present model (see Fig. 1 (a))

$$
g_P \langle P \rangle \frac{1}{M_{H_c} M_{\bar{H}_c}} 2\sqrt{2} f_{ijkl} \phi_{F i}^{(FA)} \psi_j^{(BC)} \psi_k^{(DE)} \epsilon_{ABCDE}.
$$

Notice that those in the minimum SUSY-SU(5) GUT are given by (see Fig. 1 (b))

$$
\frac{1}{M_{H_c}} \frac{1}{2\sqrt{2}} f_{ijkl} \phi_{F i}^{(FA)} \psi_j^{(BC)} \psi_k^{(DE)} \epsilon_{ABCDE}.
$$

Thus we easily see that the $D = 5$ operators in the present model are more suppressed by a factor $M_{H'_c}/M_{H_c}$ compared with in the minimum SUSY-SU(5) GUT.

We are now at the point to show a crucial difference between our model and the previous Pececi-Quinn extension [9] of the minimum SUSY-SU(5) GUT. The mass spectrum above $U(1)_{PQ}$ breaking scale contains four $SU(2)_L$-doublets Higgses and hence the success of the gauge coupling unification in the minimum SUSY-GUT may be lost in general. Thus, we have non-trivial constraints on the Higgs masses from the gauge coupling unification, which are different from the minimum SUSY-SU(5) GUT. The crucial point is that these constraints in the present
model are much different from those obtained in the previous model \[\text{[4]}\], since the 
SU(3)_C \times SU(2)_L \times U(1)_{Y} components of Σ(75) have different masses for each other 
due to the own vacuum-expectation value. This mass splitting gives large threshold 
corrections to the SU(3)_C \times SU(2)_L \times U(1)_{Y} gauge coupling constants.

The running of the three gauge coupling constants at the one-loop level is given 
by the following solutions to the renormalization group equations \[\text{[18]}\],

\[\begin{align*}
\alpha_3^{-1}(m_Z) &= \alpha_5^{-1}(\Lambda) + \frac{1}{2\pi} \left\{ \left( -2 - \frac{2}{3} N_g \right) \ln \frac{m_{SUSY}}{m_Z} \\
&\quad + (-9 + 2N_g) \ln \frac{\Lambda}{m_Z} - 4 \ln \frac{\Lambda}{M_V} \\
&\quad + 9 \ln \frac{\Lambda}{M_\Sigma} + \ln \frac{\Lambda}{0.8M_\Sigma} + 10 \ln \frac{\Lambda}{0.4M_\Sigma} + 3 \ln \frac{\Lambda}{0.2M_\Sigma} \\
&\quad + \ln \frac{\Lambda}{M_{H_c}} + \ln \frac{\Lambda}{M_{H_c'}} \right\},
\end{align*}\]

(20)

\[\begin{align*}
\alpha_2^{-1}(m_Z) &= \alpha_5^{-1}(\Lambda) + \frac{1}{2\pi} \left\{ \left( -\frac{4}{3} - \frac{2}{3} N_g - \frac{5}{6} \right) \ln \frac{m_{SUSY}}{m_Z} \\
&\quad + (-6 + 2N_g + 1) \ln \frac{\Lambda}{m_Z} - 6 \ln \frac{\Lambda}{M_V} \\
&\quad + 16 \ln \frac{\Lambda}{M_\Sigma} + 6 \ln \frac{\Lambda}{0.4M_\Sigma} \\
&\quad + \ln \frac{\Lambda}{M_{H_c'}} \right\},
\end{align*}\]

(21)

\[\begin{align*}
\alpha_1^{-1}(m_Z) &= \alpha_5^{-1}(\Lambda) + \frac{1}{2\pi} \left\{ \left( -\frac{2}{3} N_g - \frac{1}{2} \right) \ln \frac{m_{SUSY}}{m_Z} \\
&\quad + \left( 2N_g + \frac{3}{5} \right) \ln \frac{\Lambda}{m_Z} - 10 \ln \frac{\Lambda}{M_V} \\
&\quad + 10 \ln \frac{\Lambda}{0.8M_\Sigma} + 10 \ln \frac{\Lambda}{0.4M_\Sigma} \\
&\quad + 2 \ln \frac{\Lambda}{M_{H_c}} + \frac{2}{5} \ln \frac{\Lambda}{M_{H_c'}} + \frac{3}{5} \ln \frac{\Lambda}{M_{H_c'}} \right\},
\end{align*}\]

(22)

where \(\alpha_5 \equiv g_5^2/4\pi\) is the SU(5) gauge coupling constant, \(M_V\) the heavy gauge 
boson mass (\(M_V = 2\sqrt{15} g_5 V_\Sigma\)), and \(\Lambda\) the renormalization point which is taken 
\(\Lambda \gg M_{GUT}\). Here, we have assumed that all superparticles in the SUSY-standard 
model have a SUSY-breaking common mass \(m_{SUSY}\) for simplicity, and the mass 
splitting of Σ(75) has been included. The each SU(3)_C × SU(2)_L × U(1)_Y components
of $\Sigma$ have the following masses;

\begin{align*}
(SU(3)_C \times SU(2)_L \times U(1)_Y) \quad & \text{mass} \\
(8, 3, 0) \quad & M_\Sigma \\
(3, 1, \frac{5}{3}), (\bar{3}, 1, -\frac{5}{3}) \quad & \frac{4}{5}M_\Sigma \\
(6, 2, \frac{5}{6}), (\bar{6}, 2, -\frac{5}{6}) \quad & \frac{2}{5}M_\Sigma \\
(1, 1, 0) \quad & \frac{2}{5}M_\Sigma \\
(8, 1, 0) \quad & \frac{1}{5}M_\Sigma \\
(3, 2, -\frac{5}{6}), (\bar{3}, 2, \frac{5}{6}) \quad & 0 \text{ (Nambu-Goldstone multiplets)}
\end{align*}

where $M_\Sigma = 10\lambda_7 V_\Sigma/3$. By eliminating $\alpha^{-1}_5$ from Eqs. (20-22), we obtain simple relations [5, 9]:

\begin{align*}
(3\alpha_2^{-1} - 2\alpha_3^{-1} - \alpha_1^{-1})(m_Z) &= \frac{1}{2\pi} \left\{ \frac{12}{5} \ln \frac{M_{H_\mu}}{M_{H_\mu}} - 2 \ln \frac{m_{SUSY}}{m_Z} \\
& \quad - \frac{12}{5} \ln(1.7 \times 10^4) \right\}, \\
(5\alpha_1^{-1} - 3\alpha_2^{-1} - 2\alpha_3^{-1})(m_Z) &= \frac{1}{2\pi} \left\{ 12 \ln \frac{M_{V_\mu}^2}{m_Z^3} + 8 \ln \frac{m_{SUSY}}{m_Z} \\
& \quad + 36 \ln(1.4) \right\}.
\end{align*}

Notice that the last terms in Eqs. (24,25) come from the mass splitting of $\Sigma(75)$, which makes a crucial difference between the previous and the present models [5, 9].

To perform a quantitative analysis, we use the two-loop renormalization group equations between the weak and the GUT scales. Instead of the common mass $m_{SUSY}$ of superparticles we have used the mass spectrum estimated from the minimum supergravity [5, 19] to calculate the one-loop threshold correction at the SUSY-breaking scale. Using the experimental data $\alpha^{-1}(m_Z) = 127.9 \pm 0.2$, $\sin^2 \theta_W(m_Z) =$

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8The Nambu-Goldstone multiplets are absorbed to the gauge multiplets forming massive vector multiplets $V$ at the GUT scale.

9The mass splitting of $\Sigma(24)$ in the minimum SUSY-SU(5) GUT and in its Peccei-Quinn extension does not produces these constant terms as pointed out in Ref. [5, 9].
\[0.2326 \pm 0.0008, \text{and } \alpha_3(m_Z) = 0.118 \pm 0.007\] we obtain
\[
3.7 \times 10^{17} \text{ GeV} \leq \frac{M_{H_c}M_{\Sigma}}{M_{H'}f} \leq 3.8 \times 10^{21} \text{ GeV, (26)}
\]
\[
6.8 \times 10^{15} \text{ GeV} \leq (M_VM_{\Sigma})^{1/3} \leq 2.4 \times 10^{16} \text{ GeV. (27)}
\]

This should be compared with the previous result in Ref. [9, 5],
\[
2.2 \times 10^{13} \text{ GeV} \leq \frac{M_{H_c}M_{\Sigma}}{M_{H'}f} \leq 2.3 \times 10^{17} \text{ GeV, (28)}
\]
\[
9.5 \times 10^{15} \text{ GeV} \leq (M_VM_{\Sigma})^{1/3} \leq 3.3 \times 10^{16} \text{ GeV. (29)}
\]

The main reason for the different results comes from the presence of the constant terms in Eqs. (24, 25) which originate from the mass splitting of \(\Sigma(75)\). Notice that Eq. (26) suggests \(M_{H_c} \sim M_{H'} \sim 10^{(13-16)}\text{GeV}\) for \(M_{H'}f \approx 10^{10}\text{GeV}\). This is very much consistent with Eq. (8) with \(V_{\Sigma} \approx 10^{(15-16)}\text{GeV}\) and \(G_{H^{'}}G_{H^{'}} \sim 1.\)

The \(D=5\) operator in the minimum SUSY-SU(5) model is proportional to \(1/M_{H_c}\) as shown in Eq. (19). The detailed analysis on the nucleon-decay experiments gives the lower limit on the color-triplet Higgs mass as \(M_{H_c} \geq 5 \times 10^{15}\text{GeV}\) in the minimum model. On the other hand in the present model one can easily estimate the nucleon decay rate by replacing \(1/M_{H_c}\) by \(M_{H'}f/M_{H_c} \leq M_{H_c} \leq 5 \times 10^{20}\text{GeV}\) (see Eq. (26)). This mass ratio is nothing but one in Eq. (26) derived from the requirement of the gauge-coupling unification. We find that the constraint Eq. (26) is much weaker than that in Eq. (28) and hence this model is still consistent with the lower limit on the nucleon lifetime even for the case of large \(\tan \beta_H \equiv \langle H_f \rangle/\langle \overline{H}_f \rangle^\circ\)

\[10\] If one uses the recent experimental data \(\alpha^{-1}(m_Z) = 127.9 \pm 0.2, \sin^2 \theta_W(m_Z) = 0.2314 \pm 0.0004, \text{and } \alpha_3(m_Z) = 0.118 \pm 0.007\), one gets
\[
1.4 \times 10^{17} \text{ GeV} \leq \frac{M_{H_c}M_{\Sigma}}{M_{H'}f} \leq 5.5 \times 10^{20} \text{ GeV,}
\]
\[
8.4 \times 10^{15} \text{ GeV} \leq (M_VM_{\Sigma})^{1/3} \leq 2.6 \times 10^{16} \text{ GeV.}
\]

However, we use the old data in the text for a comparison with the previous result.

\[11\] Comparing the upper limits on \(M_{H_c}M_{\Sigma}/M_{H'}f\) in Eq. (26) and in Eq. (28), we see that the \(D=5\) operators in our model can be suppressed by a factor \(\sim 10^{-4}\) compared with in the previous model.

\[12\] When \(\tan \beta_H\) is large, the \(D=5\) operators for the nucleon decay are proportional to \(\tan \beta_H\).
We show the evolution of the SU(3)$_C \times$SU(2)$_L \times$U(1)$_Y$ and SU(5) gauge coupling constants in Fig. 2 taking $M_{H_c} = M_{\tilde{H}_c} = 10^{15}$GeV and $M_{H'_{\tau}} = 10^{10}$GeV for a demonstrational purpose. We see that the unification of three gauge coupling constants occurs around $10^{16}$GeV and the SU(5) gauge coupling constant stays in the perturbative regime below the gravitational scale $M \simeq 2.4 \times 10^{18}$GeV.

In Fig. 3 we show the constraint from the present nucleon-decay experiments, taking the same parameters for $M_{H_c}$ and $M_{H'_{\tau}}$ as above. To demonstrate how safe our model is, we have chosen even the large tan $\beta_H = 60$ and the largest hadron matrix element $\beta = 0.03$GeV$^3$ (see Ref. [5] for notations). One can see that the superparticle masses below 1TeV are still allowed in the present model. Thus, we stress that the intriguing idea of the Yukawa coupling unification $h_t(M_{GUT}) = h_b(M_{GUT}) = h_{\tau}(M_{GUT})$ [22] (which implies the large tan $\beta_H \simeq 50 - 60$) is consistent not only with the observed fermion masses, $m_b = (4.2 - 4.4)$GeV [11] and $m_t = (160 - 190)$GeV [23], but also with the present lower limit on the nucleon lifetime.

In this letter we have shown that the Peccei-Quinn extension of the missing-partner model in SUSY-SU(5) GUT is consistent with the observed stability of the proton, even if the masses of the unwanted $50 + \bar{50}$ are lifted up to the gravitational scale so that the SU(5) gauge coupling constant remains small enough for the perturbative description of GUT’s. We believe that the present model is worth being pursued, since it is only a known, perturbative SUSY-GUT model which is phenomenologically consistent and naturally realizes the doublet-triplet splitting of Higgs multiplets.
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Figure Captions

Fig. 1
(a) The Feynman diagram of the baryon-number violating $D = 5$ operators in the present model. (b) The corresponding Feynman diagram in the minimum SUSY-SU(5) GUT.

Fig. 2
The flows of the running gauge coupling constants of SU(3)$_C$×SU(2)$_L$×U(1)$_Y$ and SU(5). Here, $M_{H_c}$ and $M_{H'_c}$ are taken at $10^{15}$GeV, and $M_{H'_f}$ at $10^{10}$GeV. We use $\alpha^{-1}(m_Z) = 127.9 \pm 0.2$, $\sin^2 \theta_W(m_Z) = 0.2326 \pm 0.0008$, and $\alpha_3(m_Z) = 0.118 \pm 0.007$ for the initial condition. We assume the SUSY-breaking scale $\sim 1$TeV.

Fig. 3
The lower bound of the superparticle masses from the negative search of the nucleon decay [7]. The horizontal axis is the wino mass and the vertical line the sfermion mass. Here, $M_{H_c}$ and $M_{H'_c}$ are taken at $10^{15}$GeV, and $M_{H'_f}$ at $10^{10}$GeV. We take $\tan \beta_H = 60$ for the ratio of the vacuum-expectation values, $\langle H_f \rangle / \langle \overline{H}_f \rangle$, and $\beta = 0.03$GeV$^3$ for the hadron matrix element. We also show the lower bound in the minimum SUSY-SU(5) GUT by the dashed line, taking $10^{17}$GeV for $M_{H_c}$ and the same values for the other parameters. The dotted line is the lower bound of the chargino mass from LEP [24] and the dash-dotted line that of the squark mass from CDF [25].
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