BARYON DECAY: FLIPPED $SU(5)$ SURMOUNTS ANOTHER CHALLENGE

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ABSTRACT

Effective four-dimensional field theories derived from string may contain superpotential terms that violate baryon number. Terms that are cubic in the Standard Model particles are well known to lead to excessively rapid proton decay, and the same is true for terms that are quartic in these light fields. This holds even if the terms contain many powers of fields with large intermediate vacuum expectation values, corresponding to asymmetric manifolds of compactification. We show that no such B-violating superpotential interaction arises in flipped $SU(5)$, thanks to an extra symmetry of the effective cubic $d = 4$ and quartic $d = 5$ superpotential terms, induced at short distances. We conclude that in this model baryon decay proceeds predominantly through conventional $d = 6$ gauge boson exchange, and we recalculate its rate using the latest LEP data to estimate the heavy gauge boson masses.

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Our experimental colleagues' lack of success in finding baryon decay [1] makes life difficult for theorists trying to go beyond the Standard Model, and their difficulties grow the further beyond the Standard Model they try to go. In the basic Standard Model, it is impossible to write down a B-violating renormalizable interaction [2]. These appear in a Grand Unified Theory (GUT) [3], but are suppressed by the large masses of the gauge and Higgs bosons that violate baryon number. Minimal non-supersymmetric GUTs do not have large enough masses for the B-violating gauge bosons [4] but this particular shortcoming can be resolved by going supersymmetric [5]. However, even the minimal supersymmetric extension of the Standard Model could in principle admit renormalizable B- and L-violating cubic superpotential terms that, if present, would have caused baryons to decay around the epoch of cosmological nucleosynthesis. These dangerous interactions have to be forbidden by postulating an extra matter parity symmetry, which strikes many theorists as inelegant. Moreover, although supersymmetric GUTs have gauge boson masses sufficiently heavy for the $d = 6$ four-fermion B-violating operators they generate to give an acceptably long baryon lifetime, there are new $d = 5$ Higgs supermultiplet exchange operators [6] that give an unacceptably short baryon lifetime unless $m_H \gtrsim 10^{17}$ GeV [7]. This latter problem is avoided most elegantly in the flipped $SU(5)$ model [8], which has a missing partner mechanism that suppresses the Higgs supermultiplet mixing required to generate the dangerous $d = 5$ operators. However, even if these are not generated by Higgs exchanges, nonrenormalizable $d = 5$ superpotential terms could appear with coefficients $\mathcal{O}(1/M_{Pl})$, generated by dynamics at the Planck scale [9].

Many of the above difficulties are present in effective field theories derived from the string. For example, early in the days of string-inspired phenomenology, it was pointed out that matter parity was not automatic in most string-derived models, and that renormalizable trilinear $d = 4$ superpotential terms violating baryon number were present in generic string vacua [10]. These could often be avoided by going to a particular symmetric vacuum [11], but this vacuum would often have additional global symmetries rendering it unusable for phenomenology.

Then it could be tricky to control the absence of $d = 4$ B-violating interactions when one perturbed away from the symmetric point. Nonrenormalizable quartic B-violating superpotential terms [9] could also be expected to be generic features of string vacua, but were usually ignored, in part because the technology required to calculate them was lacking.

In previous papers we have shown how the flipped $SU(5)$ model derived from the string [12,13] avoided the dangerous $d = 4$ and Higgs-induced $d = 5$ operators, and we gave a first estimate of the lifetime for baryon decay induced by $d = 6$ gauge boson exchange operators [14]. However, we have not previously discussed possible $d \geq 5$ operators generated at the Planck scale [9], some of which could generate effective trilinear qqq and $qq\bar{l}$ $d = 4$ operators *via* large intermediate scale vacuum expectation values (*vevs*) corresponding to distortions of the flipped $SU(5)$ symmetric vacuum, and also effective quadrilinear $qqq\bar{l}$ $d = 5$ operators *via* an analogous mechanism.

In this paper, we first demonstrate that these Planck-scale $d = 5$ operators are indeed a very real danger, potentially violating the present experimental lower limit on the baryon lifetime by many orders of magnitude. Hence these and the prospective $d = 4$ operators must be controlled to very high orders in the large intermediate-scale *vevs*. Using recently-developed technology [15] for calculating nonrenormalizable superpotential terms in free world-sheet fermionic theories, we prove that the dangerous $qqq/qq\bar{l}$ and $qqq\bar{l}$ operators in fact *vanish to all orders in the large vevs* in the flipped $SU(5)$ model. This is because of a custodial symmetry beyond $SU(5) \times U(1)$ which arises when one calculates point-like higher-order nonrenormalizable interactions using the conformal field theory on the world-sheet. Since these results confirm our previous claims [14] that $d = 6$ gauge boson exchange is the dominant mechanism for baryon decay in the flipped $SU(5)$ model, we recalculate its rate using an estimate of the heavy gauge boson masses provided by an analysis of recent LEP data [16]. We find

$$\tau(p) \sim \tau(n) \sim 10^{31} \text{ to } 10^{37} \text{ y}, \quad (1)$$

with the dominant decay modes being $p \rightarrow e^+ \pi^0$, $\bar{\nu}_e \pi^+$ and $n \rightarrow e^+ \pi^-$, $\bar{\nu}_e \pi^0$, as predicted previously [14].

The catastrophic consequences [10] of trilinear $qqq/qq\bar{l}$ superpotential terms do not need any belaboring, but it is worth emphasizing how undesirable quartic $qqq\bar{l}$ superpotential terms [9] also are. We assume an $SU(5)$ GUT framework so that quarks and leptons appear in chiral $\bar{\mathbf{5}}$ (f^α) and $\mathbf{10}$ ($F_{\beta\gamma}$) superfields. The only B-violating quartic interactions are those contained in

$$\bar{f}_a^\alpha F_{\alpha\beta}^\dagger F_{\gamma\delta}^c F_{\gamma\delta}^d F_{\lambda\mu}^\dagger \epsilon^{\beta\gamma\delta\lambda\mu}, \quad (2)$$

where a, b, c, d are flavor indices. This term contains two types of $\Delta B = \pm 1$ interaction:

$$l_a q_b q_c q_d \quad \text{and} \quad q_a^c l_b^c q_c^c q_d^c. \quad (3)$$

Because of color antisymmetrization, the three quark generation indices cannot all be the same, and none of the $q_a^c l_b^c q_c^c q_d^c$ operators can contribute to baryon decay when dressed by gaugino exchange, whereas $lqqq$ can [17]. A complete catalog of such operators is given in Ref. [9]: the linearly-independent operators made of the two lightest generations are

$$\mathcal{L}_{eff} = \int d^2\theta \epsilon^{ijk} \sum_{l=\epsilon, \mu, \tau} g_{cnu}^l (l c_i u_j d_k^l - \nu_l s_i^l u_j d_k^l) + g_{ucc}^l (l u_i c_j s_k^l - \nu_l d_i^l c_j s_k^l), \quad (4)$$

where i, j, k are color indices and the primes denote Cabibbo-rotated charge $-1/3$ quarks. An operator of the second type with $l = \mu$ is present [17, 7] in the minimal $SU(5)$ supersymmetric GUT with a coefficient

$$g_{ucc}^\mu = 2\sqrt{2} K G_F m_c m_s \sin \theta_c / M_{H_u}, \quad (5)$$

where $K \equiv (v_1^2 + v_2^2)/v_1 v_2 : v_i \equiv \langle H_i^0 \rangle$. The dressing of (4) by $U(1) \tilde{B}$ and $SU(3) \tilde{g}$ gaugino exchanges give symmetric combinations of four-fermion operators that

vanish by a Fierz identity [9]. Dressing the ucc and cuu operators in (4) by $\tilde{W}^\pm / \tilde{H}^\pm$ exchange gives four-fermion operators proportional to

$$g_{ucc}^l : O_W = (us)_L (d\nu)_L + (ud)_L (s\nu)_L, \quad (6)$$

$$g_{cuu}^l : O'_W = (su)_L (ul - d\nu)_L + 2(du)_L (s\nu)_L, \quad (7)$$

respectively. Taking the first operator as an example, we see that $p \rightarrow \bar{\nu} K^+$ and $n \rightarrow \bar{\nu} K^0$ decays are preferred, and cannibalizing the calculations of Ref. [7] we estimate

$$\tau(B \rightarrow \bar{\nu} K) \approx 4 \times 10^{15} \text{ y} \times \frac{1}{(10^{16} \text{ GeV} \times g_{ucc}^l)^2}. \quad (8)$$

Taking $\tau(B \rightarrow \bar{\nu} K) > 10^{32} \text{ y}$ [1], we therefore infer that

$$g_{ucc}^l \lesssim 7 \times 10^{-25} \text{ GeV}^{-1}, \quad (9)$$

which is much smaller than $1/M_{Pl}$. Expressing this apparent conflict another way, for $g_{ucc}^l \approx 1/10^{18} \text{ GeV}$, which is the magnitude of a typical $d = 5$ nonrenormalizable superpotential term found in previous calculations [18], we estimate

$$\tau(B \rightarrow \bar{\nu} K) \approx 10^{19} \text{ to } 10^{20} \text{ y}. \quad (10)$$

A similar estimate would apply to the g_{cnu}^l operator in (4). We conclude that effective $d = 5$ nonrenormalizable superpotential terms are a real, major danger for string models.

We now turn to the analysis of these effective $d = 4, 5$ operators in the string-derived flipped $SU(5)$ model [12]. The relevant $SU(5) \times U(1)$ fields are:

$$F(\mathbf{10}, 1/2), \quad \bar{f}(\bar{\mathbf{5}}, -3/2), \quad \bar{F}(\mathbf{10}, -1/2), \quad l^c(\mathbf{1}, 5/2). \quad (11)$$

Of the four quartic invariants that can be formed with these fields, namely $FFF\bar{f}, F\bar{f}\bar{f}l^c, (F\bar{F})^2, \bar{f}l^c\bar{F}\bar{F}$, the last two clearly do not contain relevant $\Delta B \neq 0$

operators. Moreover, the only B-violating operator in the $SU(3) \times SU(2) \times U(1)$ decomposition of $F\bar{f}\bar{f}f^c$, namely $d^c u^c u^c l^c$, does not contribute to baryon decay [17], as discussed above. The remaining invariant, $FFF\bar{f}$, contains four pieces in its decomposition,

$$FFF\bar{f}: d^c d^c u^c N^c, QLd^c N^c, QQQL, QQd^c u^c, \quad (12)$$

where $Q(L)$ has the quantum numbers of a quark (lepton) doublet, and N^c those of the conjugate neutrino. If one of the F fields in (12) contains an $SU(5)$ breaking vev (N^c), then the first two terms would produce effective $d = 4$ operators and *together* mediate very rapid proton decay [10]. The third operator in (12) could *a priori* contain the $d = 5$ $qqql$ pieces, discussed above.

In the flipped $SU(5)$ model derived [12] from the string, the possible presence of such terms is constrained by the extra four $U(1)$ charges present in the model, whose assignments are given in Ref. [12]. The lightest generation of quarks and leptons is made out of the fields *F_3 , $\bar{f}_3$, and l_3^c , modulo small Cabibbo angle rotations. It is easy to check that, because of the special values these fields have of the third and fourth extra $U(1)$ hypercharges, any effective $FFF\bar{f}$ operator must contain an even number (0, 2, or 4) of the light fields $F_3, \bar{f}_3$, possibly multiplied by any number of *neus* of the gauge singlet fields that correspond to deformations of the $SU(5) \times U(1)$ model. We have indeed found fifth-order superpotential terms of the types

$$F_3 F_i F_i \bar{f}_3 \phi, \quad i = 2, 4. \quad (13)$$

These terms are consistent with all the gauge symmetries and the known rules for calculating nonzero nonrenormalizable superpotential terms, and give the

* In this model the quark and lepton fields are contained in [12, 18]: $F_i(10, 1/2), \bar{f}_i(5, -3/2)$, and $l_i^c(1, -5/2)$, $i = 2, 3, 4$. The second and third generation fields come from linear combinations of $i = 2, 4$.

following potential $d = 5$ $\Delta B \neq 0$ operators

$$Q_3 Q_i Q_i L_3, \quad i = 2, 4, \quad (14)$$

after the gauge-singlet fields ϕ acquire *vevs* of order 10^{17} GeV. These operators appear to pose a severe danger for flipped $SU(5)$.

Let us therefore look more carefully at these possible nonrenormalizable superpotential terms, which can be written explicitly as

$$c_{IJK} \epsilon^{\beta\gamma\delta\lambda\mu} (\bar{f})^\alpha (F_I)_{\alpha\beta} (F_J)_{\gamma\delta} (F_K)_{\lambda\mu}, \quad (15)$$

where the c_{IJK} are generation-dependent coefficients to be calculated. From this expression one can obtain the $SU(3) \times SU(2) \times U(1)$ components of interest as follows

$$d^c d^c u^c N^c: c_{IJK} \epsilon^{abd} u_a^c d_b^c (d_{Jd}^c N_K^c - N_K^c d_{Kd}^c), \quad (16)$$

$$QLd^c N^c: c_{IJK} \epsilon_{\alpha\beta} L^\alpha Q_{I\alpha}^\beta (d_{Ja}^c N_K^c - N_K^c d_{Ka}^c), \quad (17)$$

$$QQQL: c_{IJK} \epsilon_{\alpha\beta} L^\alpha Q_{I\alpha}^\beta \epsilon_{\gamma\delta} Q_J^\gamma Q_K^\delta \epsilon^{abc}, \quad (18)$$

where $\alpha, \beta, \gamma, \delta$ are $SU(2)$ indices and a, b, c, d are color indices. Three observations are in order: (i) the first two operators are antisymmetric in J, K , (ii) the third operator is symmetric in J, K , and (iii) the sum of all three cyclic permutations of I, J, K vanishes. Hence, if the c_{IJK} were to be totally symmetric in its indices, then all three operators would vanish identically.

We now study the coefficient c_{IJK} in the context of nonrenormalizable interactions in this class of models [15]. As discussed above, in general, some product of fields that acquire intermediate scale *vevs* will still need to be tacked on to $F_I F_J F_K \bar{f}$ to yield a nonvanishing term, since these models contain further $U(1)$ and other symmetries beyond $SU(5) \times U(1)$, and in the particular case of the string-derived flipped $SU(5)$, these terms first appear with one singlet field attached to them (see (13)).

Following the technological recipes of [15], the correlator that could distinguish between the c_{IJK} is among the $SU(5) \times U(1)$ vertex operators,

$$\left\langle e^{W_I \cdot J} (1) e^{W_{F_I} \cdot J} (2) e^{W_{F_I} \cdot J} (3) e^{W_{F_K} \cdot J} (4) \right\rangle = (1 - \bar{z})^{W_{F_I} \cdot W_{F_I}(\bar{z})} W_{F_I} \cdot W_{F_K} C_{IJK}, \quad (19)$$

where the W 's are the weight vectors for the participating $SU(5) \times U(1)$ representations, and C_{IJK} ensures that a gauge singlet is picked out in the product, *i.e.*, $W_{\bar{f}} + W_{F_I} + W_{F_J} + W_{F_K} = 0$. The specific values of the W 's depend on how the singlet is formed. From (15) we see that the F_I field is contracted with the $\bar{f}$ field. This constraint gives specific values to the dot products in (19). In fact, these dot products will remain unchanged as long as in $W_{F_I} \cdot W_{F_J}$, one of the fields is the one contracted with $\bar{f}$, and in $W_{F_J} \cdot W_{F_K}$ none of the fields is contracted with $\bar{f}$.

A different ordering of the I, J, K (*e.g.*, $I \leftrightarrow J$) can always be brought to the same form (19). That is, the field at the point 2 can be taken to be the one contracted with $\bar{f}$. In the low-energy limit, where Planck-mass particle exchanges are all absorbed into an effective point interaction, this reshuffling does not affect the end result, *i.e.*, the value of c_{IJK} . Hence, the coefficient c_{IJK} is totally symmetric in its indices and the effective $d = 4, 5$ operators vanish identically.

We remind the reader that the above result has been established in the free fermionic formulation of the heterotic string, which corresponds to an exact conformal field theory calculation [15]. This fact and the powerful nonrenormalization theorems of supersymmetric theories [19] imply that our result is unchanged by perturbative or nonperturbative σ -model type corrections, and holds to all orders in string loops.

Next we consider whether higher $SU(5) \times U(1)$ invariants could possibly contain the above or other potentially dangerous operators. To be of interest, these would have to have fields acquiring $SU(5) \times U(1)$ *vevs* in order to keep the effective dimension of the light-field operator at 4 or 5. Let us write the

general $SU(5) \times U(1)$ invariant involving the fields in (11) as $F^n \bar{F}^m \bar{f}^p l^q$, with $N = n + m + p + q$, and

$$SU(5) : \quad 2n - 2m - p = 5k, \quad k \in \mathbf{Z}, \quad (20)$$

$$U(1) : \quad n - m - 3p + 5q = 0. \quad (21)$$

Solving these constraints, we find the general invariant

$$F^n \bar{F}^{m-q-3k} \bar{f}^{2q+k} l^q, \quad (22)$$

with $N = 2(n + q - k)$ an even integer and

$$0 \leq n < \frac{N}{2} : \quad -\frac{2}{3} \left(\frac{N}{2} - n \right) \leq k \leq \frac{1}{2} \left(n - \frac{N}{4} \right), \quad (23)$$

$$\frac{N}{2} \leq n \leq \frac{3N}{4} : \quad n - \frac{N}{2} \leq k \leq \frac{1}{2} \left(n - \frac{N}{4} \right). \quad (24)$$

The fact that N is even is in contrast with regular $SU(5)$ where no such constraint exists. Moreover, an invariant with $N > 4$ will be suppressed by $(\langle N^c, \bar{N}^c \rangle / M)^{N-4} \approx 10^{-2(N-4)}$ relative to the basic invariant ($N = 4$). This is in addition to any other singlet-*vev* suppression $\sim \langle \phi \rangle / M_{Pl}$ that might be present in order to get a nonvanishing nonrenormalizable term.

One can now produce the higher- N invariants that follow from (22)–(24). If we denote the $N = 4$ invariants by

$$A : FFF\bar{f}, \quad B : (F\bar{F})^2, \quad C : F\bar{f}\bar{f}l^c, \quad D : \bar{f}l^c\bar{F}\bar{F}, \quad (25)$$

we then get for the $N = 6, 8$, and 10 invariants*

$$N = 6 : \quad \{A, B, C, D\}(F\bar{F}), \quad E : \bar{f}\bar{f}\bar{f}l^c\bar{F}, \quad G : l^c\bar{F}^4, \quad (26)$$

* For $N \geq 11$ the effective $d = 4$ operators already give $\tau_B > 10^{32} \text{ y}$.

$$N = 8 : \quad \{A, B, C, D\}^2, G(F\bar{F}), H : \bar{f}^5 t^3, \quad (27)$$

$$N = 10 : \quad \{A, B, C, D\}^2(F\bar{F}), BG, DG, DE, H(F\bar{F}). \quad (28)$$

All these higher- N invariants are either irrelevant, or products of $N = 4$ invariants or $F\bar{F}$ products, or have $d \geq 6$. We therefore conclude that they make a smaller contribution to baryon decay amplitudes than do the conventional $d = 6$ operators generated by gauge boson exchange.

Another source of effective $d = 4, 5$ operators is through nonrenormalizable superpotential terms containing the $SU(5) \times U(1)$ invariants discussed above, but this time supplemented by condensates of the hidden sector $SO(10) \times SO(6)$ fields that appear in the flipped $SU(5)$ model [12, 18]. However, these condensates must be at least bilinear and amount to a suppression greater than $\langle \hat{H}\hat{H} \rangle / M^2 \sim 10^{-6}$ [18] relative to the basic invariants. In any rate, the presence of the condensates does not affect the calculation of the c_{IJK} coefficients and hence the above analysis applies as well.

Yet another source of baryon decay is through $d = 6$ scalar Higgs exchange operators. The analysis of these operators is greatly model-dependent and would take us away from the scope of this paper. However, in the flipped $SU(5)$ model, including up to quintic superpotential terms [18], we have found that these operators are always much suppressed relative to $d = 6$ gauge boson exchange.

Since we have found no operator contributing to baryon decay in the string-derived flipped $SU(5)$ model that is as large as conventional $d = 6$ gauge boson exchange, we now estimate more carefully the baryon lifetime it implies. We found in Ref. [14] that

$$\Gamma(p \rightarrow \bar{\nu}\pi^+) = 2\Gamma(p \rightarrow e^+\pi^0) = \Gamma(n \rightarrow e^+\pi^-) = 2\Gamma(n \rightarrow \bar{\nu}\pi^0) \gtrsim 20\Gamma(n \rightarrow \bar{\nu}\eta), \quad (29)$$

and

$$\tau(p \rightarrow e^+\pi^0) \approx 3 \times 10^{31} \left(\frac{M_{GUT}}{10^{15} \text{ GeV}} \right)^4 \text{ y}, \quad (30)$$

where M_{GUT} is the scale at which the $SU(3)$ and $SU(2)$ couplings become equal. For comparison, in Ref. [16] we found that the following range of M_{GUT} was compatible with the ranges $\sin^2 \theta_W^{\overline{MS}}(M_Z) = 0.233 \pm 0.001$, $\alpha_3^{\overline{MS}} = 0.115 \pm 0.016$ allowed by LEP:

$$1 \times 10^{15} \text{ GeV} \leq M_{GUT} \leq 3 \times 10^{16} \text{ GeV}. \quad (31)$$

We therefore conclude that

$$3 \times 10^{31} \text{ y} \lesssim \tau(p \rightarrow e^+\pi^0) \lesssim 3 \times 10^{37} \text{ y}, \quad (32)$$

in the flipped $SU(5)$ model. For comparison, the present experimental lower limit is $\tau(p \rightarrow e^+\pi^0) > 3 \times 10^{32} \text{ y}$ [1]. It would clearly be very interesting to push experimental sensitivity further into the range (32), and even be prepared to check the characteristic predictions (29).

We have shown in this paper how flipped $SU(5)$ surmounts another *prima facie* fatal challenge due to baryon decay. We cannot exclude the possibility that the model will be struck down tomorrow by some other disaster. However, the string sends us miracles in order that we believe. We leave it to the reader to decide how many more miracles she wants to see before abandoning her doubts about flipped $SU(5)$.

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