STRANGENESS IN NUCLEON DECAY

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ABSTRACT

We develop different GUT predictions for baryon decays into strange particles. It seems unlikely that minimal conventional GUTs could yield $p \rightarrow \mu^+ K^0$ at a faster rate than $p \rightarrow e^+ \pi^0$. We show that minimal SUSY GUTs, alternative $d = 5$ operators and $d = 6$ Higgs exchanges give characteristically different ratios for $p \rightarrow \mu^+ K^0$, $p \rightarrow \bar{\nu} K^+$ and $n \rightarrow \bar{\nu} K^0$. We give bounds on light sparticle masses in minimal SUSY GUTs and on the masses of B-violating Higgs bosons. These are still marginally consistent with constraints from particle phenomenology and from cosmology, and suggest that baryon decays into strange particles may be detectable in the near future.

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It is by now notorious that baryons seem to suffer from stage fright, and so far decline to decay¹⁾. We confess to being disappointed, but are not yet discouraged. There is still a remaining window of opportunity to catch a baryon decay into strange particles, such as $p \rightarrow \mu^+ K^0$ or $\bar{\nu} K^+$ or $n \rightarrow \bar{\nu} K^0$, before the atmospheric neutrino background sabotages further experimental searches. Indeed, the frenzy of recent theoretical work on SUSY GUTs²⁾ seems to favour strangeness in baryon decays³⁻⁷⁾, and recent experimental candidates⁸⁾ have fostered the suspicion that such possibilities may not only be strictly for the birds. In this paper we analyze the information that can be gleaned from experimental searches for baryon decays into strange particles, within different GUT contexts.

We first argue that flavour SU(3) symmetry breaking is unlikely to enhance $p \rightarrow \mu^+ K^0$ much above $p \rightarrow e^+ \pi^0$ in conventional minimal GUTs⁹⁾. We then pass to the $d = 5$ operators¹⁰⁾ mediated by Higgsino exchange in minimal SUSY GUTs¹¹⁾ and estimate the relative rates for $p \rightarrow \bar{\nu} K^+$, $n \rightarrow \bar{\nu} K^0$ and $p \rightarrow \mu^+ K^0$ (see also Refs. 12 and 13). We then use the experimental lower bounds⁶⁾ on $\tau(B \rightarrow \bar{\nu} K)$ and $\tau(p \rightarrow \mu^+ K^0)$ to constrain the masses of light gauginos in such models. We find small allowed regions which are consistent with other particle physics constraints and with cosmology¹⁴⁾, and in which W decays into pairs of supersymmetric particles are possible¹⁵⁾. We also obtain the ratios of $p \rightarrow \bar{\nu} K^+$, $n \rightarrow \bar{\nu} K^0$ and $p \rightarrow \mu^+ K^0$ or $e^+ K^0$ favoured by alternative $d = 5$ operators⁷⁾ allowed by SUSY but not generated in minimal SUSY GUTs. Finally, we estimate the ratios of $p \rightarrow \bar{\nu} K^+$, $n \rightarrow \bar{\nu} K^0$ and $p \rightarrow \mu^+ K^0$ to be expected from $d = 6$ Higgs exchange⁵⁾ if the $d = 5$ operator vanishes, and give a lower limit on the mass of the B-violating triplet Higgs boson which is consistent with values expected in SUSY cosmology models. Measurements of the relative rates for $p \rightarrow \bar{\nu} K^+$, $n \rightarrow \bar{\nu} K^0$ and $p \rightarrow \mu^+ K^0$ or $e^+ K^0$ would be able to distinguish between these different theoretical models.

SHADOW OF A DOUBT?

In minimal conventional (i.e. not supersymmetric) SU(5) à la Georgi-Glashow⁹⁾ the conventional expectation¹⁷⁾ is that $\tau_p = 10^{29 \pm 2}$ years, with $p \rightarrow e^+ \pi^0$ an important decay mode of branching ratio 0(30)%. These predictions seem to conflict

with the present limit $\tau(p \rightarrow e^+ \pi^0) > 2 \times 10^{32}$ years¹⁾. Is there any way in which Georgi and Glashow⁹⁾ can be resuscitated, assuming one finds this desirable? One possibility¹⁸⁾ is to introduce additional low-mass particles from split GUT multiplets, which have the effect of pushing up m_X and hence $\tau_p \propto m_X^4$. This appears rather artificial and unmotivated. Another possibility might be that up to now τ_p for a given value of m_X has been underestimated. Recent calculations^{13,19)} of $\tau(p \rightarrow e^+ \pi^0)$ using current algebra, PCAC and phenomenological estimates of the 3-quark overlap inside a baryon disfavour this possibility. Finally, it might be that the $p \rightarrow e^+ \pi^0$ branching ratio has been grossly over-estimated, and perhaps $p \rightarrow \mu^+ K^0$ dominates. In conventional SU(5), current algebra and PCAC²⁰⁾ give

$$\frac{\Gamma(p \rightarrow \mu^+ K^0)}{\Gamma(p \rightarrow e^+ \pi^0)} \approx 2 \left(\frac{1+D-F}{1+D+F} \right)^2 \frac{\left(1 - \frac{m_K^2}{m_p^2}\right)}{\left(1 - \frac{m_\pi^2}{m_p^2}\right)} \approx \frac{1}{2} \quad (1)$$

if we neglect SU(3) violating effects in the 3-quark annihilation amplitudes. It could well be that the PCAC extrapolation is in error, especially for kaons, but model calculations²¹⁾ suggest that it should be reasonably reliable, at least for pions. The first terms in both the denominator and the numerator of Eq. (1) are related by soft pseudoscalar theorems to the same $\langle 0 | uud | p \rangle$ matrix element. However, the D and F terms in Eq. (1) are related in the denominator to the same $\langle 0 | uud | p \rangle$ matrix element, but in the numerator to a $\langle 0 | uus | \Sigma^+ \rangle$ matrix element which could *a priori* differ by an SU(3) breaking factor λ :

$$\frac{\Gamma(p \rightarrow \mu^+ K^0)}{\Gamma(p \rightarrow e^+ \pi^0)} \xrightarrow{?} 2 \left(\frac{1 + \lambda(D-F)}{1+D+F} \right)^2 \frac{\left(1 - \frac{m_K^2}{m_p^2}\right)}{\left(1 - \frac{m_\pi^2}{m_p^2}\right)} \quad (2)$$

Numerically, it is different to get a large enhancement factor from λ in Eq. (2) because $(D-F) \approx 0.28 \ll 1$. Furthermore, it is not possible to introduce the desirable SU(3) breaking $\lambda > 1$ into the 3-quark operator matrix elements $\langle 0 | qq\bar{q} | B \rangle$ using the formalism of chiral Lagrangians²²⁾. For minimal conventional SU(5) operators^{*}) $(s_R u_R)_{u_L}$, $(s_R u_R)_{d_L}$ and $(s_L u_L)_{u_R}$, these are controlled in the SU(3) limit

^{*}) We use the notation $(f_L f'_L) \equiv \epsilon^{\alpha\beta} f_\alpha f'_\beta$, $(f_R f'_R) \equiv \epsilon_{\dot{\alpha}\dot{\beta}} f^{\dot{\alpha}} f'^{\dot{\beta}}$.

by $\text{Tr } (\tilde{O} B_{L,R})$ and by $\text{Tr } (\tilde{O}' B_R)$ where (we use the notation of Ref. 22):

$$\tilde{O} \equiv \begin{pmatrix} 0 & 0 & 0 \\ 1 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad \tilde{O}' \equiv \begin{pmatrix} 0 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad B \equiv \begin{pmatrix} \sqrt{\frac{1}{2}} \Sigma^0 + \sqrt{\frac{1}{6}} \Lambda^0 & \Sigma^+ & p \\ \Sigma^- & -\sqrt{\frac{1}{2}} \Sigma^0 + \sqrt{\frac{1}{6}} \Lambda^0 & n \\ \Xi^- & \Xi^0 & -\sqrt{\frac{2}{3}} \Lambda^0 \end{pmatrix} \quad (3)$$

These clearly give $\Sigma^{+,0}$ and Λ^0 to vacuum matrix elements. Now, SU(3) symmetry breaking can be introduced into these traces by incorporating factors of

$$m \equiv \begin{pmatrix} m_u & 0 & 0 \\ 0 & m_d & 0 \\ 0 & 0 & m_s \end{pmatrix} \quad (4)$$

e.g. $\text{Tr } (\tilde{O} m^n B m^{n'})$, $\text{Tr } (\tilde{O}' m^n B m^{n'})$. However, within this formalism, the $\Sigma^{+,0}$ and Λ^0 to vacuum matrix elements of interest are not sensitive to the ratio of $m_s/m_{u,d}$. Hence, one cannot make $\lambda > 1$, except by the seemingly artificial expedient of enforcing a partial cancellation between SU(3) invariant and m_s dependent terms in $\langle 0 | (d_R u_R) u_L \rangle$ or $\langle d_L u_L | u_R | p \rangle$, which are proportional to

$$\text{Tr } (O m^n B m^{n'}) : O \equiv \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix} \quad (5)$$

in the chiral Lagrangian formalism. By way of contrast, SU(3) breaking which enhances $n \rightarrow \bar{\nu} K^0$, $p \rightarrow \bar{\nu} K^+$ can be incorporated naturally into matrix elements of $\langle 0 | (d_R u_R) s_L | \Lambda^0 \rangle$ and $\langle d_L u_L | s_L | \Lambda^0 \rangle$ which are proportional to

$$\text{Tr } (\tilde{O}'' m^n B m^{n'}) : \tilde{O}'' \equiv \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix} \quad (6)$$

in the chiral Lagrangian formalism²²⁾. However, these operators do not appear in minimal conventional SU(5). We conclude that the chiral Lagrangian formalism²²⁾, equivalent to current algebra and PCAC²⁰⁾, does not offer the prospect of SU(3) breaking which enhances $p \rightarrow \mu^+ K^0$ relative to $p \rightarrow e^+ \pi^0$ above the naïve prediction (1).

THE TROUBLE WITH HIGGSINOS

In minimal supersymmetric SU(5) the dominant mechanism of B decay involves a supersymmetric $d = 5$ operator¹¹⁾ generated by the exchanges of superheavy colour triplet Higgsino fermions and Higgs bosons H_3 . This operator must be dressed by supersymmetric fermion exchange as in Fig. 1 in order to give $d = 6$ four-fermion operators. The exchanges of neutral supersymmetric fermions $\tilde{g}$, $\tilde{W}^3$, $\tilde{B}^0$ and $\tilde{H}_{1,2}^0$ do not contribute^{6,7)} in leading, Cabibbo-unsuppressed order. Charged supersymmetric fermions $\tilde{W}^\pm$ and $\tilde{H}_{1,2}^\pm$ do contribute, and in general these charged fermions mix through the mass matrix¹⁵⁾

$$\begin{pmatrix} \tilde{W}^+ & \tilde{H}_1^+ \end{pmatrix}_L \begin{pmatrix} M_2 & g_2 v_2 \\ g_2 v_1 & -\epsilon \end{pmatrix} \begin{pmatrix} \tilde{W}^- \\ \tilde{H}_2^- \end{pmatrix}_L \quad (7)$$

where M_2 is the SUSY violating SU(2) gaugino mass, $v_{1,2} \equiv \langle 0 | H_{1,2}^0 | 0 \rangle$: $m_W^2 = \frac{1}{2}g_2^2 (v_1^2 + v_2^2)$, and ϵ is an $(\tilde{H}_1 \tilde{H}_2)$ mixing term. The $\tilde{H}^\pm$ components of the mass eigenstates $m_{1,2}$ of the matrix (7) do not contribute significantly to B decay, since their Yukawa couplings are suppressed by factors $O(m_f/m_W)$. The mass eigenstates are related to the interaction eigenstates $(\tilde{W}^+, \tilde{H}_1^+)$ and $(\tilde{W}^-, \tilde{H}_2^-)_2$ by orthogonal rotations^{15,23)}

$$R_\pm \equiv \begin{pmatrix} \cos \Theta_\pm & \sin \Theta_\pm \\ -\sin \Theta_\pm & \cos \Theta_\pm \end{pmatrix} : \tan \Theta_\pm = \frac{-b_\pm \pm \sqrt{b_\pm^2 + 4a_\pm^2}}{2a_\pm} \quad (8)$$

where

$$a_\pm = \begin{cases} M_2 g_2 v_1 & -\epsilon g_2 v_2 \\ M_2 g_2 v_2 & -\epsilon g_2 v_1 \end{cases} ; b_\pm = M_2^2 - \epsilon^2 \pm g_2^2 (v_2^2 - v_1^2) \quad (9)$$

Therefore, incorporating factors coming from the heavy triplet Higgsino and Higgs couplings, short distance renormalization, and light charged supersymmetric fermion loops, we find^{4,7,23)} the following effective Lagrangian for B decay in minimal SUSY SU(5):

$$\mathcal{L} = -\frac{g_2^2}{16\pi^2} A \frac{1}{m_{H_3}} \left(\frac{v_1^2 + v_2^2}{v_1 v_2} \right) K O_W \quad (10)$$

where K is a kinematic function of $m_{\tilde{q}}$ (assumed for simplicity to be the same as $m_{\tilde{\ell}}$), m_W and $m_{1,2}$:

$$K = 2m_W \left[\cos\Theta_+ \cos\Theta_- \left(\frac{m_1}{m_{\tilde{q}}^2 - m_1^2} \left[1 - \frac{m_1^2}{m_{\tilde{q}}^2 - m_1^2} \ln \frac{m_{\tilde{q}}^2}{m_1^2} \right] \right) + \sin\Theta_+ \sin\Theta_- \left(\frac{m_2}{m_{\tilde{q}}^2 - m_2^2} \left[1 - \frac{m_2^2}{m_{\tilde{q}}^2 - m_2^2} \ln \frac{m_{\tilde{q}}^2}{m_2^2} \right] \right) \right] \quad (11)$$

and

$$\begin{aligned} \frac{O_W}{2} = & m_c m_s \cos\Theta_c \sin^2\Theta_c \left[(us)_L (d\nu_\mu)_L + (ud)_L (s\nu_\mu)_L \right] \\ & + m_c m_s \sin^3\Theta_c \left[(ud)_L (d\nu_\mu)_L \right] \\ & + m_u m_s \cos\Theta_c \left[(su)_L (u\mu - d\nu_\mu \cos\Theta_c)_L \cos\Theta_c + 2\cos^2\Theta_c (du)_L (s\nu_\mu)_L \right] \end{aligned} \quad (12)$$

For numerical purposes we will assume $m_c = 1.5$ GeV, $m_s = 0.15$ GeV, $m_d/m_s = 1/20$ and $m_u/m_d = 1/1.8$. In writing (12), we have neglected contributions coming from Higgs couplings to the third generation, since recent bounds²⁴⁾ on the second and third Kobayashi-Maskawa angles mean that heavy quark contributions can be no larger than the light quark contributions (12).

There is⁷⁾ an alternative form of SUSY $d = 5$ operator which cannot be simply obtained from superheavy triplet Higgs(ino) exchange, but could be generated by dynamics at the Planck scale, for example. It gives a four-fermion operator proportional to

$$O'_W = (su)_L (ul - d\nu_\ell)_L + 2(du)_L (s\nu_\ell)_L \quad (13)$$

The chiral Lagrangian formalism²²⁾ can be used^{25,13)} to obtain from the forms (12) and (13) predictions for the relative decay rates for $p \rightarrow \bar{\nu}K^+$, $n \rightarrow \bar{\nu}K^0$ and $p \rightarrow \mu^+K^0$ [or e^+K^0 in the case of (13)]. Algebraic expressions and numerical results are shown in Table 1. We see that the operators (12) and (13) give qualitatively different results, with μ^+K^0 not expected^{3,4)} to have a large relative rate if superheavy Higgsino/Higgs exchange dominates. The relative rates for $p \rightarrow \bar{\nu}K^+$ and $n \rightarrow \bar{\nu}K^0$ are also different in the two cases, and our ratios based on current

algebra and PCAC agree qualitatively with those previously obtained using non-relativistic SU(6) technology¹²⁾.

There has been considerable discussion in the literature^{3,4,6,10,12,13)} questioning whether the predicted lifetime for $B \rightarrow \bar{\nu} K$ decay in minimal SUSY GUTs is compatible with the ever-increasing experimental lower limit⁸⁾. For orientation, we estimate⁴⁾ $m_{H_3} \approx 10^{16}$ GeV and take an estimate of the relevant $\langle 0 | qqq | B \rangle$ matrix element from Ref. 13:

$$\langle 0 | (du)_L u_L | p \rangle = \beta \quad : \quad \beta = 0.03 \text{ GeV}^3 \quad (14)$$

The value of m_{H_3} is rather model-dependent -- it may or may not be larger or smaller than m_X in SUSY GUTs, and the best estimate of m_X itself depends sensitively on the light particle spectrum. Also, estimates of β in the literature vary widely, with (14) being a relatively large value. Baryon decay rates depend on the ratio $(\beta/m_{H_3})^2$ which we parametrize using R :

$$R^2 \left(\frac{\beta}{m_{H_3}} \right)^2 = \left(\frac{0.03 \text{ GeV}^3}{10^{16} \text{ GeV}} \right)^2 \quad (15)$$

We then find

$$\tau(B \rightarrow \bar{\nu} K) = 3.6 \times 10^{26} \text{ yr} \times \left(\frac{2v_1 v_2}{v_1^2 + v_2^2} \right)^2 \times \frac{R^2}{K^2} \quad (16)$$

where models suggest

$$1 \leq v_1/v_2 \leq 4 \quad (17)$$

as a plausible range^{14,15)}. Taking for example the case $v_1 = v_2$, the experimental limit $\tau(n \rightarrow \bar{\nu} K^0) > 8 \times 10^{31}$ years requires in Eq. (16) that

$$|K| < 2 \times 10^{-3} |R| \quad (18)$$

In view of the uncertainties in β and m_{H_3} , we take $1 \leq R \leq 10$ as a typical range. Equation (18) therefore tells us that experiment requires $|K| \ll 1$. The value of $K(11)$ depends sensitively on the assumed spectrum of SUSY particles via $m_{\tilde{q}}$, M_2 and ε . Small values of K can be obtained if $m_{\tilde{q}} \gg m_W$ and/or $|M_2|$, $|\varepsilon| \ll m_W$. In the plausible case $m_{\tilde{q}} \approx m_W$ we find:

$$|K| = \left| \frac{M_2 + 2\epsilon}{3m_W} \right| + O\left(\frac{M_2, \epsilon}{m_W}\right)^3 \quad (19)$$

which shows explicitly that, as expected, a Majorana mass term (either for $\tilde{W}$ or for $\tilde{H}$) is required in order for Fig. 1 to be non-zero. Formula (19) has the interesting feature that a partial or complete cancellation is possible if M_2 and ϵ have opposite signs. Indeed, there is a locus of zeroes of K in the (M_2, ϵ) plane, in the neighbourhood of which the bound (18) is satisfied. Figures 2 to 4 exhibit the (M_2, ϵ) plane for different signs of ϵ , ratios v_1/v_2 and values of $m_{\tilde{q}}/m_W$. The domains in which $|K|$ obeys the constraint (18) for $|R| = 1, 10$ and 100 are indicated. Also shown in Figs. 2 and 3 are the domains of the (M_2, ϵ) plane which are cosmologically acceptable, in the sense that the relic density of the lightest neutral sparticle is sufficiently small¹⁴⁾. We do not have available analogous calculations for the case $m_{\tilde{q}} = 2m_W$, but we expect the cosmologically acceptable domains of (M_2, ϵ) to be further reduced by analogy to the changes between Figs. 2 and 3. We see from Figs. 2 to 4 that while the Higgsino exchange diagrams generally cause baryons to decay too rapidly, there are regions which are not yet excluded by experiment. Interestingly, these regions are in the domain where $W^\pm$ decays into pairs of supersymmetric fermions are kinematically accessible¹⁵⁾. These domains would, of course, be substantially extended if m_{H_3} were significantly larger than 10^{16} GeV and we feel it is premature to exclude the existence of $d = 5$ operators, given the present uncertainties. The trouble with Higgsinos may not be fatal. Even if acceptable ranges of the parameters M_2 and ϵ are eventually excluded by experiment, it is possible to suppress the $d = 5$ operators entirely and still obtain an acceptable cosmological baryon asymmetry and observable baryon decay, as we now discuss.

THE $d = 5$ OPERATOR VANISHES?

There are variants both of conventional and of supersymmetric SU(5) in which the "superheavy" Higgs triplets H_3 are significantly lighter than 10^{16} GeV, perhaps of order 10^{10} GeV for reasons associated with cosmology and the gauge hierarchy problem^{5,16)}. It is a generic feature of supersymmetric and especially

supergravity GUTs that the grand unified phase transition takes place when the temperature $T \sim \sqrt{m_W m_P} \sim 10^{10}$ GeV, and such a delay offers a mechanism for suppressing the number of monopoles created. This temperature $T \sim 10^{10}$ GeV is comparable with the scale at which non-perturbative SU(5) GUT effects may become important and disfavour the SU(5) invariant minimum of the Higgs potential. It is evident that if a cosmological baryon asymmetry is to be produced by Higgs decays after such a phase transition, the relevant Higgs bosons cannot weigh more than about 10^{10} GeV. This is indeed the natural mass-scale for such bosons in models exploiting the missing partner mechanism and a pair of 24 Higgs representations²⁶⁾ instead of a 75 representation: $m_{H_3} = 0(m_X^4/m_P^3)$. In this case, surely the $d = 5$ operator must vanish, and there exist variants of SUSY SU(5) in which the errant $d = 5$ operator is banished by a discrete symmetry¹⁰⁾. In these cases, the dominant source of baryon decay may be the $d = 6$ Higgs boson exchange operator⁵⁾. In theories where only two Higgs multiplets contribute significantly to the quark and lepton mass matrices, the dominant $d = 6$ operators are:

$$\begin{aligned} \mathcal{L}_{\text{eff}} = & \frac{1}{v_1^2 m_{H_3}^2} \left\{ \sin^2 \theta_c m_S^2 (su)_R (u\mu)_L - m_d m_S (du)_R (s\nu_\mu)_L + \dots \right\} \\ & + \frac{1}{v_2^2 m_{H_3}^2} \left\{ m_u^2 (ud)_L (ue)_R + \dots \right\} \end{aligned} \quad (20)$$

The first two terms in Eq. (20) give essentially equal amplitudes, which result in the relative rates for $p \rightarrow \bar{\nu} K^+$, $n \rightarrow \bar{\nu} K^0$ and $p \rightarrow \mu^+ K^0$ shown in Table 1. We note that the relative rates to be expected are characteristically different from those due to Higgsino or other SUSY $d = 5$ operators. In addition, the $p \rightarrow e^+ \pi^0$ decay rate is given by

$$\frac{\Gamma(p \rightarrow e^+ \pi^0)}{\Gamma(n \rightarrow \bar{\nu}_\mu K^0)} = \frac{1}{2} \left(\frac{v_1}{v_2} \right)^4 \left(\frac{m_u^2}{m_d m_S} \right) \left(\frac{1+D+F}{1+\frac{D}{3}+F} \right)^2 \approx \frac{1}{60} \left(\frac{v_1}{v_2} \right)^4 \quad (21)$$

Therefore the best bounds on m_{H_3} probably come (almost equivalently) from the experimental lower limits on $n \rightarrow \bar{\nu} K^0$ and on $p \rightarrow \mu^+ K^0$:

$$m_{H_3} \left(\frac{g_2 v_1}{m_W} \right) \geq 2.3 \times 10^{11} \text{ GeV} \quad (22)$$

This is within the range of values of m_{H_3} which are suggested by cosmological considerations.

We conclude from Figs. 2 to 4 and from Eq. (22) that while the present lower limits on baryon decays into strange decay modes push very hard on the $d = 5$ Higgsino exchange model and on the $d = 6$ Higgs exchange model, these are not yet excluded experimentally. Looking on the optimistic side, both models suggest that baryon decays into strange particles should be detectable in the near future. Furthermore, chiral Lagrangian calculations which embody current algebra and PCAC predict characteristically different ratios of rates for $p \rightarrow \bar{\nu}K^+$, $n \rightarrow \bar{\nu}K^0$ and $p \rightarrow \mu^+K^0$ in the different models. Thus, if baryons do decay into the available strange window before the curtain of the neutrino background falls, we may be able to disentangle the mechanisms responsible for their decay.

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Table 1

Relative ratios for baryon decays into strange particles

	$\tilde{d} = 5$ $\tilde{H}, H$ exchange	Alternative $d = 5$ SUSY operator	$d = 6$ H exchange
$p \rightarrow \nu K^+$	$\left(\frac{1+D+F}{2(1+F)}\right)^2 = 0.55$	$\left(\frac{2(1+F)}{1+D+F}\right)^2 = 1.75$	1
$n \rightarrow \bar{\nu} K^0$	1	1	1
$p \rightarrow \mu^+ K^0$	$\left(\frac{1-D+F}{2(1+F)}\right)^2 \left(\frac{m_u \cos \theta_c}{m_c \sin^2 \theta_c}\right)^2$ $\approx 2 \times 10^{-3} \text{ to } 2 \times 10^{-4}$	$\left(\frac{1-D+F}{1+D+F}\right)^2 \approx 0.10$	$\left(\frac{\sin^2 \theta_c m_s}{m_d}\right)^2 \left(\frac{1+D-F}{1+\frac{D}{3}+F}\right)^2 \approx 0.55$

We use $D = 0.76$ and $F = 0.48$ (Ref. 27).

Figure captions

- Fig. 1 : a) $d = 5$ operator generated by Higgsino/Higgs exchange, dressed b) by gaugino fermion exchanges to give a four-fermion operator.
- Fig. 2 : Domains of (M_2, ϵ) space consistent with cosmology (---), the absence of a charged sparticle lighter than 20 GeV (-.-), and the absence of baryon decay (18) with $R = 1$ (—), $R = 10$ (----) and $R = 100$ (.....): (a) $v_1 = v_2$, $\epsilon > 0$; (b) $v_1 = v_2$, $\epsilon < 0$; (c) $v_1 = 4v_2$, $\epsilon > 0$; (d) $v_1 = 4v_2$, $\epsilon < 0$; all taking $m_{\tilde{q}} = m_{\tilde{\ell}} = \frac{1}{2}m_W$.
- Fig. 3 : As for Fig. 2, but with $m_{\tilde{q}} = m_{\tilde{\ell}} = m_W$.
- Fig. 4 : As for Fig. 2, but with $m_{\tilde{q}} = m_{\tilde{\ell}} = 2m_W$.

Fig. 1

Fig. 2

Fig. 3

Fig. 4