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OBSERVABLE GRAVITATIONALLY INDUCED BARYON DECAY*

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ABSTRACT

We find that in supersymmetric theories gravitodynamic effects scaled by the inverse of the Planck mass can induce baryon decay at an observable rate. In a minimal supersymmetric (susy) grand unified theory (GUT) the dominant gravitationally induced baryon decay mode is $B \to \bar{\nu} + K$, with a likely admixture of $p \to (e^+$ or $\mu^+) + K$. As a by-product, we present an improved estimate of Higgs-mediated baryon decay branching ratios in minimal susy GUTs. We consider the possibility that a loss of quantum coherence may be observable in gravitationally induced baryon decay, but argue that this would be difficult to reconcile with successful experimental tests of quantum mechanics.

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Many theoretical frameworks for baryon decay have been discussed in the last few years. These include conventional grand unified theories (GUTs) [1], variants of supersymmetric GUTs [2]-[5], non-perturbative effects [6] in the electroweak gauge theory, and baryon number violation catalyzed by grand unified monopoles [7]. Zel'dovich [8] and others [9],[10] have suggested that non-perturbative quantum gravitational effects such as virtual black holes may also lead to baryon decay, albeit with an unobservably long lifetime of order 10^{50} years or more [10]. Three of us [11] have recently pointed out that in a supersymmetric theory such a "gravitodynamic" baryon decay amplitude may be scaled by $O(1/m_P)$, leading to faster baryon decay which could be observable in the present generation of experiments, even if the putative grand unification mass is much larger than the canonical $O(10^{16})\,GeV$ encountered in minimal susy GUTs.

In this paper we explore the phenomenology of gravitodynamic baryon decay, assuming a susy GUT framework. We show that only a few distinct dimension 5 operators can be important, and quantify the baryon decay branching ratios to be expected in each case. We find that $B \to \bar{\nu} + K$ dominates, with a likely admixture of $p \to (e^+ \, or \, \mu^+) + K$. As a by-product, we correct previous estimates [2],[3],[12],[13] of baryon decay branching ratios in conventional susy GUTs, confirming the dominance [3] of $B \to \bar{\nu} + K$ induced by $\tilde{W}$ exchange. We show that the baryon decay rate can be observable even if the dimensional coefficient of the gravitationally induced dimension 5 operators is much smaller than $m_P^{-1} = O(10^{-19})\,GeV^{-1}$. We discuss the possibility [9],[10] that quantum coherence may be lost in baryon decay and other processes induced by non-perturbative quantum gravitational phenomena such as virtual black holes.

We assume that particle physics is described by a renormalizable spontaneously broken local susy gauge theory at energies much below the Planck mass, augmented by non-renormalizable interactions scaled by the appropriate inverse power of the Planck mass [11]. We allow these non-renormalizable interactions to violate all global symmetries consistent with gauge invariance and susy. The requirement of susy is essential in order to protect the squarks $\tilde{q}$ and sleptons $\tilde{\ell}$ from acquiring masses $O(m_P)$, and also to prevent the appearance of quartic scalar interactions such as $\tilde{q}\tilde{q}\tilde{q}\tilde{\ell}$ with coefficients of order unity, which would lead to catastrophically rapid baryon decay. A general effective phenomenological action framework for these interactions is provided by the work of Cremmer *et al.* [14]. We expect such interactions by analogy with low energy effective interactions in strong interaction physics, and are agnostic about their detailed origins. Maybe there are many particles with masses $O(m_P)$ which can mediate new interactions. Or maybe our known "elementary" particles are in fact composite at the Planck scale [15], with novel interactions corresponding to constituent interchanges. Another possible source is non-perturbative quantum gravitational effects such as virtual black holes [8] or space-time foam [9],[10].* We would certainly expect an effective Lagrangian approach to be applicable in the first two cases, but it has been suggested [9] that in the latter case quantum coherence might be lost, entailing a revision [10] of conventional quantum field theoretical rules. Later on, however, we will encounter reasons for discounting the observability of incoherence, and for playing according to the field-theoretical rules, at least as a first approximation.

We will be concerned with possible non-renormalizable terms in the chiral superpotential, and the lowest-dimensional terms of interest are quartic. We assume an SU(5) GUT framework so that quarks and leptons appear in chiral $\bar{5}$ (F^α) and $\underline{10}$ ($T_{\beta\gamma}$) superfields. The only interesting quartic combinations which can violate B and L conservation are

$$F^\alpha_a T^b_{\alpha\beta} T^c_{\gamma\delta} T^d_{\lambda\mu}\, \epsilon^{\beta\gamma\delta\lambda\mu} \tag{1}$$

where the Latin indices denote different generations of quarks and leptons. These give the same dimension 5 operators [16] as Higgs exchange in minimal susy GUTs [17], but in a potentially different algebraic combination. One can extract from (1) $\Delta B = \Delta L = \pm 1$ interactions of the form $\ell^a_L q^b_L q^c_L q^d_L$ and $q^{c^a}_L \ell^{c^b}_L q^{c^c}_L q^{c^d}_L$. Because of colour antisymmetrization, the three quark generation indices cannot all be the same, and none of the $q^c_L \ell^c_L q^c_L q^c_L$ operators can contribute to baryon decay when dressed by gaugino exchange. However the following $\ell_L q_L q_L q_L$ operators can contribute when dressed by gaugino exchange:

*We note that extant foam [9],[10] calculations respect discrete reflection symmetries on the matter fields: $\phi \to -\phi$, and we assume such a symmetry in the following. This has the effect of forbidding catastrophic B and L-violating trilinear qqq and $q\bar{q}\ell$ interactions.

$$
\begin{aligned}
\mathcal{L}_{eff} = \int d^2\theta \epsilon^{ijk} \sum_{\ell=e,\mu,\tau} \Big[& g^{\ell}_{cuu}(\ell c_i u_j d'_k - \nu_\ell s'_i u_j d'_k) + g^{\ell}_{ucc}(\ell u_i c_j s'_k - \nu_\ell d'_i c_j s'_k) \\
& + g^{\ell}_{tuu}(\ell t_i u_j d'_k - \nu_\ell b'_i u_j d'_k) + g^{\ell}_{utt}(\ell u_i t_j b'_k - \nu_\ell d'_i t_j b'_k) \\
& + \frac{1}{2} g^{\ell}_{uct}(\ell u_i c_j b'_k - \ell u_i s'_j t_k - \nu_\ell d'_i c_j b'_k + \nu_\ell d'_i s'_j t_k) \\
& + \frac{1}{2} g'^{\ell}_{uct}(\ell c_i t_j d'_k - \ell c_i b'_j u_k - \nu_\ell s'_i t_j d'_k + \nu_\ell s'_i b'_j u_k) \Big].
\end{aligned}
\tag{2}
$$

where the indices i,j,k denote SU(3) colour and the Cabibbo-rotated charge -1/3 quarks are denoted by primes. We might naively expect on dimensional grounds that the coefficients g in (2) would be $O(m_P^{-1}) = O(10^{-19})GeV^{-1}$. The superspace integration $\int d^2\theta$ picks out pairs of left-handed fermion components from each of the products of four superfields in the sum (2). The resulting two fermion-two boson products can then be dressed by SU(2) $\tilde{W}^{\pm,0}$, U(1) $\tilde{B}$ or SU(3) $\tilde{g}$ gaugino exchange to give four-fermion operators [16]. The flavour-conserving $\tilde{B}$ and $\tilde{g}$ exchanges cannot possibly give operators relevant to baryon decay, except from the terms $\epsilon^{ijk}\nu_\ell s'_i u_j d'_k$, $\epsilon^{ijk}\nu_\ell b'_i u_j d'_k$ and $\epsilon^{ijk}\nu_\ell s'_i b'_j u_k$. However, $\tilde{B}$ and $\tilde{g}$ exchanges both generate from these operators symmetric combinations of four-fermion operators such as

$$
\epsilon^{ijk}\Big[(\nu_\ell s'_i)_L(u_j d'_k)_L + (\nu_\ell d'_k)_L(s'_i u_j)_L + (\nu_\ell u_j)_L(s'_i d'_k)_L\Big] \tag{3}
$$

where $(ff')_L$ denotes the Lorentz scalar product of two left-handed fermion fields. However, this combination vanishes because of a simple algebraic identity. Therefore $\tilde{B}$ and $\tilde{g}$ dressings of dimension 5 operators do not contribute to baryon decay. However, dressing the dimension 5 operators with $\tilde{W}^{\pm}$ gauginos can contribute to baryon decay. Exchanges of the complete isotriplet $\tilde{W}^{\pm,0}$ gives rise to the following unrenormalized $SU(3) \times SU(2) \times U(1)$ invariant four-fermion operators (in the notation of ref. [3]):

$$\mathcal{L}^0_{eff} = b^0 \sum_{\ell=e,\mu,\tau} \left[g^{\ell}_{cuu}\left\{\frac{1}{2}\left[O^4_{cuu\ell} - 3O^3_{cuu\ell} + 3O^3_{uuc\ell}\right]\right\}\right.$$

$$+ g^{\ell}_{ucc}\left\{\frac{1}{2}\left[O^4_{ucc\ell} - 3O^3_{ucc\ell} + 3O^3_{ccu\ell}\right]\right\} \tag{4}$$

$$\left. + \; similar \; terms \right] .$$

where b^0 is a loop integral:

$$b^0 = \frac{G_F}{\sqrt{2}} \frac{m^2_W m_{\tilde{W}}}{32\pi^2}\left[f(m_{\tilde{q}}, m_{\tilde{q}}, m_{\tilde{W}}) + f(m_{\tilde{q}}, m_{\tilde{\ell}}, m_{\tilde{W}})\right] \tag{5a}$$

with

$$f(m_1, m_2, m_3) \equiv \frac{1}{m_2^2 - m_3^2}\left(\frac{m_2^2}{m_1^2 - m_2^2}\, \ell n \frac{m_1^2}{m_2^2} - \frac{m_3^2}{m_1^2 - m_3^2}\, \ell n \frac{m_1^2}{m_3^2}\right). \tag{5b}$$

If we guess

$$\left(\frac{m_{\tilde{W}}}{m_W}\right)\left[f(m_{\tilde{q}}, m_{\tilde{q}}, m_{\tilde{W}}) + f(m_{\tilde{q}}, m_{\tilde{\ell}}, m_{\tilde{W}})\right] \simeq \frac{1}{m^2_W} \tag{6}$$

we get

$$b^0 \simeq 2.2 \times 10^{-6}\, GeV^{-1} . \tag{7}$$

This result is renormalized in the usual way by gauge loop corrections. Ignoring renormalization between distances m_P^{-1} and m_X^{-1}, the renormalization at larger distance scales is

$$A = \left[\frac{\alpha_3(1GeV)}{\alpha_3(m_c)}\right]^{2/9}\left[\frac{\alpha_3(m_c)}{\alpha_3(m_b)}\right]^{6/25}\left[\frac{\alpha_3(m_b)}{\alpha_3(m_t)}\right]^{6/23}\left[\frac{\alpha_3(m_t)}{\alpha_3(m_W)}\right]^{2/7}$$

$$\left[\frac{\alpha_3(m_W)}{\alpha_{SUM}}\right]^{4/3}\left[\frac{\alpha_2(m_W)}{\alpha_{SUM}}\right]^{-3}\left[\frac{\alpha_1(m_W)}{\alpha_{SUM}}\right]^{-1/66} \tag{8}$$

which takes the value

$$A \simeq 15 \tag{9}$$

when we make the illustrative choices

$$\alpha_3(m_W) = 0.12, \alpha_2(m_W) = \frac{1}{31}, \alpha_1(m_W) = \frac{1}{50}, \alpha_{SUM} = \frac{1}{24}\,. \tag{10}$$

The coefficient of the renormalized operators is therefore

$$b \equiv Ab^0 \simeq 3.2 \times 10^{-5}\, GeV^{-1}\,. \tag{11}$$

In order to see whether this susy gravitodynamic mechanism can lead to baryon decay at an observable rate, we exploit the non-relativistic SU(6) analysis by Salati and Wallet [12] of the baryon decay modes induced by a general $\mathcal{L}_{eff}$. In their notation we have

$$\begin{aligned} A^L(\bar{\nu}_R; \Delta S = 1) \;&= b\Big[g^{\ell}_{cuu}(3U_{cs}U_{ud} - U_{cd}U_{us}) - 2g^{\ell}_{ucc}U_{cs}U_{cd} \\ &+ g^{\ell}_{tuu}(3U_{ts}U_{ud} - U_{td}U_{us}) - 2g^{\ell}_{utt}U_{ts}U_{td} \\ &- g^{\ell}_{uct}(U_{ts}U_{cd} + U_{td}U_{cs}) + \frac{1}{2}g'^{\ell}_{uct}(3U_{td}U_{cs} - U_{ts}U_{cd})\Big] \end{aligned} \tag{12a}$$

$$\begin{aligned} A'^L(\bar{\nu}_R; \Delta S = 1) \;&= b\Big[g^{\ell}_{cuu}(3U_{cd}U_{us} - U_{cs}U_{ud}) - 2g^{\ell}_{ucc}U_{cs}U_{cd} \\ &+ g^{\ell}_{tuu}(3U_{td}U_{us} - U_{ts}U_{ud}) - 2g^{\ell}_{utt}U_{ts}U_{td} \\ &- g^{\ell}_{uct}(U_{ts}U_{cd} + U_{td}U_{cs}) + \frac{1}{2}g'^{\ell}_{uct}(3U_{ts}U_{cd} - U_{td}U_{cs})\Big] \end{aligned} \tag{12b}$$

$$
\begin{aligned}
B(\bar{\nu}_R; \Delta S = 1) \quad &= b\Big[-g^{\ell}_{cuu}(U_{cs}U_{ud} - U_{cd}U_{us}) \\
&- g^{\ell}_{tuu}(U_{ts}U_{ud} - U_{td}U_{us}) \\
&+ \frac{1}{2} g'^{\ell}_{uct}(U_{ts}U_{cd} - U_{td}U_{cs})\Big]
\end{aligned}
\tag{12c}
$$

$$
\begin{aligned}
A^L(\bar{\nu}_R; \Delta S = 0) \quad &= b\Big[2g^{\ell}_{cuu}U_{cd}U_{ud} - 2g^{\ell}_{ucc}U^2_{cd} \\
&+ 2g^{\ell}_{tuu}U_{td}U_{ud} - 2g^{\ell}_{utt}U^2_{td} \\
&-2g^{\ell}_{uct}U_{td}U_{cd} + g'^{\ell}_{uct}U_{td}U_{cd}\Big]
\end{aligned}
\tag{12d}
$$

$$
A^L(\ell^+_R; \Delta S = 1) = b\Big[2g^{\ell}_{cuu}U_{cs} + 2g^{\ell}_{tuu}U_{ts}\Big] \tag{12e}
$$

$$
A^L(\ell^+_R; \Delta S = 0) = b\Big[2g^{\ell}_{cuu}U_{cd} + 2g^{\ell}_{tuu}U_{td}\Big] \tag{12f}
$$

where the $U_{qq'}$ are the appropriate entries in the Cabibbo-Kobayashi-Maskawa matrix which we assume [3] to be the same for squarks as for quarks. Evidently the $\Delta S = 0$ decay modes are Cabibbo-suppressed relative to $\Delta S = 1$ modes, irrespective of which of the g^{ℓ}_{abc} may dominate. For the Cabibbo-favoured decay modes $p, n \rightarrow \bar{\nu} + K, \bar{\nu} + K^*, e + K, \mu + K$ and $e + K^*$, the results of Salati and Wallet [12] give

$$
\begin{aligned}
\Gamma(p \rightarrow \bar{\nu} + K^+) \quad &= 1.39 \times 10^{27} \\
&\cdot \Big|A^L(\bar{\nu}_R; \Delta S = 1) - B(\bar{\nu}_R; \Delta S = 1)\Big|^2 y^{-1}
\end{aligned}
\tag{13a}
$$

$$
\begin{aligned}
\Gamma(p \rightarrow \bar{\nu} + K^{*+}) \quad &= 1.23 \times 10^{25} \\
&\cdot \Big|3A^L(\bar{\nu}_R; \Delta S = 1) + 2A'^L(\bar{\nu}_R; \Delta S = 1) \\
&- B(\bar{\nu}_R; \Delta S = 1)\Big|^2 y^{-1}
\end{aligned}
\tag{13b}
$$

$$\Gamma(p \to e^+ + K^0) = 1.39 \times 10^{27} \cdot \left| A^L(e_R^+; \Delta S = 1) \right|^2 y^{-1} \tag{13c}$$

$$\Gamma(p \to \mu^+ + K^0) = 1.34 \times 10^{27} \cdot \left| A^L(\mu_R^+; \Delta S = 1) \right|^2 y^{-1} \tag{13d}$$

$$\Gamma(p \to e^+ + K^{*0}) = 1.23 \times 10^{25} \cdot \left| A^L(e_R^+; \Delta S = 1) \right|^2 y^{-1} \tag{13e}$$

$$\Gamma(n \to \bar{\nu} + K^0) = 1.37 \times 10^{27} \cdot \left| A^L(\bar{\nu}_R; \Delta S = 1) + A'^L(\bar{\nu}_R; \Delta S = 1) \right|^2 y^{-1} \tag{13f}$$

$$\Gamma(n \to \bar{\nu} + K^{*0}) = 1.28 \times 10^{25} \cdot \left| -3A^L(\bar{\nu}_R; \Delta S = 1) - A'^L(\bar{\nu}_R; \Delta S = 1) + 2B^L(\bar{\nu}_R; \Delta S = 1) \right|^2 y^{-1} \tag{13g}$$

Inserting the coefficients (12) into these formulae for the partial rates, we obtain the results of our table showing the baryon decay branching ratios which follow from dominance by the different operators in $\mathcal{L}_{eff}$ (2).* In the absence of strong generation-dependence in the coefficients g, we would expect the operators g^{ℓ}_{cuu} to dominate, since they contribute to baryon decay without Cabibbo-Kobayashi-Maskawa suppression. As mentioned above, we see that $B \to \bar{\nu} + K$ dominates, with an admixture of $p \to \ell^+ + K$ if g^{ℓ}_{cuu} or g^{ℓ}_{tuu} are the dominant operators, and with relatively small branching ratios into $\bar{\nu} + K^*$. It is interesting to note that because the operators (12) do not all have definite strong isospin, the decay rates of the p and n into $\bar{\nu} + K^{+,0}$ (or $\bar{\nu} + K^{*+,0}$) differ in general.

Comparing the decay rates in (13) with the present lower limit [18] on the nucleon lifetime of order 3×10^{30} years, we infer that

* We note in passing that conventional minimal supersymmetric GUTs should have the same pattern of decay modes with $g^{\mu}_{ucc} : g^{e}_{ucc} : g^{\mu}_{cuu} : g^{e}_{cuu} = m_c m_s \sin\theta : m_c m_d \cos\theta : m_u m_s \cos\theta : -m_u m_d \sin\theta$ These results differ from those of ref. [12] because the analogue of equation (4) was misprinted in ref. [3].

$$g^{e,\mu}_{cuu} < O(10^{-6})\frac{1}{m_P} \tag{14}$$

and similarly for the other operator coefficients in equation (2). We do not know *a priori* what the values of these coefficients g might be. Perhaps they should be $O(1/m_P)$, or perhaps they are suppressed by analogy with the trilinear Yukawa couplings of chiral superfields. Equation (14) suggests that gravitodynamic baryon decay could even be overdue in susy theories, whereas it would be unobservably slow in theories without light spin-zero squark and slepton fields. It is perhaps surprising that the expected hierarchy of decay modes in the table is so similar to that coming [3] from conventional minimal susy GUTs, whereas one might have expected that no hard and fast predictions could be made about gravitationally induced baryon decay modes.

It has been suggested [9],[10] that quantum coherence might be lost in gravitationally induced interactions, and it is reasonable to ask whether this could be detected in baryon decay, whether or not a satisfactory framework exists [10] for describing such a breakdown of quantum mechanics. A natural idea is to study the distribution of spins in baryon decay through an experiment analogous to those [19],[20] inspired by Einstein, Podolsky and Rosen [21]. Such an experiment could in principle proceed in the following manner. We will consider for concreteness the case of $p \rightarrow \mu^+ + \rho$. Once the nonrelativistic ρ decays into $\pi\pi$, one can define the $\pi\pi$ axis to be the $\hat{z}$ axis. Since a ρ with spin $S_{\hat{z}} = \pm 1$ cannot give rise to final state mesons whose momenta lie exactly along the $\hat{z}$ axis, one can deduce that the ρ was precisely $S_{\hat{z}} = 0$. However, quantum mechanics is incompatible with the interpretation that the ρ was produced with any definite spin state. It is the decay (or more precisely the observability of the $\pi\pi$ decay products) which determines the spin state precisely. At this point, if the ρ and the muon are produced "coherently" (i.e., in a pure state) the muon also assumes a definite spin state (i.e., "collapses") such that angular momentum is conserved in the quantum mechanical sense. Thus, for example, if the initial baryon state polarization were known and one simply assumes that the baryon decays in an S-wave, one obtains a definite set of predictions for the muon's spin distribution. If, instead, the μ^+ and ρ are produced incoherently and behave afterwards as independent particles (as in thermal radiation from black hole evaporation), one would expect the muon spin wave function to collapse independently of the ρ, which could change the predicted spin distribution. Unfortunately, baryon decay modes such as $B \rightarrow \ell^+ +$ vector meson that have analyzable spin states are Cabibbo-and phase space-suppressed, and have branching ratios of order 1%. Moreover, the experiment requires polarized nucleons and the ability to measure the lepton spins, neither of which seems particularly feasible in a large scale baryon decay experiment. Furthermore, a change in spin distribution would entail a microscopic breakdown of angular momentum conservation in the quantum mechanical sense. A more palatable possibility might be the appearance of interactions which respect quantum mechanical gauge invariance and angular momentum conservation, but are mixed in flavor space. This would imply treating interactions of the type of eq. (1) *incoherently* in the generation indices, and hence discarding interferences between the various g^{ℓ}_{qqq} in eq. (12), instead of adding them coherently. This could lead to baryon decay branching ratios different from those obtained from eqs. (12) and (13) if two or more of the interactions have comparable magnitude. Given the current stage-fright of unstable baryons, it is unlikely that the precision measurements required to discern such an effect will be forthcoming. However, in addition to these practical problems, we have the following theoretical reasons to believe that quantum mechanics would not be violated observably in baryon decay.

In the same way that one might expect interactions like (1) to be generated, one might also expect [11],[22] to encounter interactions like

$$O\left(\frac{1}{m_P}\right)(FH\Sigma\ T)\ ,\ O\left(\frac{1}{m_P}\right)(TT\Sigma\ H) \tag{15}$$

where H, $\bar{H}$ and Σ are respectively a $\underline{5}$, $\bar{\underline{5}}$ and $\underline{24}$ of Higgs in SU(5). Putting in $\langle 0|H,\bar{H}|0\rangle = 0(100)GeV$ and $\langle 0|\Sigma|0\rangle = O(10^{16})GeV$ one easily gets from (15) contributions to the e, d and u masses which are of the same order as their actual values [11],[22]. However, we know that non-relativistic electrons and neutrons obey quantum mechanics with very high precision. Therefore there are severe upper limits on the amount of incoherence that can be introduced by gravitational effects analogous to mass terms arising from interactions such as (15). The best limit may come from experiments [23] on quantum interference using slow neutrons. Coherence could not be maintained over the times and distances used in these experimental tests of quantum mechanics unless

$$\frac{\delta m_{incoherent}}{m_n} \leq O(10^{-18}) \tag{16}$$

which can be translated into

$$\delta m_q(incoherent) \leq O(10^{-18}) GeV\ . \tag{17}$$

A similar, if somewhat less restrictive, constraint follows from the coherent behavior of kaons in $K^0 - \overline{K^0}$ mixing. We therefore deduce from (17) that gravitodynamic interactions such as (15) must be predominantly coherent if they are present at anywhere near the expected [22] level. One may expect the same to be true of the baryon decay operator (1): if it is observable it will probably be coherent. Note that although our remarks about quantum mechanics were motivated by the suggestion [9],[10] that non-renormalizable interactions scaled by $O(1/m_P)$ might not be coherent, the restrictions (16) and (17) also constrain severely any incoherent contribution to conventional renormalizable interactions including Yukawa or gauge couplings. We should mention that some arguments in ref. [10] suggest that some incoherent non-perturbative gravitational effects at low energies E may be suppressed by additional powers of $O(E/m_P) = O(10^{-19})$. * In this case, a modest improvement of the limit (16) by just a few orders of magnitude could be very revealing.

In view of the interest [9],[10] in a possible breakdown of quantum mechanics at the Planck scale, we think that tests of quantum coherence should be developed to the greatest possible extent, and improved quantum interference experiments with slow particles propagating over long distances seem particularly desirable.

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* This suppression might be absent in a more realistic theory where the space-time foam is supersymmetric. Also, we know there are matter interactions violating global symmetries, which could for example generate $O(1/m_P)$ interactions via radiative corrections to the effects discussed in ref. [10].

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TABLE I

Relative Decay Rates

Dominant Operator	$p \to \bar{\nu}K^+$	$p \to e^+K^o (\mu^+K^o)$	$p \to \bar{\nu}K^{*+}$	$p \to e^+K^{*o}$	$n \to \bar{\nu}K^o$	$n \to \bar{\nu}K^{*o}$
$g_{cuu}^{e(\mu)}$ (s^0), $g_{tuu}^{e(\mu)}$ (s^2)	1	$\frac{1}{4}$	$\mathcal{O}\left(\frac{1}{25}\right)$	$\mathcal{O}\left(\frac{1}{400}\right)$	$\frac{1}{4}$	$\mathcal{O}\left(\frac{1}{16}\right)$
$g_{ucc}^{e(\mu)}$ (s^2), $g_{uct}^{e(\mu)}$ (s^4), $g_{utt}^{e(\mu)}$ (s^6)	$\frac{1}{4}$	0	$\mathcal{O}\left(\frac{1}{16}\right)$	0	1	$\mathcal{O}\left(\frac{1}{25}\right)$
$g_{uct}^{\prime e(\mu)}$ (s^4)	$\mathcal{O}(1)$	0	$\mathcal{O}\left(\frac{1}{100}\right)$	0	$\mathcal{O}(1)$	$\mathcal{O}\left(\frac{1}{100}\right)$

The decay rates into K^* are suppressed by kinematic factors. The relative decay rates for the $g_{uct}^{\prime\ell}$ case depend on unknown Cabibbo-Kobayashi-Maskawa mixing angles. Indicated in parentheses at the left are the number of small angle factors suppressing the decay rates due to the different operators. Decay rates into non-strange final states are suppressed by an additional factor of $sin^2\theta_c$.