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RARE K DECAYS AND NEW PHYSICS BEYOND THE STANDARD MODEL [a]

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Rare decays provide a complementary window to direct collider searches as probes of physics beyond the Standard Model. We present an overview of the New Physics sensitivity provided by existing and future measurements of rare leptonic and semileptonic K decays within the context of several classes of models.

1 Introduction

Rare K decays have historically provided strict constraints on the construction of New Physics models. Amongst these many decays those involving leptons are the most theoretically clean. Here we will provide an overview of what these modes are telling us about the parameter spaces of several New Physics scenarios. We can categorize these rare K decays into several distinct classes, which we will discuss in turn, as follows:

- Non-Standard Model Modes. These consist of a handful of decays such as $K_L \to e\mu$, $K \to \pi e\mu$ and $K \to \pi X$ which just don't fly in the Standard Model (SM). The simple observation of *any* of these modes would signal New Physics.

- Precision Modes. These decay modes occur in the SM but have associated with them T-odd correlations which possibly lead to CP-violating observables that can only be probed by high precision experiments. Examples are the measurements of the transverse μ polarization in either $K \to \pi\mu\nu$ or $K \to \mu\nu\gamma$.

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- Short Distance Dominated Modes. These particularly clean modes occur in the SM but rate enhancement or suppression would signal New Physics. $K^+ \to \pi^+ \nu\nu$ and the CP-violating $K_L \to \pi^0 \nu\nu$ channels provide the best examples and are very well understood in the SM.

- Long-Distance Dominated Modes. These include such decays as $K \to \ell^+\ell^-$, $K \to \ell^+\ell^-\gamma$ and $K \to \pi\ell^+\ell^-$ which provide fertile testing grounds for chiral perturbation theory(ChPT) [1] but which have somewhat diminished sensitivity to New Physics due to long distance uncertainties. (Perhaps $K \to \mu^+\mu^-$ is an exception.) We will have nothing to say about these modes in the discussion below.

2 Non-Standard Model Modes

Both $K_L \to e\mu$ and $K \to \pi e\mu$ proceed through operators of the form

$$\mathcal{O}_{V,A} = \frac{g_X^2}{2M_X^2}\bar{d}\gamma_\mu[C_{Lq}P_L + C_{Rq}P_R]s \cdot \bar{\mu}\gamma^\mu[C_{L\ell}P_L + C_{R\ell}P_R]e + h.c.\,,$$
$$\mathcal{O}_{S,P} = \frac{g_X^2}{2M_X^2}\bar{d}[C'_{Lq}P_L + C'_{Rq}P_R]s \cdot \bar{\mu}[C'_{L\ell}P_L + C'_{R\ell}P_R]e + h.c.\,, \qquad (1)$$

but are complementary in that $< 0|\mathcal{O}_{A,P}|K >\neq 0$ and $< \pi|\mathcal{O}_{V,S}|K >\neq 0$ with all other matrix elements being zero. As it stands, the form of the above operators conserve generation number. It is important to remember that other operators of the same kind may exist wherein the roles of e and μ are interchanged so that generation number is no longer conserved. In principle, all operators such as those above may be generated either through loops or via tree-level exchanges. Branching fractions for the above decays can be immediately calculated from these operators and the well-known SM matrix elements apart from small isospin corrections, *e.g.*, the branching fraction for $K_L \to e\mu$ due to $\mathcal{O}_A$ is

$$B_A = 11.24 \cdot 10^{-12}\ [\frac{g_X}{g}\frac{100TeV}{M_W}]^4\ (C_{Lq} - C_{Rq})^2\ (C_{L\ell}^2 + C_{R\ell}^2)\,, \qquad (2)$$

with qualitatively similar results holding for the other modes and operators, apart from chiral enhancement factors in the cases of S- and P-type couplings. These results are summarized in Fig. 1 assuming $g_X = g$, the weak coupling constant, and the product of the coefficients C_i's is set to unity. Since the present limit [2] on the branching fraction for $K_L \to e\mu$ is $3 \cdot 10^{-12}$ at 90% CL, we see that for tree-level exchanges with weak couplings, mass scales $>$ 100 TeV, beyond the reach of any planned collider, are already being probed.

Somewhat lower scales, but no less impressive, are probed by the corresponding $K_L \to \pi^0 e\mu$ and $K^+ \to \pi^+ e\mu$ modes with branching fraction upper limits [2] of $3.2 \cdot 10^{-9}$ and $4 \cdot 10^{-11}$, respectively, at the 90% CL. Unfortunately, due to the $\sim M_X^4$ scaling of the branching fractions for both $K_L \to e\mu$ and $K \to \pi e\mu$ decay modes, the mass range being probed will improve quite slowly as bounds on these branching fractions improve.

R-parity violating interactions that also violate lepton number can mediate both $K_L \to e\mu$ and $K \to \pi e\mu$ processes through tree-level exchanges. The relevant trilinear terms in the superpotential are given by $W_R = \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c$ (where i, j, k are family indices and symmetry demands that $i < j$ in the term proportional to λ). In the case of, *e.g.*, $K_L \to e\mu$, the reaction occurs through both s-channel $\tilde{\nu}$ exchange(which induces S,P-type operators) as well as t-channel $\tilde{u}$ exchange(which induces V,A-type operators via a Fierz transformation)[3]. Interestingly if the $\tilde{\nu}$ exchange dominates, both the processes $K_L \to e^+\mu^-$ and $K_L \to e^-\mu^+$ can occur but need not have equal branching fractions. If the $\tilde{\nu}(\tilde{u})$ channel dominates then the very strong existing bound on the branching fraction severely constrains the product of the relevant Yukawa couplings to be $< 7 \cdot 10^{-9}(1 \cdot 10^{-7})$ for 100 GeV SUSY particle masses.

Similarly, leptoquark(LQ) exchange [4] can also be responsible for these lepton number violating processes by inducing V,A-type operators. Using the current branching fraction bound results in a roughly model-independent limit of $M_{LQ} \geq 150\sqrt{\frac{\lambda\lambda'}{g^2}}$ TeV where $\lambda\lambda'$ is the product of the relevant Yukawas. This is one of the major reasons for believing that LQ couplings are family diagonal.

The two-body decay $K \to \pi X$ where X is invisible can occur in many scenarios, *e.g.*, the familon model[5]. If one describes such decays via an effective interaction of the form

$$\mathcal{L} = \frac{1}{F}\partial^\mu X(\bar{u}\gamma_\mu s) + h.c.\,, \tag{3}$$

then the lack of observation of this decay by E-787[6] at the 90% CL of $3.0 \cdot 10^{-10}$ implies $F > 3.1 \cdot 10^{11}$ GeV, an impressively large scale.

3 Precision Modes

In the absence of final state interactions(FSI) the T-odd transverse μ polarization, $P_\mu^T \sim \mathbf{s}_\mu \cdot (\mathbf{p}_{\pi,\gamma} \times \mathbf{p}_\mu)$ in either $K \to \pi\mu\nu$ or $K \to \mu\nu\gamma$ is a new measure of CP violation as it is effectively zero in the SM. Fortunately, the effects of FSI have been shown to be reasonably small (or at least not dangerously large) [7] particularly for the $K \to \pi$ case. For $K \to \pi\mu\nu$ one expects FSI at the level of

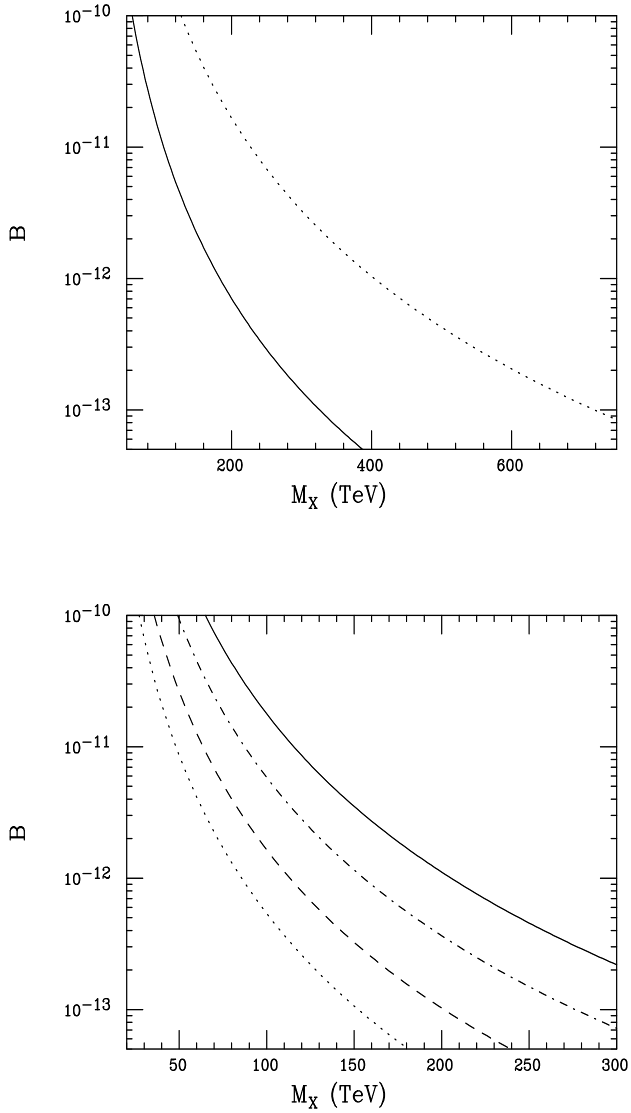

Figure 1: Branching fractions for $K_L \to e\mu$(top) and $K \to \pi e\mu$(bottom) as a function of M_X assuming $g_X = g$. In the top panel the solid(dotted) curve corresponds to A(P) exchange. In the bottom panel, the K_L mode is represented by the solid(S) and dashed(V) curves while the K^+ mode is represented by the dash-dotted(S) and dotted(V) curves.

$\sim 10^{-6}$ whereas for $K \to \mu\nu\gamma$ they are expected to be $\sim 10^{-3}$. Both processes can be analyzed in a model-independent manner using the formalism of Wu and Ng [8].

The most general dimension-six interaction responsible for $K \to \pi\mu\nu$ assuming only left-handed neutrinos can be written as

$$\mathcal{L} = \mathcal{L}_{SM} + 2[\bar{u}(G_S + G_P\gamma_5)s](\bar{\mu}P_L\nu) + 2[\bar{u}\gamma_\mu(G_V + G_A\gamma_5)s](\bar{\mu}\gamma^\mu P_L\nu)\,, \quad (4)$$

but the A, P terms do not contribute to this decay having zero matrix elements due to parity conservation. (They will contribute in the case of the $K \to \mu\nu\gamma$ decay.) From $\mathcal{L}$ one obtains an effective Hamiltonian which has the same structure as in the SM:

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}} V_{us}\bar{u}(\mu)\gamma^\lambda(1-\gamma_5)v(\nu)\big[f'_+(p_K+p_\pi)_\lambda + f'_-(p_K-p_\pi)_\lambda\big]\,, \quad (5)$$

where $f'_\pm = f_\pm(1+\delta_\pm)$, $f_\pm$ being the well-known SM form-factors and

$$\delta_+ = \frac{\sqrt{2}G_V}{G_F V_{us}}\,,$$
$$\delta_- = \delta_+ + \frac{\sqrt{2}G_S m_K^2}{G_F V_{us}(m_s - m_d)m_\mu}\big[\frac{f_+}{f_-}(1-r) + \frac{q^2}{m_K^2}\big]\,, \quad (6)$$

with $r = (m_\pi/m_K)^2$. A short analysis then shows that for this mode $P_\mu^T(\pi) \sim Im(\xi) = Im[f'_-/f'_+] \sim ImG_S$, so that only this quantity needs to be calculated in any given model [9]. If G_S cannot be generated or if it is real then $P_\mu^T(\pi)$ remains zero. The current experimental results summarized by the PDG [10] are $P_\mu^T(K^+) = -0.0031 \pm 0.0053$ and $P_\mu^T(K^0) = 0.0021 \pm 0.0048$. In addition, KEK-246 has recently reported [11] a new preliminary result of $P_\mu^T(K^+) = -0.0025 \pm 0.0058$. In the future, KEK-246 expects to reach an experimental sensitivity of 0.0013 whereas AGS-936 and AGS-923 expect to reach sensitivities of 0.00035 and 0.00013, respectively [12].

$K \to \mu\nu\gamma$ can be analyzed in a similar manner. The matrix element has two contributions, one due to Inner Bremsstrahlung(IB) where the photon comes off one of the charged legs or the $K\mu\nu$ vertex, and a Structure Dependent(SD) term arising from loops. THese can be written symbolically as

$$M_{IB} \sim f'_K K^\mu \epsilon^*_\mu\,,$$
$$M_{SD} \sim L_\nu H^{\mu\nu}\epsilon^*_\mu\,, \quad (7)$$

with L_ν and K^μ being vectors formed from the leptonic momenta, ϵ_μ is the photon polarization vector and, with $p(q)$ being the $K(\gamma)$ momenta, the hadronic tensor is

$$H^{\mu\nu} = F'_A\big[-g^{\mu\nu}p\cdot q + p^\mu q^\nu\big] + iF'_V\epsilon^{\mu\nu\tau\sigma}q_\tau p_\sigma\,. \qquad (8)$$

The New Physics is encoded in the primed quantities: $f'_K = f_K(1-\Delta_P-\Delta_A)$, $F'_A = F_A(1-\Delta_A)$ and $F'_V = F_V(1+\Delta_V)$, with the unprimed quantities corresponding to their SM values. ($F_V = -0.0945$ and $F_A = -0.0425$ in chiral perturbation theory at the one-loop level [1].) The Δ_i are given by

$$\Delta_{P,A,V} = \frac{\sqrt{2}}{G_F V_{us}}\left[\frac{G_P m_K^2}{(m_s+m_d)m_\mu},\ G_A,\ G_V\right], \qquad (9)$$

with the G_i as defined in the effective Lagrangian above. A short analysis then shows that $P_\mu^T(\gamma) = \mathcal{F}_1 Im(\Delta_A+\Delta_V)+\mathcal{F}_2 Im(\Delta_P)$ where $\mathcal{F}_{1,2}$ are phase space factors of order ~ -0.1. Here we note that G_S does *not* contribute; nor will a common overall shift in the V-A coupling since then the sum $\Delta_R = \Delta_V + \Delta_A$ will be zero. Thus, for example, in the Left-Right Symmetric Model [13], $P_\mu^T(\pi) = 0$ but $P_\mu^T(\gamma) \sim -2\cdot 10^{-3}\ [\phi/10^{-2}]\ Im[e^{i\omega}V_{us}^R/|V_{us}^L|]$, where $\phi(\omega)$ is the $W_L - W_R$ mixing angle(phase) and $V^{L,R}$ are the left- and right-handed CKM matrices which need not be directly related. As shown by Wu and Ng [8,9], a significant Δ_R can also be generated by large stop mixing. In models with leptoquarks, R-parity violation or an enriched Higgs sector [8,9,14], one generates both $Im(G_{S,P}) \neq 0$ with comparable magnitudes leading to values of $P_\mu^T(\pi,\gamma)$'s in the $few\cdot 10^{-3}$ range and with the same sign. This level of polarization should be observable in the future.

This discussion demonstrates that the two modes $K\to\pi\mu\nu$ and $K\to\mu\nu\gamma$ are complementary in their sensitivity to operators with different tensor structures generated by New Physics.

4 Short-Distance Dominated Modes

$K^+\to\pi^+\nu\bar\nu$ and $K_L\to\pi^0\nu\bar\nu$ are two of the most well understood rare decays in the SM [15] with anticipated branching fractions of $(9.1\pm3.2)\cdot10^{-11}$ and $(2.8\pm1.7)\cdot10^{-11}$, respectively. (Note most of the uncertainty in these values arises from our poor knowledge of the input parameters and not from theoretical uncertainties.) E-787 [6] has recently obtained evidence for the charged decay mode with a branching fraction of $B = (4.2^{+9.7}_{-3.5})\cdot10^{-10}$, consistent with the SM but leaving room for New Physics. In the case of the neutral mode there is presently only an upper bound [16] from KTEV which is $B < 1.6\cdot10^{-6}$ at 90% CL.

In almost all models with only left-handed ν's the effective Hamiltonian describing this decay can be written as

$$\mathcal{H}_{eff} = \frac{2G_F}{\sqrt{2}} \frac{\alpha}{\pi \sin^2\theta_W} \lambda_t \bar{s}\gamma^\mu (X_L P_L + X_R P_R) d \; \bar{\nu}_\ell \gamma_\mu P_L \nu_\ell \; + \; h.c. \, , \qquad (10)$$

with $\lambda_i = V_{is}^* V_{id}$ and $X_{L,R}$ being model-dependent. (Note that it is possible that more general coupling structures which include scalar and pseudoscalar terms may be generated via box diagrams in some more exotic models not discussed below.) In the SM, $X_R = 0$ and $X_L = X_t + \frac{\lambda^4 \lambda_c}{\lambda_t} P_c$, with the top contribution being $X_t \sim 1.5$, λ being the usual Wolfenstein parameter and $P_c = 0.40 \pm 0.06$ being the highly suppressed charm contribution. From this Hamiltonian one obtains the following branching fractions:

$$B(K^+ \to \pi^+ \nu\bar{\nu}) = 4.11 \cdot 10^{-11} \Big| \frac{\lambda_t (X_L + X_R)}{\lambda^4} \Big|^2 \, ,$$
$$B(K_L \to \pi^0 \nu\bar{\nu}) = 1.80 \cdot 10^{-10} \Big[Im \frac{\lambda_t (X_L + X_R)}{\lambda^4} \Big]^2 \, , \qquad (11)$$

and thus only the additional contributions to $X_{L,R}$ need to be calculated within a given New Physics model to obtain constraints.

Let us quickly survey a few New Physics models which may modify the branching fractions for these decays. (i) One simple scenario is that of a fourth generation [17] t' which gives a contribution $\Delta X_L = \frac{\lambda_{t'}}{\lambda_t} X_{t'}$ with $1.5 \leq X_{t'} \leq 5$, depending on the mass of t'. This implies that $\frac{\lambda_{t'}}{\lambda_t} \geq \lambda^2 - \lambda$ to make any reasonable contribution and, hence, is sensitively dependent on what is assumed for the generalized CKM matrix structure. (ii) A second example is the set of Two-Higgs Doublet Models [18] with natural flavor conservation. Here it is well-known that the new contributions to X_L beyond the SM are small, of order a few per cent, for all charged Higgs masses and values of $\tan\beta$ when other experimental constraints are imposed on the model. (iii) Various anomalous WWZ couplings consistent with LEP and Tevatron direct bound measurements have been shown to make only a small [19] contribution to X_L. (iv) Any leptoquark [4] mediating these decay modes through tree-level exchange has been shown to be quite heavy, *i.e.*, $M_{LQ} > (25-50)\sqrt{\lambda\lambda'}$ TeV, depending on the detailed nature of the leptoquark. (v) R-parity and lepton-number violating $\tilde{d}_{L,R}$ exchanges in the t-channel have been shown [3] to contribute to both $X_{L,R}$ and allow for final states with mixed neutrino flavors so that the $X_{L,R}$'s would in this case pick up a pair of generation labels! For $\tilde{d}$ masses of order 100 GeV, the E-787 result forces the products of the relevant Yukawa couplings in this case to be $\leq 10^{-(4-5)}$.

As a last example we consider the case of R-parity conserving SUSY [20]. There are many potential contributions to the decay amplitude to consider arising from both boxes and Z-penguins containing SUSY partners in addition to those arising from the usual charged Higgs/top exchange. These include the intermediate states of (i) $\tilde{d}_i \oplus \tilde{g}$, (ii) $\tilde{d}_i \oplus \tilde{\chi}^0_j$ and (iii) $\tilde{u}_i \oplus \tilde{\chi}^+_j$. In addition, box diagrams will now also involve $\tilde{\ell}$ and $\tilde{\nu}$ exchanges but $\tilde{g}$ boxes are absent due to color conservation. As is well-known [20], unlike in the case of $|\Delta S|=2$ transitions, the chargino-squark contributions are expected to be dominant. Furthermore, penguins are generally found to dominate over contributions from boxes.

Due to the complexity of the general SUSY parameter space certain approximations are employed in evaluating the various contributions; these center on the structure of the 6×6 squark mass matrices. Conventionally, these matrices are written down in the so-called super-CKM basis in which the flavor structure of the quark-squark-gaugino vertices is the same as in quark-quark-gauge boson vertices. As discussed [20] by Buras *et al.*, the calculation simplifies by a change to a different basis where $d^i_L \tilde{d}^j_L \chi^0_n$ and $d^i_L \tilde{u}^j_L \chi^\pm_n$ are flavor diagonal and the $d^i_L \tilde{u}^j_R \chi^+_n$ flavor structure is described by the CKM matrix. This basis is arrived at from the super-CKM basis by rotating the up-type left-handed squark fields as $\tilde{u}_L \to V^\dagger \tilde{u}_L$, with V bing the CKM matrix. In this case the flavor change in the left-handed sector occurs through the non-diagonality of the sfermion propagators. If the off-diagonal terms(Δ) in the sfermion mass matrices are small in comparison to the average sfermion mass, $\tilde{m}$, it is possible to expand these fermion propagators in powers of $\delta = \Delta/\tilde{m}^2$. The remaining issue is then to know how many powers of δ to keep in the evaluation of the various diagrams. Until recently, it has been believed (see Nir and Worah and Buras *et al.* [20]) that a single mass-insertion approximation [21] is sufficient for obtaining reasonable estimates for the SUSY contributions to the $K^+ \to \pi^+ \nu\bar{\nu}$ and $K_L \to \pi^0 \nu\bar{\nu}$ decay amplitudes. In that approximation, it has been shown that departures by factors of $2 \sim 3$ from the SM decay rate expectations are possible once other constraints on the SUSY parameter space are taken into account.

More recently, Colangelo and Isidori [20] have argued that this single mass insertion approximation is not sufficient to account for all potentially large SUSY effects and that one must include at least two mass insertions. (In their approach one also perturbatively expands the matrices responsible for diagonalizing the chargino mass matrix to keep track of powers of $SU(2)_L$ breaking terms.) They find that the dominant SUSY contribution to X_L can be quite large and is given by $\simeq \frac{1}{96}\frac{\tilde{\lambda}_t}{\lambda_t}$ where $\tilde{\lambda}_t = (\delta^U_{LR})^*_{ts}(\delta^U_{LR})_{td} = |\tilde{\lambda}_t| e^{i\tilde{\theta}_t}$ and the δ's are given by the LR elements of the u-squark mass matrix in the new basis: $(\delta^U_{LR})^*_{ab} = [M(LR)^2_U]_{ab}/m^2_{\tilde{u}_L}$. These authors then examine all of the other low-energy constraints on the parameter $\tilde{\lambda}_t$ from K, D, and B physics while simultaneously evaluating its maximum contribution to X_L. One finds that the branching fraction for $K^+ \to \pi^+ \nu\bar{\nu}$ can be increased by as much as an order of magnitude while that for $K_L \to \pi^0 \nu\bar{\nu}$ depends strongly on the new physics phase $\tilde{\theta}_t$. For large values of $\tilde{\theta}_t$ near 90^o the enhancement in the branching fraction for this mode may be almost as large as two orders of magnitude, which is a very exciting prospect.

It is clearly quite important to verify and improve upon the E-787 result for the $K^+ \to \pi^+ \nu\bar{\nu}$ branching fraction and to try to observe the decay $K_L \to \pi^0 \nu\bar{\nu}$ as soon as possible.

5 Summary

Rare K decays involving leptons do indeed provide a complementary window to direct collider searches as probes of New Physics beyond the Standard Model.

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