**Theoretical Issues in Rare $K$, $D$ and $B$ Decays**

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**Abstract.** I discuss some theoretical aspects of rare $K$, $D$ and $B$ decays focusing mainly on their potential as tests of the one loop structure of the standard model. I concentrate on flavor changing neutral current processes and compare our ability to extract short distance physics in the three cases. Finally, I give some examples of the sensitivity of these decays to extensions of the standard model.

**INTRODUCTION**

The continuing success of the standard model (SM) has turned into a challenge both for theorists and experimentalists alike. On the one hand, theorists believe that the SM picture of the Higgs mechanism based on one elementary scalar doublet is unnatural, trivial and has come to be viewed as an effective description of a more complicated Higgs sector, one involving perhaps additional scalars and/or fermions or even new gauge interactions. In addition, the idea that fermion masses arise as a result of the interactions with this elementary scalar, requires for instance that its dimensionless couplings to two otherwise identical fermions such as the up and the top quarks, differ by more than four orders of magnitude. Understanding the mechanism for electroweak symmetry breaking (EWSB) and the origin of fermion masses calls for physics beyond the SM. At least in the case of EWSB, it is understood that this new physics must reside at an energy scale not far beyond 1 TeV. The experimental challenge of finding new physics in direct searches may still take some time if the new states or their effects only set in at several hundred GeV. A complement of these direct signals at the highest available energies is the measurement of the effects of the new particles in loops, either through precision measurements such as the ones performed at LEP, or through the detection of processes only occurring at one loop in the SM. Among these are the transitions induced by flavor changing neutral currents (FCNC), such as $K^0\rightarrow\bar{K}^0$ or $b\rightarrow s\gamma$. These are forbidden at tree level in the SM due to the presence of the

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1) Talk presented at the Workshop on Heavy Quarks at Fixed Target, Fermilab, Batavia IL, October 10-12, 1998.
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Glashow-Iliopoulos-Maiani (GIM) mechanism. As a consequence, these processes are highly suppressed in the SM. Thus the one loop effect of a new heavy state may translate into a large effect in a branching ratio if this loop induces a FCNC transition. Here I discuss the FCNC decays of $K$, $D$ and $B$ and their potential as tools in searching for new physics at high energy scales $\gtrsim M_W$. One critical aspect in evaluating this potential is the extent to which a given process is determined by high energy scales, i.e. short distance physics, or it is contaminated by the more mundane effects of long distance dynamics such as the ones induced by propagating intermediate hadrons. Since the latter are not calculable in perturbation theory, the decay modes affected by long distance physics are less understood. Next, I will give examples of short and long distance contributions in the FCNC decays of $K$, $D$ and $B$ mesons. Once established which modes are the most likely to be testing grounds of the SM, I will turn to the discussion of two typical examples of its extensions: supersymmetry and the effects of anomalous triple gauge boson couplings.

**RARE $K$ DECAYS**

Let us start the discussion with the decay $K^+ \to \pi^+ e^+ e^-$. In the SM this process receives one loop contributions from electroweak penguin and box diagrams involving up-type quarks. These are truly short distance diagrams and although some long distance physics enters in the hadronization from $s \to K$ and from $d \to \pi$ in the way of form-factors, the theoretical uncertainties this introduces are not by themselves large enough to obscure the interesting physics. However, the process $K^+ \to \pi^+ \gamma^*$ followed by $\gamma^* \to e^+ e^-$ gives another contribution, reflecting the long distance dynamics of the $K^+ \pi^+ \gamma$ vertex, and dominates the rate [1]. We conclude that this decay mode is not well suited for a test of the short distance physics entering the loops.

We now turn to consider $K_L \to \pi^0 e^+ e^-$. The short distance contributions to this process involve direct CP violation. This is extremely interesting since the measurement of this decay rate could be in principle a direct determination of the CP violating parameter $\eta$, the complex phase in the Cabibbo-Kobayashi-Maskawa (CKM) matrix. Furthermore, the one photon intermediate state cannot contribute here since the long distance vertex to one photon is CP conserving. Still, there is contamination from long distance physics. This comes from two sources. First, there is indirect CP violation given essentially by

$$BR(K_L \to \pi^0 e^+ e^-)_{ICP} = |\epsilon_K|^2 \frac{\tau(K_L)}{\tau(K_S)} BR(K_S \to \pi^0 e^+ e^-),$$

where the $K_S$ decay, being CP conserving, is dominated by the one photon intermediate state. In addition, there is a CP conserving contribution to the decay rate from $K_L \to \pi^0 \gamma^* \gamma^*$ which gives a non interfering term resulting in

$$BR(K_L \to \pi^0 e^+ e^-)_{CPC} \simeq (1 - 2) \times 10^{-12},$$

although this could be larger [2]. The
indirect CP violating piece will be well known once the $K_S$ decay mode is measured. Progress to a better understanding of the CP conserving piece may be achieved in the near future. Thus, there is reasonable hope that in this process we may be able to disentangle the short and long distance physics, at least to some extent. For the moment, it remains in the list of long distance “polluted” modes.

Finally, we turn to the decay modes with neutrinos replacing the charged leptons in the final state. The processes $K^+ \rightarrow \pi^+\nu\bar{\nu}$ and $K_L \rightarrow \pi^0\nu\bar{\nu}$ are almost completely determined by short distance physics. The effective Hamiltonian for the charged mode is

$$
H_{\text{eff}} = \frac{4G_F}{\sqrt{2}} \frac{\alpha}{2\pi s^2 \theta_W} \sum_{\ell} \left( \lambda_c X_c^\ell + \lambda_t X(x_t) \right) \left( \bar{s}_L \gamma_\mu d_L \right) \left( \bar{\nu}_L \gamma^\mu \nu_L \right)_{\ell}, \tag{2}
$$

where $\lambda_i = V_{ls}^* V_{id}$, ($\ell = e, \mu, \tau$), $X(x_t)$ is the result of the top quark loop contribution, and $X_c^\ell$ is the charm quark contribution and carries a dependence on the lepton flavor coming from the box diagram with charged leptons. The contribution from the charm loop, in addition to the charm quark mass dependence, introduces a sizeable scale dependence which is only reduced when including next-to-leading-order corrections. The remnant dependence on the scale $m_c$ entering in the charm running mass results in an uncertainty $<10\%$ in the $BR(K^+ \rightarrow \pi^+\nu\bar{\nu})$. The hadronic matrix element needed to compute the matrix element of the exclusive mode can be obtained by isospin rotation from $K^+ \rightarrow \pi^0\ell^+\nu$ and therefore does not introduce additional hadronic uncertainties. The SM prediction for the branching ratio is obtained for $m_c = 1.3$ GeV, $m_t = 170$ GeV, $V_{cb} = 0.04$ and $V_{ub} = 0.032$, leading to $[3] \ Br(K^+ \rightarrow \pi^+\nu\bar{\nu}) = (9.1 \pm 3.8) \times 10^{-11}$, where the uncertainty is mainly from the CKM parameters. The recent observation of one event in this channel by BNL E787, translates into $[4] \ Br(K^+ \rightarrow \pi^+\nu\bar{\nu})_{\text{exp}} = (4.2 + 9.7 - 3.5) \times 10^{-10}$.

Finally, the neutral mode $K_L \rightarrow \pi^0\nu\bar{\nu}$ is even cleaner than the charged one due to the fact that it is largely dominated by direct CP violation. As a result, only the top quark loop in Eq. (2) contributes so the uncertainties associated with the charm scale are not present. Furthermore, this mode constitutes a very unique and direct way to measure the CP violating phase of the CKM matrix. The amplitude depends on $Im[\lambda_t]$, which in the Wolfenstein parametrization can be written in term of the CP violating term $\eta$ and $V_{cb}$, giving $[3]$

$$
Br(K_L \rightarrow \pi^0\nu\bar{\nu}) = 3.29 \times 10^{-5} \eta^2 |V_{cb}|^4 X^2(x_t), \tag{3}
$$

with $\eta \equiv Im[V_{ub}^* / V_{cb}] / \lambda$. The branching ratio of the neutral mode is still too small in the SM when compared with current experimental limits. Future experiments expect to be sensitive to a few $\times 10^{-11}$ branching fraction.

There seems to be a compromise between experimental accessibility and the theoretical uncertainties in any of the modes discussed above: the most accessible modes tend to be affected by larger theoretical uncertainties. This tension is always present in rare FCNC decays. In the case of $K$ decays, the neutrino modes are hard but still accessible experimentally. The charged kaon mode seems to be a good
compromise, since is a short distance dominated process and the uncertainties seem to be surmountable. The additional interest of the neutral mode is the observation of direct CP violation and the direct measurement of CKM parameters.

**RARE D VS. B DECAYS**

We now turn to a comparative discussion of FCNC in $D$ and $B$ decays. The essential aspects can be framed as external-up-quark vs. external-down-quark processes. We will make the comparison using the radiative processes $c \rightarrow u\gamma$ vs. $b \rightarrow s\gamma$, for the sake of simplicity. Most of the conclusions can be extended to other modes. The short distance contributions to the radiative FCNC process $Q \rightarrow q\gamma$ result in the decay width

$$
\Gamma^{(0)} = \frac{\alpha G_F^2 m_Q^5}{128 \pi^4} \left| \sum \lambda_i F(x_i) \right|^2,
$$

where the superscript “0” denotes the absence of QCD corrections, $Q = (c, b)$, $x_i = (m_q/M_W)^2$, and the function $F(x)$ is the result of integrating the loop contribution of the internal quark $i$. This loop function is the same in the $c$ and $b$ cases. The main difference in Eq. (4) comes from the masses of the internal quarks and the CKM factors $\lambda_i$. In order to see how this affects the widths we turn to Table I, where we show separately the contribution of each quark flavor.

|     | $i$ | $F(x_i)$  | $\lambda_i F(x_i)$   |
|-----|-----|-----------|----------------------|
| $c \rightarrow u\gamma$ | d   | $1.6 \times 10^{-9}$ | $3.4 \times 10^{-10}$ |
|     | s   | $2.9 \times 10^{-7}$ | $6.3 \times 10^{-8}$ |
|     | b   | $3.3 \times 10^{-4}$ | $3.2 \times 10^{-8}$ |
| $b \rightarrow s\gamma$ | u   | $2.3 \times 10^{-9}$ | $1.3 \times 10^{-12}$ |
|     | c   | $2.0 \times 10^{-4}$ | $7.3 \times 10^{-6}$ |
|     | t   | 0.4       | $1.6 \times 10^{-2}$ |

Table I: Contributions to $Q \rightarrow q\gamma$. From Ref. [5].

The CKM factors are $\lambda_i = V_{ci}^* V_{ui}$ for $c \rightarrow u\gamma$, and $\lambda_i = V_{bi}^* V_{si}$ for $b \rightarrow s\gamma$. As we can see, all three contributions are small in the $c \rightarrow u\gamma$ case, whereas for $b \rightarrow s\gamma$ the top quark loop gives the overwhelmingly dominant piece. The central point is that heavier quarks give the dominant contributions as long as their mixing with the external quarks is not highly suppressed. This is a consequence of the non-decoupling aspect of the SM, the fact that fermions that acquired masses from the Higgs mechanism do not decouple in loops involving the massive electroweak gauge bosons. In $c \rightarrow u\gamma$ the internal $b$ quark contribution would dominate if it was not for the fact that $V_{cb}$ and $V_{ub}$ are extremely small. In any event, the QCD
uncorrected $c \to u\gamma$ rate is very small due to the CKM dominance of the lighter intermediate states $d$ and $s$. The $b \to s\gamma$ width is large due to the presence of a heavy top!

Although the QCD corrections to Eq. (4) are generally important, their impact also varies depending on the intermediate mass in the loop. They enhance the $c \to u\gamma$ rate by five orders of magnitude on the one hand, but the $b \to s\gamma$ rate goes up by less than a factor of three or so. The main source of these large corrections is the mixing of the short distance operators such as

$$O_7 = \frac{e}{16\pi^2}m_Q(q_L\sigma_{\mu\nu}Q_R)F_{\mu\nu}, \quad (5)$$

generated by the interesting short distance physics, with the more mundane four-fermion operators such as

$$(q_L\gamma_\mu q_L')(q_L'\gamma^\mu Q_L), \quad (6)$$

that are generated at tree level by the SM charged currents. The mixing comes about when loop generated by gluons are taken into account. New physics, if present, will almost certainly appear in (5), not in (6), which is then a background for precision tests of the SM. Thus, the lesson from Table II is that in $b \to s\gamma$ there is still sensitivity to new physics affecting the operator (5), whereas even if it were possible to measure a branching ratio as low as $10^{-12}$ for $c \to u\gamma$, this would reflect the SM physics of operators such as the one in Eq. (6).

|                      | No QCD     | QCD Corrected |
|----------------------|------------|---------------|
| $Br(c \to u\gamma)$  | $1.5 \times 10^{-17}$ | $6.0 \times 10^{-12}$ |
| $Br(b \to s\gamma)$  | $1.3 \times 10^{-4}$     | $3.3 \times 10^{-4}$     |

Table II: Leading order and QCD-corrected branching ratios for $Q \to q\gamma$. The QCD corrected rates involve important QCD uncertainties.

The important point is that the overwhelming dominance of the QCD corrections in $c \to u\gamma$ not only tells us that the short distance physics is not sensitive to the one loop FCNC operators of interest, but also signals that there will be even larger long distance contributions to the rate. After all, the QCD corrections were computed perturbatively. In general, the dominance of the perturbative one loop amplitude by light quark contributions hints the existence of large long distance effects. Although these cannot be computed from first principles, it is possible to estimate them phenomenologically. For instance, the operator (6) with $q' = s$ gives rise to the dominant short distance piece in $c \to u\gamma$, through a $\bar{s}s$ loop. But one could imagine the $\bar{s}s$ pair propagating a long distance, forming a “$\phi$”, which turns into a photon via vector meson dominance. These and other similar long distance mechanisms [5] give rise, for instance, to $Br(D^0 \to \rho^0\gamma) \simeq 10^{-6}$, far above the
level of the QCD-corrected short distance rates expected for $c \to u\gamma$ processes. Many other long-distance dominated radiative $D$ decays are at this level. Similar effects are expected in the leptonic modes. There, however, the gap between short and long distance physics is less dramatic. The inclusive short and long distance branching fractions are \[ Br(c \to u\ell^+\ell^-)_{SD} \simeq 10^{-8}, \]
\[ Br(c \to u\ell^+\ell^-)_{LD} \simeq 10^{-6}. \]

Then, although the long distance contributions still dominate in the SM, it is still conceivable that new physics contributions could overcome them. For instance, the SM prediction for the exclusive mode \[ Br(D^0 \to \pi^0 e^+ e^-) \simeq 7 \times 10^{-7}, \]
\[ Br(B \to X_s \gamma)_{\text{CLEO}} = (3.15 \pm 0.35 \pm 0.32 \pm 0.26) \times 10^{-4}, \]
\[ Br(B \to X_s \gamma)_{\text{ALEPH}} = (3.11 \pm 0.80 \pm 0.72) \times 10^{-4}. \]

Thus, for the moment the potential long distance pollution is not problematic, but it should be taken into account in the future when precise enough measurements become available.

Similar considerations apply to the dilepton modes $b \to s\ell^+\ell^-$. In this case the long distance pollution comes in the form of “spill over” of the $J/\psi$ and $\psi'$ resonant peaks into the continuum. But in principle these modes are short distance dominated and together with $b \to s\gamma$ constitute a stringent test of the SM. In addition to the dipole moment operator in (5), these modes receive contributions from the operators

\[ O_9 = \frac{e^2}{16\pi^2}(\bar{s}_L\gamma_\mu b_L)(\bar{\ell}\gamma_\mu \ell), \]
\[ O_{10} = \frac{e^2}{16\pi^2}(\bar{s}_L\gamma_\mu b_L)(\bar{\ell}\gamma_\mu\gamma_5 \ell). \]
coefficients evaluated at the relevant experimental scale. New physics effects enter as additions to the values of the coefficients at the high energy scale $E > M_W$. The extraction of these quantities is a research program involving the inclusive rates $B \to X_s \gamma$, $B \to X_s \ell^+ \ell^-$, as well as exclusive modes such as $B \to K^{(*)} \ell^+ \ell^-$ among others. A lot of theoretical effort has gone into understanding the inclusive decay rates [11], and the theoretical predictions are under control. On the other hand, the exclusive modes are, in principle, affected by large theoretical uncertainties due to our poor knowledge of the non-perturbative dynamics determining form-factors. Some sound theoretical predictions can be made based on symmetries [12]. In some cases [13], this is enough to extract the short distance physics. In any event, in the future all these form-factors will be obtained from first principle calculations on the lattice [14], where a lot of progress has been made recently in computing weak matrix elements [15].

**SENSITIVITY TO NEW PHYSICS**

Here we discuss two typical examples of extensions of the SM of very different kind: supersymmetry and anomalous triple gauge boson couplings (TGC).

**Supersymmetry**

Supersymmetry is perhaps one of the most popular extensions of the SM. However, as many other extensions, it has a FCNC problem: in its most general form it does not come with an automatic GIM mechanism, and therefore it may generate large FCNC effects [16]. Most of the trouble comes from the fact that the diagonalization of fermion mass matrices does not, in general, diagonalize the squark mass matrices. Thus flavor mixing in the sfermion sector “misaligned” with the fermions, are a potential disaster in general SUSY scenarios. Even if the sfermion sector is assumed to be diagonal (or aligned), there is an additional source of FCNC effects, coming from the charged Higgs and chargino-squark contributions, arising from the standard CKM matrix. These effects, then are expected to be present at most at the SM level, since they depend only on the masses of the charged Higgs, the charginos and the third generation squarks. In any case, some assumption about the sfermion mass matrices is necessary in order to accommodate the FCNC constraints. The two possibilities are: (i) sfermion mass matrices are diagonal at some high energy scale (e.g. $M_{GUT}$) and small off diagonal elements are generated by the running down to the electroweak scale; (ii) they are (partially) aligned with the SM fermion mass matrices.

The vast parameter space of SUSY models includes these off diagonal elements, the superpartner and Higgs sector masses and mixings, and allows to accommodate the lack of deviations in FCNC processes such as $K^0 - \bar{K}^0$ mixing, $b \to s \gamma$, etc. However, one can argue that in most cases the SUSY effects should be “naturally” of the order of $(10 - 20)\%$ or larger. Such effects, for instance, in $K$, $D$ and $B$
mixing, or $b \rightarrow s$ and $s \rightarrow d$ transitions, are hard to see at the moment. But larger effects are also possible. For example, even satisfying the current $b \rightarrow s\gamma$ and mixing bounds, we could still see enhancements of $\mathcal{O}(1)$ in $K \rightarrow \pi\nu\bar{\nu}$ [17] and $b \rightarrow s\ell^+\ell^-$ [18].

Moreover, it has been recently argued [19] that if large off diagonal sfermion mixings are allowed, the next to leading order expansion in these mixings reveals the possibility of even larger effects in the $s \rightarrow dZ$ vertex. This would have a large impact in decay modes such as $K \rightarrow \pi\nu\bar{\nu}$, where the $Z$ penguin plays a dominant role, resulting in enhancements of the branching ratios of one order of magnitude or more, depending on the modes. This is an interesting possibility and deserves further study, particularly the correlation with possible enhancements in $D$ mixing and rare $D$ decays that would result from very large mass insertions in the up-squark sector.

**Anomalous Triple Gauge Boson Couplings**

We now turn to examine the potential of rare FCNC decays to constrain anomalous triple gauge boson couplings (TGC). In general, we can assume in a model independent way, that extensions of the SM might modify some of the couplings of fermions and/or gauge bosons. In particular, the TGC are of interest since they have not been measured with such precision as some of the fermion couplings. We have in mind deviations from the SM values for the couplings of a pair of $W$ to a photon or a $Z$. We would expect that the anomalous TGC encode the physics of some higher energy scale. Imposing $CP$ conservation, the most general form of the $WWN$ ($N = \gamma, Z$) couplings can be written as [20]
Similarly, from Fig. 1b we see that the only significant contribution of a nomalous normalized to the SM, is plotted against both ∆\(\gamma\)κ to the presence of ∆b in bκ not known. However, simplification is possible, when considering rare decays. For instance, in rare decays, we can neglect the contribution of ∆κZ, λZ, as well as the three WWZ CP violating anomalous TGC, since their effects are suppressed by powers of the small external momenta over \(m_Z\). This selective sensitivity is an advantage, rather than a handicap, when we view these measurements as complement of other ones made at higher energies and sensitive to all the Z TGC.

Thus up to this point, we have 6 coefficients left. However, we can ignore λγ and \(\tilde{\lambda}_\gamma\) if we assume that the dynamics producing these non-SM effects resides at a scale parametrically larger than the weak scale, say \(\Lambda \simeq \mathcal{O}(1)\) TeV. Although this is is not general, I believe this is a reasonable scenario, since if this was not the case we should take into account the states that are present with weak scale masses (e.g. superpartners, weakly coupled scalars, etc.) and not integrate them out as we do in an effective coupling approach. When we accept this, we see that in an effective Lagrangian approach, these coefficients can only be generated by next to leading order operators, which can be ignored since they are suppressed by \((M_W^2/\Lambda^2)\) with respect to the leading order ones³.

Thus, at the energies at hand in rare decays, the only relevant coefficients are the three CP conserving parameters (∆κγ, ∆\(g_1^Z\), \(g_5^Z\)); and a CP violating one, \(\tilde{\kappa}_\gamma\). Their FCNC effects are most interesting in B and K decays. For instance, in Fig. 1a we see the sensitivity of the current measurements of the \(b \to s\gamma\) rate to the presence of ∆κγ, whereas in Fig. 1b the branching ratio for \(K^+ \to \pi^+\nu\bar{\nu}\), normalized to the SM, is plotted against both ∆\(g_1^Z\) and \(g_5^Z\). The effect of ∆κγ in \(b \to s\gamma\) is obtained without any assumption other than the suppression of \(\lambda_\gamma\). Similarly, from Fig. 1b we see that the only significant contribution of anomalous TGC to \(K^+ \to \pi^+\nu\bar{\nu}\) is given by ∆\(g_1^Z\), which could give effects as large as factors

\[
\mathcal{L}_{WWN} = g_{WWN} \left\{ i\kappa_N W^\dagger_\mu W_\nu N^{\mu\nu} + i g_1^N \left( W^\dagger_\mu W^{\mu\nu} N^\nu - W_\mu W^{\mu\nu} N^\nu \right) \right. \\
\left. + g_5^N e^{\mu\nu\rho\sigma} (W^\dagger_\mu \partial_\nu W_\rho - W_\mu \partial_\nu W^\dagger_\rho) N_\sigma + i \frac{\lambda_N}{M_W^2} W^\dagger_\mu W^{\mu\nu} N^{\nu\lambda} \right\},
\]

with the conventional choices being \(g_{WWN} = -e\) and \(g_{WWZ} = -g \cos \theta\). Additionally, the are three CP violating Lorentz invariant terms, resulting in other 6 parameters: \(\tilde{\kappa}_N\) and \(\tilde{\lambda}_N\), obtained from (13) by replacing \(N_\mu\nu\) by the dual field strength; and \(g_4^N\) from a term similar to the second one in Eq. (13).

Gauge invariance implies \(g_1^\gamma = 1\), \(g_5^\gamma = 0\). Then, in principle, there are 11 new free parameters. Two CP conserving (∆κγ, \(\lambda_\gamma\)) and two CP violating (∆\(\tilde{\kappa}_\gamma\), \(\tilde{\lambda}_\gamma\)) affecting the WWγ couplings; four CP conserving (∆\(g_1^Z\), \(g_5^Z\), \(\lambda_Z\) and ∆κZ) and three CP violating (∆κZ, \(\tilde{\lambda}_Z\) and \(g_4^Z\)) shifting the WWZ vertex.

This is a typical problem of this type of approach, where the model independence is traded off by a large number of free parameters the sources of which are not known. However, simplification is possible, when considering rare B and K decays. For instance, the branching ratio for \(K^+ \to \pi^+\nu\bar{\nu}\), normalized to the SM, is plotted against both ∆\(g_1^Z\) and \(g_5^Z\). The effect of ∆κγ in \(b \to s\gamma\) is obtained without any assumption other than the suppression of \(\lambda_\gamma\). Similarly, from Fig. 1b we see that the only significant contribution of anomalous TGC to \(K^+ \to \pi^+\nu\bar{\nu}\) is given by ∆\(g_1^Z\), which could give effects as large as factors

³ This corresponds to the so called non-linear realization of the EWSB sector. However, it is also possible to imagine a scenario where there is a light scalar similar to the SM Higgs, with all other new states above the scale \(\Lambda\). In these linear realization scenarios, the power counting requires the consideration of \(\lambda_N\) and \(\lambda_N\) on the same footing with the other anomalous TGC.
of $(2 - 3)$ in the branching ratio. The CP violating parameter $\tilde{\kappa}_\gamma$ gives smaller effects that its CP conserving counterpart since it does not interfere with the SM. In $b \to s\ell^+\ell^-$ decays again these two coefficients ($\Delta\kappa_\gamma$, $\Delta g_1^Z$) give the dominant contributions. In Fig. 2a we see the effects of the $WWZ$ couplings on the SM normalized branching ratio, taking $\Delta\kappa_\gamma = 0$. Somewhat less dramatic effects are given by this $WW\gamma$ coupling by itself. However, an interesting feature of these decay modes, is that the additional information given by the lepton asymmetry can be used to disentangle the two contributions. As it can be seen in Fig. 2b, the effect of $\Delta g_1^Z$ on the forward-backward lepton asymmetry $A_{FB}(s)$ in $B \to K^*\ell^+\ell^-$ is such that it does not change the position of the zero [21], but affects the shape of the asymmetry as well as the rate. On the other hand, the zero of $A_{FB}(s)$ is shifted significantly by accessible values of $\Delta\kappa_\gamma$. Thus, a measurement of this asymmetry, as well as the rate, provides enough information to constrain the two relevant anomalous TGC without assuming that one of them vanishes. Then, the bounds obtained from FCNC decays are not only competitive with those from high energy colliders, but also complementary to them due to their rather selective sensitivity. For comparison, we show in Table III the projected sensitivities of LEP II [22] at
\[ \sqrt{s} = 190 \text{ GeV} \] and \[ 500 \text{pb}^{-1} \] integrated luminosity, the upgraded Tevatron [23] with \[ 1 \text{fb}^{-1} \], and a guess of the \( 3\sigma \) sensitivity to be reached in the next round of \( B \) and \( K \) experiments for FCNC decays.

### Table III. Comparison of bounds on Anomalous TGC.

| \( \Delta \kappa_{\gamma} \) | \( \Delta g^Z_1 \) | \( \tilde{\kappa}_{\gamma} \) |
|-----------------------------|-----------------|-------------------|
| (-0.25,0.40)               | (-0.08,0.08)    | -                 |
| (-0.38,0.38)               | (-0.18,0.48)    | (-0.33,0.33)      |
| (-0.20,0.20)               | (-0.10,0.10)    | (-0.50,0.50)      |

### CONCLUSIONS AND OUTLOOK

Some of the FCNC decays we discussed are largely dominated by short distance physics, a fact that makes them very sensitive to extensions of the SM entering at one loop. This is particularly true of the \( K \rightarrow \pi \nu \bar{\nu} \) modes as well as for the \( B \rightarrow X \nu \bar{\nu} \) decays. The former are accessible at experiments planned for the near future, such as KAMI and CKM, whereas is not clear how to get SM sensitivity for the neutrino modes in \( B \) decays. On the other hand, \( b \rightarrow s \gamma \) and \( b \rightarrow s \ell^+ \ell^- \) are short distance dominated modes. They may contain some long distance pollution as large as 20%, although this is theoretically very uncertain. At the moment this is not a limiting uncertainty, but it may become an issue when experiments such as LHC-B and BTeV start running.

Rare charm decays are mostly dominated by long distance dynamics. This, in general, would prevents us from using this physics to test the short distance structure of the SM, but on the other hand it constitutes a laboratory where we could improve our understanding of these effects, something we may need in rare \( B \) decays. Also, there are some exceptions in charm physics, where one may still constrain considerably new physics scenarios. We mentioned the \( c \rightarrow u \ell^+ \ell^- \) modes, and also \( D^0 - \bar{D}^0 \) mixing, have the potential to receive large non-standard effects that would appear somewhere between the current experimental limits and the most conservative estimates of the long distance effects.

Finally, we have seen how rare FCNC decays complement the searches for new physics at high energy colliders. It is possible to imagine that, with the wealth of data on these physics that will be available in the near future (BELLE, BaBar, \( K \) and \( B \) physics at the Tevatron main injector, BNL and CERN \( K \) experiments, etc.) a program similar to the electroweak precision measurements of the 1990’s could emerge. This program would then serve as guidance to the high \( p_T \) physics to be carried out at the Tevatron, the LHC and beyond.

**Acknowledgments**
This work was supported by the U.S. Department of Energy under Grant No. DE-FG02-95ER40896 and the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation.

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