Effects of the Electroweak Symmetry Breaking Sector in Rare $B$ and $K$ Decays

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Abstract

The effects of the electroweak symmetry breaking sector in rare $B$ and $K$ decays are considered in the presence of new strong dynamics at a large energy scale. In the context of the low energy effective lagrangian of the symmetry breaking sector, we focus on the contributions that are not constrained by bounds on oblique corrections or anomalous triple gauge boson vertices. We find that these have very little effect in $b \to s\gamma$, but potentially large contributions appear in $b \to s\ell^+\ell^-$ and $b \to s\nu\bar{\nu}$ as well as in $b \to sg$ processes. Deviations from the standard model in rare $K$ decays such as $K^+ \to \pi^+\nu\bar{\nu}$ and $K_L \to \pi^0\nu\bar{\nu}$ are shown to be correlated with the ones in the $B$ modes. The distinct resulting pattern of deviations from the standard model in rare decays is a characteristic feature of a strongly interacting electroweak symmetry breaking sector.

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1 Introduction

If one considers the standard model (SM) as an effective theory, its remarkable success when confronted with experiment suggests the possibility that the dynamics underlying the mechanism for electroweak symmetry breaking resides at the TeV scale or above. In particular, in a scenario without a light Higgs boson the corresponding strong dynamics might manifest itself only at very high energies as the ones to be probed by the LHC. In some cases, however, the new dynamics can produce sizeable quantum effects in low energy observables. In this letter we study the effects of a strongly coupled symmetry breaking sector in rare $B$ and kaon decays. In the context of effective field theories, we concentrate on non-decoupling effects induced by operators not contributing to oblique corrections or to shifts in the couplings of three gauge bosons. One of the remaining operators, which induces $m_t^2$ corrections in charged interactions, was recently considered in [1] where its effects in $R_b$ and $B^0_d - \bar{B}^0_d$ mixing were estimated. The logarithmic dependence on the high energy scale $\Lambda$ is used to obtain approximate bounds on the corresponding coefficient in the low energy expansion. We will show that the bounds from these observables still allow for the possibility of large effects in transitions involving flavor changing neutral currents (FCNC), most notably rare $B$ and $K$ decays.

In the absence of a light Higgs boson the symmetry breaking sector is represented by a non-renormalizable effective lagrangian corresponding to the non-linear realization of the sigma model. The essential feature is the spontaneous breaking of the global symmetry $SU(2)_L \times SU(2)_R \to SU(2)_V$. To leading order the interactions involving the Goldstone bosons associated with this mechanism and the gauge fields are described by

$$L_{LO} = -\frac{1}{4} B_{\mu\nu} B^{\mu\nu} - \frac{1}{2} \text{Tr} [W_{\mu\nu} W^{\mu\nu}] + \frac{v^2}{4} \text{Tr} [D_\mu U^\dagger D^\mu U],$$

(1)

where $B_{\mu\nu}$ and $W_{\mu\nu} = \partial_\mu W_\nu - \partial_\nu W_\mu + ig [W_\mu, W_\nu]$ are the the $U(1)_Y$ and $SU(2)_L$ field strengths respectively, the electroweak scale is $v \approx 246$ GeV and the Goldstone bosons enter through the matrices $U(x) = e^{i\pi(x)^a \tau_a/v}$. The covariant derivative acting on $U(x)$ is given by $D_\mu U(x) = \partial_\mu U(x) + ig W_\mu(x) U(x) - \frac{i}{2} g B_\mu(x) U(x) \tau_3$. To this order there are no free parameters once the gauge bosons masses are fixed. The dependence on the dynamics underlying the strong symmetry breaking sector appears at next to leading order. A complete set of operators at next to leading order includes one operator of dimension two and nineteen operators of dimension four [2,3]. The effective lagrangian to next to leading order in the basis of Ref. [3] is given by

$$L_{\text{eff.}} = L_{LO} + L'_1 + \sum_{i=1}^{19} \alpha_i L_i,$$

(2)

where $L'_1$ is a dimension two custodial symmetry violating term absent in the heavy Higgs limit of the SM. If we restrict ourselves to CP invariant structures, there remain fifteen
operators of dimension four. The coefficients of some of these operators are constrained by low energy observables. For instance precision electroweak observables constrain the coefficient of $L'_1$, which gives a contribution to $\Delta \rho$. The combinations $(\alpha_1 + \alpha_8)$ and $(\alpha_1 + \alpha_{13})$ contribute to $S$ and $U$. The coefficients $\alpha_2$, $\alpha_3$, $\alpha_9$ and $\alpha_{14}$ modify the triple gauge boson couplings and are probed at LEP II at the few percent level. The remaining operators either contribute to oblique corrections to one loop only or do not contribute. To the last group belong $L_{11}$ and $L_{12}$ given that their contributions to the gauge boson two-point functions only affect the longitudinal piece of the propagators. Of particular interest is the operator $L_{11}$ defined by:

$$L_{11} = \text{Tr} \left[ (D_\mu V^\mu)^2 \right],$$

with $V_\mu = (D_\mu U)U^\dagger$ and the covariant derivative acting on $V_\mu$ defined by $D_\mu V_\nu = \partial_\mu V_\nu + ig [W_\mu, V_\nu]$. The equations of motion for the $W_{\mu\nu}$ field strength imply:

$$D_\mu V^\mu = \frac{2i}{v^2} D_\mu J^\mu_w,$$

where the $SU(2)_L$ current is $J^\mu_w = \sum_\psi \left( \bar{\psi}_L \gamma^\mu \frac{\tau^a}{2} \psi_L \right) \tau^a$. The dominant effect appears in the quark sector due to the presence of terms proportional to $m_t$. After the quark fields are rotated to the mass eigenstate basis, the operator $L_{11}$ can be written as:

$$L_{11} = \frac{m_t^2}{v^2} \left\{ (\bar{t}_5 t)^2 - 8 \sum_{i,j} V^*_{ti} V_{tj} (\bar{q}_{iL} t_R) (\bar{t}_R q_{jL}) \right\} + \ldots$$

where $i, j = d, s, b$, the $V_{ti}$ are Cabibbo-Kobayashi-Maskawa (CKM) matrix elements and the dots stand for terms suppressed by small fermion masses. Simple inspection of this form of $L_{11}$ as a four-fermion operator makes clear where these corrections will appear. In Ref. the effects of $L_{11}$ in $R_b$ and $B^0_d - \bar{B}^0_d$ mixing were studied and found to be binding on the value of $\alpha_{11}$. There, it was concluded that the mixing limit is slightly more stringent then the one coming from $R_b$, even when the former is plagued with theoretical uncertainties. More generally, transitions involving FCNC are likely to be more sensitive to effects of this kind. In what follows we show that $L_{11}$ gives a very distinct pattern of potentially large deviations from the SM in rare $B$ and $K$ decays.

Operators of this type are also present in scenarios with a light Higgs boson. In this case the SM symmetry breaking sector is linearly realized. Using naive dimensional analysis we conclude that the analogous operators generating non-oblique corrections and expressible as four-fermion interactions are of dimension eight or higher and therefore are suppressed by additional powers of $v^2/\Lambda^2$. Thus the effects we discuss in this paper are of no significance in the presence of a light Higgs boson.

In the next section we compute the effects induced in rare $B$ and $K$ processes by the four-fermion operators contained in $L_{11}$. Additional contributions to these decays from
the next-to-leading order effective lagrangian in (2) could come in the form of anomalous triple gauge boson couplings. We do not consider these corrections here. Moreover, we assume the coupling of the new physics to fermions to be small and therefore neglect contributions from induced anomalous couplings of fermions to gauge bosons [7], which in any case are constrained by the bounds on oblique corrections. We summarize and discuss the results in the last section.

2 Effects in Rare $B$ and $K$ Decays

The four-fermion operator in (5)

$$\mathcal{L}_{11} = \frac{-8m_t^2}{v^4} \left\{ V_{ts}^* V_{tb} \bar{s}_L t_R \bar{t}_R b_L + V_{td}^* V_{tb} \bar{d}_L t_R \bar{t}_R b_L + V_{ts}^* V_{td} \bar{d}_L t_R \bar{t}_R s_L \right\} + \ldots$$

induces contributions to several FCNC processes when inserted in a top loop, as shown in Fig. 1. Among them are $b \to q\gamma$, $b \to qZ$ and $b \to qg$ transitions ($q = d, s$) as well as the analogous vertices relevant for rare $K$ decays. The insertion of this non-renormalizable interaction results in a logarithmically divergent amplitude when computed using dimensional regularization. The divergences are absorbed by suitable counterterms. We will estimate the size of the effects by computing the logarithmic dependence on the high energy scale $\Lambda$. In principle, there could be finite effects coming from the counterterms. However, unless large cancellations occur, the logarithmic dependence should give a good approximation of the effects [6,8]. We will come back to discuss this point later in the paper.

We begin with the modification of the $b \to s\gamma$ vertex. This can be generically written
The second term in $(7)$ gives the on-shell amplitude whereas the term proportional to $A$ only contributes to the vertex when the photon is off-shell. The computation of the corresponding diagram in Fig. 1 (a) gives

$$A = \frac{8Q_t m_t^2}{3} \frac{m^4_t}{v^4} \alpha_{11} \log \left( \frac{\Lambda^2}{m_t^2} \right),$$

$$B = 0,$$

where $Q_t$ denotes the top quark electric charge. As it was mentioned before, the value of $A$ in $(8)$ corresponds to the logarithmic $\Lambda$ dependence and does not include possible finite counterterm contributions. However it gives the correct dependence on the high energy scale $\Lambda$ $(8)$. The leading corrections to $b \to s\ell^+\ell^-$ processes are therefore governed by $A$. From $(8)$ we conclude that the operator $\mathcal{L}_{11}$ does not induce an effect in $b \to s\gamma$ processes at the one loop level. However, next to leading order corrections induced by the strong interactions will give a non-zero value of $B$. The need to go to two loops is characteristic of the mixing of four-quark operators with the electromagnetic dipole operator in $(7)$. A typical contribution of this type is shown in Fig. 1 (b). We compute this mixing after Fierz rearranging $(6)$ into a product of currents and evolving the corresponding coefficient from the scale $\Lambda$ down to $M_W$, where the matching to the effective weak hamiltonian is made. The dependence with the high energy scale now enters through the factor $\frac{\alpha_s}{4\pi} \log \Lambda^2/M_W^2$, to be compared with $(8)$ where there is no $\alpha_s$ suppression. The effect of the resummation of these logarithms through the renormalization group equations is small. The renormalization group running from $M_W$ down to $\mu \approx m_b$ produces the mixing of the four-quark operator with the second term in $(7)$. This gives, approximately

$$B \approx -\frac{1}{3} \left[ \sum_{i=1}^{8} h_i \eta^{a_i} \right] \alpha_{11} \left( \frac{m_t}{v} \right)^2,$$

where we have neglected a $\approx 10\%$ effect from the running between $\Lambda$ and $M_W$, we define $\eta \equiv \alpha_s(M_W)/\alpha_s(m_b)$ and the coefficients $h_i$ and $a_i$ can be found in $(12)$. As we will discuss below, this two-loop contribution results in a small effect in $b \to s\gamma$ processes.

Next we consider the $b \to sZ$ vertex. The operator $(8)$ adds the contribution

$$\delta\Gamma_{\mu}^{b\to sZ} = (\frac{g_W}{c\theta_W}) V_{tb}^{*} V_{ts} \frac{1}{4\pi^2} \frac{m_t^4}{v^4} \alpha_{11} \log \left( \frac{\Lambda^2}{m_t^2} \right) \bar{s}_L \gamma_\mu b_L,$$

with $g_W$ the $SU(2)_L$ gauge coupling and $c\theta_W$ the cosine of the Weinberg angle. The additional factor of $m_t^2$ in $(11)$ when compared to $q^2$ in $(7)$ and $(8)$ comes from the axial-vector coupling of the $Z$ boson, which here gives the leading contribution.
At this point and before estimating the effects of $\mathcal{L}_{11}$ in rare decays, we incorporate the constraints on the quantity $y \equiv \alpha_{11} \log \frac{A^{2}}{m^{2}}$, characterizing the new physics contributions. Implicit in this procedure is the assumption that the contributions of finite counterterms do not change significantly the size of the contribution. Furthermore, there are extensions of the SM where there are no finite counterterms at the scale $\Lambda$. Neglecting the counterterm contributions allows one to correlate the effects of the new dynamics in various observables, given that these are proportional to $y$. The upper limit for $y$ that is obtained from the analysis of $K^0 - \bar{K}^0$ and $B_d^0 - \bar{B}_d^0$ mixing in $[4]$ corresponds to $y < 0.4$. A more conservative estimate using a higher value for $B_K = 0.80 \pm 0.20$ $[4]$ and assuming a 50% uncertainty in the value of $|V_{ub}/V_{cb}|$ yields $y < 0.50$. A lower limit on $y$ can be derived by comparing the experimental value $[11]$ $R_b = 0.2178 \pm 0.0011$ with the SM expectation of $R_{b,SM} = 0.2158$. Noting that $[4]$

$$\frac{\delta \Gamma_b}{\Gamma_b^{SM}} = 4.58 \times \frac{1}{4\pi^2} \left(\frac{m_t}{v}\right)^4 y,$$

(12)

the 95% C.L. lower limit is $y > -0.40$. This limit is largely driven by the fact that $R_b$ is larger that the SM prediction by about $2\sigma$. For instance, had the central value of $R_b$ agreed with the SM, the 95% C.L. lower limit would be $y > -0.80$.

The induced vertices of eqns. (9) and (11) translate into deviations from the SM predictions for $B$ decays governed by the $b \rightarrow s \gamma$ and $b \rightarrow s \ell^+\ell^-$ transitions. They can be expressed as shifts in the Wilson coefficients in the weak effective hamiltonian

$$H_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} \sum_{i=1}^{10} C_i(\mu)O_i(\mu),$$

(13)

with the operator basis defined in Ref. [9]. From eqn. (10) we see that there is a small shift in the coefficient $C_7(M_W)$ of the electromagnetic dipole operator, $O_7 = \frac{e}{16\pi^2} m_b \bar{s}_L \gamma_\mu b_R F^{\mu\nu}$. The renormalization group induced mixing of $\mathcal{L}_{11}$ with $O_7$ translates into a shift for this coefficient of the order of $\approx -0.05\alpha_{11}$. For instance, for a typical value of $\Lambda = 2$ TeV the bounds on $y$ discussed above give $|\alpha_{11}| < 0.1$ which gives a shift of at most 2% in the coefficient $C_7(m_b)$ determining the $b \rightarrow s \gamma$ amplitude.

The coefficient $C_9(M_W)$ of the operator $O_9 = \frac{e}{16\pi^2} (\bar{s}_L \gamma_\mu b_L)(\bar{\ell}\gamma^\mu \ell)$ is affected almost exclusively by the shift in the off-shell photon vertex (8) given that the contribution to it from the $Z$ exchange is proportional to $-1/4 + s^2\theta_W \approx 0$. On the other hand, the induced vertex in (11) produces a significant shift in $C_{10}(M_W)$, the Wilson coefficient corresponding to the operator $O_{10} = \frac{e}{16\pi^2} (\bar{s}_L \gamma_\mu b_L)(\bar{\ell}\gamma^\mu \gamma_5 \ell)$. We make use of next-to-leading order QCD results for $C_9(m_b)$ [12]. On the other hand, the coefficient $C_{10}(m_b) = C_{10}(M_W)$ and can be found in [4]. The entire analysis above is also valid for $b \rightarrow d \ell^+\ell^-$ with the replacements $s \rightarrow d$ and $V_{ts} \rightarrow V_{td}$, given that we have neglected $(m_s/m_b)$ effects. In Fig. 2 (solid line)
Figure 2: Ratio of the modified branching ratio to the standard model expectation as a function of $y = \alpha_{11} \log \frac{\Lambda^2}{m_t^2}$. The solid line corresponds to the ratio $R_\ell$ for $B \to X_{(s,d)} \ell^+ \ell^-$ inclusive decays, the dashed line to $R_\nu$ for $B \to X_{(s,d)} \nu \bar{\nu}$ and the dot-dashed line to $R_g$ for $b \to s \bar{s}s$ decays.

we estimate the size of the effect in $b \to (s,d) \ell^+ \ell^-$ decays by plotting the ratio

$$R_\ell = \frac{Br(B \to X_{(s,d)} \ell^+ \ell^-)}{Br(B \to X_{(s,d)} \ell^+ \ell^-)_{SM}} \tag{14}$$

as a function of the quantity $y$. We consider a range of values of $y$ compatible with the limits from $R_b$ and $B^0_d - \bar{B}^0_d$ mixing as discussed above. We observe that for positive (negative) values of $\alpha_{11}$ the branching ratio can be enhanced (suppressed) by up to a factor of about two. For instance, the SM prediction for the $B \to X_{s} \mu^+ \mu^-$ branching fraction is about $(5 - 6) \times 10^{-6}$. Thus, the maximum enhancement in $R_\ell$ would take this mode to $\approx 1 \times 10^{-5}$. However, it should be noted that the effect does not necessarily affect all exclusive branching ratios equally. For instance, the $B \to K \ell^+ \ell^-$ branching fraction follows $R_\ell$, whereas this is not the case for $B \to K^* \ell^+ \ell^-$, where the interplay of the various combinations of Wilson coefficients with the helicity amplitudes gives a quantitatively different answer (e.g. for $y = 0.50$ the enhancement is $\approx 50\%$, for $y = -0.50$ the suppression is small). On the other hand, in these modes the forward-backward asymmetry for leptons is very sensitive to these type of changes in the Wilson coefficients [13]. Current experimental upper limits on exclusive modes are already at the $1 \times 10^{-5}$ level [14] and sensitivity to SM branching ratios will be achieved in the near future.

The modification of the $b \to sZ$ vertex [11] also induces an effect in $b \to s \nu \bar{\nu}$ processes.
Defining the ratio $R_\nu$ analogously to (14) and plotting it versus $y$ in Fig. 2 (dashed line) we see that the effect of $L_{11}$ in these decays approximately follows $R_\ell$. The SM expectation is $Br(B \to X_s \nu\bar{\nu}) = (4.50 \times 10^{-5})$. The current 90\% C.L. upper limit from LEP \cite{15} is $7.7 \times 10^{-4}$, which is still not constraining when compared with $R_\nu$.

The operator $L_{11}$ also induces corrections to the Wilson coefficients of QCD penguin operators in the effective weak hamiltonian. For instance the modification of the $b \to s g$ vertex due to (6) is

$$
\delta \Gamma_{b \to s g} = (-i) g \left( \frac{m_t^2}{v^4} \right) \frac{y}{3\pi^2} s_L \left( \not{q}_\mu - q^2 \gamma_\mu \right) b_L
$$

(15)

leading to shifts in the coefficients of the QCD penguin operators $C_3$, $C_4$, $C_5$ and $C_6$ defined in Ref. \cite{9}, but not affecting the gluonic dipole operator. In order to estimate the size of the effect we compute the branching fraction for the pure QCD penguin process $b \to s\bar{s}s$. The ratio $R_g$, defined analogously to (14) is plotted in Fig. 2 (dot-dashed line). The deviation with respect to the SM is slightly reduced in this case when compared to the semileptonic cases. Furthermore, the theoretical uncertainties associated with the observable exclusive modes are rather large and might obscure any new physics effects \cite{10}. In any case, we see the same correlation with the sign of $y$ as in $R_\ell$ and $R_\nu$.

Finally, we turn to study the effects of the operator $L_{11}$ in the kaon system. These are induced by the last term in (3). We focus on the semileptonic decays affected by the modification of the $s \to dZ$ vertex, which allows us to concentrate on the theoretically cleaner modes. For the CP violating mode $K_L \to \pi^0\nu\bar{\nu}$ the ratio to the SM branching fraction, $R_{K_L}$, is identical to $R_\nu$ and is therefore given by the dashed line in Fig. 2. On the other hand, the equivalent ratio for the $K^+ \to \pi^+\nu\bar{\nu}$ decay depends mildly on the CP violation parameters $\eta$ and $\rho$ as they appear in the Wolfenstein parametrization \cite{17} of the CKM matrix, as well as on $V_{cb}$. The ratio $R_{K^+}$ is shown in Fig. 3 (solid line) as a function of $y$. It can be seen that the pattern of deviation from the SM is very similar to that of the neutral mode as well as to the one observed in rare $B$ decays. The current experimental limit on the charged mode is $Br(K^+ \to \pi^+\nu\bar{\nu}) < 2.40 \times 10^{-9}$ \cite{18} whereas the SM predictions are in the vicinity of $1 \times 10^{-10}$. Thus it is likely that experiments in the near future will have sensitivity at the SM level. Similar effects will be present in other decay modes, for instance $K_L \to \mu^+\mu^-$, $K^+ \to \pi^+\mu^+\mu^-$, etc. However, these decays are contaminated by potentially large long distance contributions \cite{19}.

3 Summary and Discussion

In a scenario without a light Higgs boson the electroweak sector is non-linearly realized and most likely underlied by strong dynamics. We have studied the effects of such a
scenario in rare $B$ and $K$ decays. These originate almost exclusively in one operator in the next-to-leading order effective lagrangian for the Goldstone bosons, given that its effects are proportional to $m_t^2$. The dependence on the high energy scale $\Lambda$ resulting from the insertion of the operator $\mathcal{L}_{11}$ given in (6) in FCNC loops is given by the factor $y = \alpha_{11} \log \frac{\Lambda^2}{m_t^2}$. We further assume this dependence to be the leading contribution to the amplitudes. In doing so we neglect possible finite contributions from counterterms coming from the matching at the scale $\Lambda$. This allows us to correlate the effects in several observables. We have shown that, even after imposing the bounds on $y$ from $R_b$ and $B^0_d - \bar{B}^0_d$ mixing, large effects on rare $B$ and $K$ decays remain, as can be seen in Figs. 2 and 3. A very distinct pattern of deviations from the SM emerges. Positive values of $y$ give similar enhancements of up to a factor of almost two in $B \to X_{(s,d)} \ell^+\ell^-$, $B \to X_{(s,d)} \nu\bar{\nu}$, $b \to s\bar{s}s$ and other QCD gluonic penguin decays, $K^+ \to \pi^+\nu\bar{\nu}$ and $K_L \to \pi^0\nu\bar{\nu}$, whereas negative values of $\alpha_{11}$ imply a similar suppression of up to a factor of two in all these processes. Other modes, such as $B_{(s,d)} \to \ell^+\ell^-$, $K_L \to \ell^+\ell^-$ and $K^+ \to \pi^+\ell^+\ell^-$, are similarly affected but we concentrated on those processes that are both theoretically clean and most accessible to present and future experiments. The effects of the operator $\mathcal{L}_{11}$ in on-shell photon processes as $b \to s\gamma$, $s \to d\gamma$, etc. are only induced by the two-loop strong interaction mixing and are expected to be rather small when compared with those in the leptonic and semileptonic modes for the same values of the parameter $y$.

Both the correlation among different processes as well as the size of the effects are well approximated by the logarithmic behavior as long as there are no large cancellations due to counterterms. Moreover, in extensions of the SM where the new dynamics associated with the electroweak symmetry breaking does not couple appreciably to fermions, there will be no counterterms and the observed correlation among processes is not just an expected pattern but a quantitatively valid statement stemming from the calculation of the effects. This is consistent with our assumption of no anomalous couplings of fermions to gauge bosons.

Regarding the size of the coefficient $\alpha_{11}$ we note that, for a high energy scale of e.g. $\Lambda = 2$ TeV, its value is allowed to be $|\alpha_{11}| < 0.1$. We have not made an attempt to estimate the size of this coefficient in extensions of the SM. In the scenario considered in this work, the new dynamics is strongly coupled and the coefficients of the effective lagrangian can be large.

Finally, we point out that this is one more example of deviations occurring at the same level in both rare $B$ and $K$ decays. This is characteristic of new physics entering in loop-induced FCNC as long as the new dynamics couples equally to all generations. In this way, $m_t^2$ effects associated with the electroweak symmetry breaking affect similarly FCNC with external down-type quarks. Therefore the experimental information from both types of processes is essential in disentangling the source of the effect in scenarios.
Figure 3: Ratio of the modified branching ratio to the standard model expectation for $K^+ \to \pi^+ \nu \bar{\nu}$ (solid line) and $K_L \to \pi^0 \nu \bar{\nu}$ (dashed line).

like the present one. Experiments designed to achieve sensitivity to the SM branching fractions, such as leptonic and hadronic $B$ factories and high intensity kaon experiments can have interesting consequences for the physics at the high energy scale $\Lambda$ even when this can only be directly probed at the LHC.

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References

[1] J. Bernabéu, D. Comelli, A. Pich and A. Santamaria, Phys. Rev. Lett. 78, 2902 (1997).

[2] A. Longhitano, Phys. Rev. D22, 1166 (1980), Nucl. Phys. B188, 118 (1981).

[3] T. Appelquist and G. Wu, Phys. Rev. D48, 3235 (1993).

[4] K. Hagiwara, K. Hikasa, R. D. Peccei and D. Zeppenfeld, Nucl. Phys. B282, 253 (1987); K. Hagiwara, S. Ishiara, R. Szalapski and D. Zeppenfeld, Phys. Lett. B283, 353 (1992) and Phys. Rev. D48, 2182 (1993).

[5] F. Feruglio, Int. J. Mod. Phys. A8, 4937 (1993).

[6] A. Manohar and H. Georgi, Nucl. Phys. B234, 189 (1984); H. Georgi, “Weak Interactions and Modern Particle Theory”, Benjamin/Cummings, Menlo Park, California, 1984.

[7] R. D. Peccei and X. Zhang, Nucl. Phys. B337, 269 (1990); R. D. Peccei, S. Peris and X. Zhang, Nucl. Phys. B349, 305 (1991).

[8] C. P. Burgess and D. London, Phys. Rev. Lett. 69, 3428 (1992); Phys. Rev. D48, 4227 (1993).

[9] B. Grinstein, R. Springer and M. B. Wise, Nucl. Phys. B339, 269 (1990); M. Ciuchini et al., Phys. Lett. B316, 127 (1993); M. Ciuchini, E. Franco, L. Reina and L. Silvestrini, Nucl. Phys. B421, 41 (1994); A. Buras, M. Misiak, M. Münz and S. Pokorski, Nucl. Phys. B424, 374 (1994). For a review see G. Buchalla, A. J. Buras and M. E. Lautenbacher, Rev. Mod. Phys. 68, 1125 (1996).

[10] P. B. Mackenzie, proceedings of the 8th meeting of the Division of Particles and Fields, Albuquerque, New Mexico, August 2-6 1994; A. Soni, hep-lat/9510036.

[11] J. Alcaraz et al., Combined LEP and SLD Electroweak Working Groups, preprint LEPEWWG/96-02, SLD Physics Note 52.

[12] A. J. Buras and M. Münz, Phys. Rev. D52, 186 (1995).

[13] G. Burdman, Phys. Rev. D52, 6400 (1995); D. Liu, Phys. Lett. B346, 355 (1995).

[14] C. Albajar et al., the UA1 collaboration, Phys. Lett. B262, 163 (1991); F. Abe et al., the CDF collaboration, Phys. Rev. Lett. 76, 4675 (1996); R. Balest et al., the CLEO collaboration, preprint CLEO-CONF-94-4, contributed paper to the International Conference on High Energy Physics, Glasgow, Scotland, 1994, unpublished.
[15] The ALEPH collaboration, preprint PA10-019, paper contributed to the International Conference on High Energy Physics, Warsaw, Poland, 25-31 July 1996.

[16] N. G. Deshpande and X. He, *Phys. Lett.* **B336**, 471 (1994).

[17] L. Wolfenstein, *Phys. Rev. Lett.* **51**, 1945 (1983).

[18] S. Adler *et al.*, the BNL 787 collaboration, *Phys. Rev. Lett.* **76**, 1421 (1996).

[19] L. Littenberg and G. Valencia, *Annu. Rev. Nucl. Part. Sci.* **43**, 729 (1993).