Rare B decays in left-right-symmetric models

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The manifestation of the higher-order effects induced by the flavor-changing neutral currents in B decays is analyzed in the left-right-symmetric generalization of the electroweak unification. In such left-right models the effects of the right-handed currents and the mixing between the two gauge bosons give rise to some enhancement relative to the standard model in the branching ratio of the rare processes $b \to s\gamma$, $b \to sg$, $b \to sl^+l^-$, and $b \to s\nu\bar{\nu}$, in particular if the present bounds on the W_R mass and on the mixing ξ turn out to be too severe. We can conclude that this kind of rare B physics can put interesting constraints on these models.

I. INTRODUCTION

The studies of the weak decays of the heavy flavors,[1,2,3] in particular the rare modes, are a powerful way to probe the fundamental interactions, since the rare decay processes appear quite sensitive to the eventual presence of a new interaction beyond the standard electroweak model. In particular, the rare B decays are important when one looks for flavor-changing neutral currents and can yield evidence for nonstandard processes: a fourth generation, new scalars, new gauge bosons, or supersymmetric partners of the usual particles. Comparing to the s quark, loop decays of the b quark are especially promising as (1) they are generally not as rare, (2) they are more readily amenable to perturbative QCD analysis, and (3) the b quark, being a member of the third family, is likely to be more sensitive to the presence of the fourth family. The loop effects in the B meson involve the element $V_{ts}^* V_{tb}$, which is much larger than the kaon counterpart. So rare-B-decay branching ratios can be enhanced relative to the analogous K decay modes by a factor $|V_{ts}^* V_{tb}/V_{cb}|^2/|V_{td}^* V_{ts}/V_{ub}|^2 \sim 10^5$.

Moreover, the one-loop processes are enhanced by the presence of a heavy top quark and recent evidence suggests that the top-quark mass may be larger than half of the Z^0 mass. The UA1 Collaboration has set the limit $m_t > 44$ GeV using a conservative calculation for $\sigma(p\bar{p} \to t\bar{t}X)$, and $m_t > 56$ GeV using a reasonable calculation for this cross section.[4] The ARGUS Collaboration[5] has reported the observation at the Υ''' of B_d-$\bar{B}_d$ mixing with a value much greater than the prediction of the standard model.[6] This result is still compatible provided that the mass of the t quark is much greater than 50 GeV (Ref. 7).

The gradually evolving theoretical techniques lead to successive treatments of these decays: from symmetry concepts[8] and the pole models,[9] as well as sophisticated models on hadronic levels,[10] to the modern approach of the effective Hamiltonians of quarks and gluons, which is based on the Weinberg-Salam theory of the electroweak interactions to which QCD corrections are applied.[11] The behavior of the weak amplitudes for different rare-B-decay modes in the presence of a virtual top quark with mass m_t satisfying $m_t/M_W > 1$ appears phenomenologically exciting.

A study of the rare B decays in the context of a left-right-symmetric (LRS) model of the electroweak theory, including an analysis of the effects of a heavy top quark in the flavor-changing neutral currents (FCNC's) at the one-loop level, constitutes the main topic of this paper.

First we examine the features of the minimal LRS models and put in evidence the new peculiar parameters of the theory, i.e., the mass of the new charged gauge boson and its mixing with the standard one. We consider the Higgs-boson masses, except one, at the TeV scale and similarly with suppressed effects in the rare B decay.

In the successive part of the paper we deal with the calculations of the transition form factors of the typical relevant B-decay modes in the LRS models. In such models, the effects of the right-handed current and of the mixing give rise to some enhancement in the branching ratios of the different rare processes $b\to s\gamma$, $b\to sl^+l^-$, $b\to s\nu\bar{\nu}$, and $b\to sg$, even if such branching ratios are in general not really competitive with those found in other scenarios such as in two-doublet models or in supersymmetry.[12]

A detailed study of the different processes is performed by looking at the general contributions of the one-loop corrections for the decays of the heavy mesons in the range $40<m_t<250$ GeV. In particular, the calculation of the semileptonic $b\to sl^+l^-$ decay mode includes the effects of the photon, the Z vertex, and the box-diagram contributions.

Further, we draw a conclusion for the case of a (virtual) gluon exchange which would seem to warrant further study. Finally, in Sec IV we briefly summarize our results.

II. THE MINIMAL LEFT-RIGHT-SYMMETRIC MODEL

In the standard model (SM) the gauge symmetry is broken spontaneously, while discrete symmetries such as parity and charge conjugation are broken explicitly. It is certainly theoretically appealing to consider extensions of the SM where parity is conserved in the Lagrangian and is broken spontaneously together with the gauge symmetry.

A model involving both $V-A$ and $V+A$ currents was suggested before the advent of gauge theories.[13] Since then, the left-right-symmetric (LRS) theories have been discussed extensively in the literature.[14] The minimal electroweak gauge model based on the group

$$\mathcal{G}^{w}_{LR}=\mathrm{SU}(2)_L\otimes\mathrm{SU}(2)_R\otimes\mathrm{U}(1) \tag{1}$$

involves three fermion generations, and a single quadruplet of scalars

$$\Phi=\begin{bmatrix}\phi_1^{(0)} & \phi_1^{(+)}\\ \phi_2^{(-)} & \phi_2^{(0)}\end{bmatrix}, \tag{2}$$

coupled to the quarks together with the charge-conjugate field $\tilde{\Phi}=\tau_2\Phi^*\tau_2$. With only this Higgs content, the symmetry would be broken down to $\mathrm{U}(1)\otimes\mathrm{U}(1)$. To break an extra unwanted U(1) symmetry, the Higgs sector must be enlarged with the introduction either of the doublets $(\frac{1}{2},0)$, $(0,\frac{1}{2})$,

$$\chi_{L,R}=\begin{bmatrix}\chi^{(+)}\\ \chi^{(0)}\end{bmatrix}_{L,R}, \tag{3}$$

or, alternatively, of the representations $(\frac{1}{2},\frac{1}{2})$,

$$\Delta_{L,R}=\begin{bmatrix}\delta^{(+)}/\sqrt{2} & \delta^{(++)}\\ \delta^{(0)} & -\delta^{(+)}/\sqrt{2}\end{bmatrix}_{L,R}, \tag{4}$$

which make parity spontaneously broken at a mass higher than that of the standard charged gauge boson. The breaking of the parity and the predominantly $V-A$ feature of the charged currents at low energy can be achieved spontaneously by imposing an asymmetry structure of the vacuum

$$\begin{aligned}v_R=\langle\chi_R^0\rangle \text{ or } \langle\delta_R^0\rangle &\gg\langle\phi_1^{(0)}\rangle\\ &=k\gg\langle\phi_2^{(0)}\rangle=k'\sim\langle\delta_L^0\rangle\simeq 0=\langle\chi_L^0\rangle\ .\end{aligned} \tag{5}$$

This hierarchical feature emerges as a wider solution from the study of the most general potential and upholds the motivation of connecting parity-violation processes to the suppression of the right-handed weak interactions.[15]

The Lagrangian which describes the interaction of the gauge bosons with quarks

$$\begin{aligned}\mathcal{L}_{fW}&=i\frac{g}{\sqrt{2}}(\bar{Q}^{(w)}_{uL}\gamma^\mu Q^{(w)}_{dL}W^{(+)}_{L\mu}+\bar{Q}^{(w)}_{uR}\gamma^\mu Q^{(w)}_{dR}W^{(+)}_{R\mu})+\text{H.c.}\\ &=i\frac{g}{\sqrt{2}}(\bar{Q}^{(s)}_{uL}\gamma^\mu U_{cL}Q^{(s)}_{dL}W^{(+)}_{L\mu}\\ &\qquad+\bar{Q}^{(s)}_{uR}\gamma^\mu U_{cR}Q^{(s)}_{dR}W^{(+)}_{R\mu})+\text{H.c.}\ ,\end{aligned} \tag{6}$$

where the renormalized coupling constants g_L and g_R can be taken to be equal, and a single g can be used, provided that quartic couplings such as $\mathrm{tr}(\Phi\Phi^\dagger)(\chi_L^\dagger\chi_L+\chi_R^\dagger\chi_R)$ are chosen in the LRS form.

The most important undetermined parameters of the LRS models are the mass scale $M_R\equiv M_{W_R}$ of the left-handed gauge bosons, the mixing ξ between $W_L^\pm$ and $W_R^\pm$, and the Cabibbo-Kobayashi-Maskawa- (CKM-) type angles[16] for the right-handed sector of the theory. We choose to work within the framework of a specific version of the theory in which all right-handed CKM-type angles and phases are equal to the corresponding left-handed parameters ("manifest" left-right symmetry). We are led to this choice mainly by the criterion of simplicity; other variants of the theory, for instance, the interesting one which requires charge conjugation to be also spontaneously broken,[15] may lead to results that, in some cases, are rather different from the ones we obtain.

Let us indicate the mass eigenstates of the charged gauge bosons as W_1 and W_2. W_L and W_R are in general given by orthogonal linear combinations

$$W_L=\cos\xi\ W_1+\sin\xi\ W_2,\quad W_R=-\sin\xi\ W_1+\cos\xi\ W_2\ , \tag{7}$$

the L-R mixing parameter ξ being related to the vacuum expectation values (VEV's) of the Higgs bosons according to

$$\xi=\begin{cases}\frac{1}{2}\arctan\dfrac{4kk'}{v_R} & \text{in the } \chi \text{ case}\ ,\\[2ex] \sim\dfrac{kk'}{k^2+k'^2+8v_R}\simeq\dfrac{kk'}{k^2+k'^2}\dfrac{M_1}{M_2} & \text{in the } \Delta \text{ case}\ .\end{cases} \tag{8}$$

It is worth mentioning that if one imposes a symmetry which assures massless Dirac neutrinos and the absence of the W_L^--W_R^+ mixing at the tree level, this scenario cannot remain unaltered in higher-order perturbation theory.[17]

On the other hand, the dominant $V-A$ structure of μ and β decays suggests a calculable and naturally small mixing angle: a lower bound on M_R and an upper bound on ξ can be derived from the e^+ spectrum measured in polarized μ^+ decay.[18] This result, however, does not appear so restrictive if compared to later estimates. In fact, current-algebra analysis of strangeness changing purely nonleptonic decays of hadrons severely restricts the W_L-W_R mixing: $\xi<0.004$ (Ref. 19). If one adopts the Adler-Weissberger relation in the case of the manifest version of LRS one obtains $\xi<0.0055$ (Ref. 20) [by including the 5% of the fractional error in V_{us} due to the SU(3) breaking, the radiative corrections in semileptonic hyperon and K_{l3} decays and the bounds on the b-decay branching ratio]. We shall perform our study adopting this latter much stronger bound. However, in view of the theoretical uncertainties which are present in its derivation, we shall mention also the results that one obtains taking the milder limit on ξ.

The most severe bound on M_R comes from the requirement that the W_L-W_R box diagram yield a contribution to the K_S-K_L mass difference smaller than that of the standard Gaillard-Lee box diagram involving two W_L bosons.[21] This gives $M_R\geq 1.6$ TeV. Apart from that of manifest LR symmetry (which here plays a crucial role), some other assumptions are involved in deriving this bound. In particular one must bar accidental and substantial cancellations between the W_L-W_R contribution and the t-quark contribution (the above limit was obtained considering only two generations of quarks). We shall impose the above strict bound throughout all our analysis.

As we already stressed, the effect of the Higgs scalars is not competitive with that of the gauge bosons. If we consider the Higgs-field multiplets $\chi_{L,R}$ and Φ, the model has 16 real scalar degrees of freedom (20 in the case of the $\Delta_{L,R},\Phi$) eight natural real fields and eight charged ones. Six of them give mass to the six vector gauge bosons: four charged ($W^{\pm}_{L,R}$) and two neutral ($Z^0_{1,2}$). The remaining fields are four charged and six neutral massive physical Higgs scalars. In the $\Delta_{L,R},\Phi$ case, of the 14 physical scalar fields four are doubly charged, four are singly charged, and six are neutral. All of them are expected to be heavy except a neutral combination.

It has been pointed out[22] that Higgs-boson-induced FCNC's play a crucial role in the K_S-K_L problem even at the tree level and that their contributions to the K_S-K_L mass difference are dangerously large unless the masses of the relevant Higgs scalars (which are expected to be of the same order as M_R) are very large. In fact, the exchange of the following physical Higgs boson

$$\frac{k\phi_2^{0*}-k^*\phi_1^0}{\sqrt{k^2+k'^2}} \tag{9}$$

produces FCNC's at the tree level that are not competitive with the one-loop processes only if the mass of the neutral scalars is of the order of that of the right-handed gauge bosons.[23] Also, the physical charged Higgs bosons are heavy, and, hence, their contributions to the one-loop diagrams leading to FCNC processes are expected to be correspondingly suppressed. That is why we shall not entertain the possibility of FCNC enhancements due to charged-Higgs-boson exchange in LRS models, and in estimating the contribution to rare B decays in LRS models we shall always consider only the contributions that are obtained by exchanging the W_L and W_R gauge bosons.

Let us conclude with a comment on the semileptonic channels. Some rare B-decay processes such as $B\to Kl^+l^-$ are obscured by the not-well-known neutrino sector. In the LRS models, the charged-current interaction of the leptons can be written as

$$\mathcal{L}_{\rm CC}=i\frac{g}{\sqrt{2}}(\bar{l}_L\gamma^\mu\nu_{lL}W^{(+)}_{L\mu}+\bar{l}_R\gamma^\mu N_{lR}W^{(+)}_{R\mu})+\text{H.c.} , \tag{10}$$

where $l=e,\mu,\tau$. Without loss of generality, we can choose a basis in which the mass matrix of the charged leptons l is diagonal with eigenvalues m_l. In this same basis, the left-handed neutrinos are made light, at least in the most frequently studied versions of left-right models, by giving large mass to the right-handed neutrinos.[24]

In the $\chi_{L,R}$ case, Dirac neutrino and charged-lepton masses can be written down in the same way as quark masses. Anyway, in general the involved leptonic matrices are arbitrary and not related to the quark mass matrices: they are characterized by independent mixing angles. Hence, it does not appear easy to perform a model-independent analysis. A hint can be suggested by the embedding of $SU(2)_L\otimes SU(2)_R\otimes U(1)$ in SO(10) (Ref. 25). However, since our aim is to realize a model-independent analysis, we may phenomenologically consider two classes of neutrino masses, one with light mass eigenstates and the other class with higher neutrino masses.

III. THE RARE B-DECAY MODES

In the calculations of the one-loop graphs in the LRS gauge theories the dominant contributions, as we stressed in the previous section, come from the exchange of the gauge bosons. However, in order to compute the transition form factors in a general ξ gauge,[26] extra diagrams appear in the game. The unphysical bosons G_i (absorbed by W_i and Z_i in the unitary gauge) give small contributions and can be included in the calculation without any detailed knowledge of the scalar content of the theory, by invoking tree-level unitarity and gauge independence in the calculations.[27] Only the leptonic sector of the theory was examined at the one-loop level,[28] neglecting, however, the W_L-W_R mixing effects. In our analysis we only neglect, in general, the small charged physical and unphysical scalar contributions since the scale of their masses is of the order of M_R, but we retain a small nonvanishing mixing.

For further applications let us define in this section the branching ratio of a generic process $b\to s+X$ by means of the usual relation

$$B(b\to s+X)=\frac{\Gamma(b\to s+X)}{\Gamma(b\to(u,c)l\nu)}B(B\to l\bar{\nu}+X)$$
$$=\frac{\Gamma(b\to s+X)}{\Gamma_0[|V_{ub}|^2+|V_{cb}|^2f(m_c^2/m_b^2)]} , \tag{11}$$

where $\Gamma_0=G_F^2m_b^5/192\pi^3$. The semileptonic branching ratio has been computed both with and without QCD corrections[29] and we use the average experimental value[30] of 12% for $B(B \to l\bar{\nu}+X)$. The usual phase-space correction factor $f(m_c^2/m_b^2)$ accounts for a non-negligible mass of the c quark. Here this factor is equal to 0.447.

A. $b \to s\gamma$

The photon-emission process $b \to s\gamma$ is a typical example of a one-loop flavor-changing neutral-current process. At the quark level, this decay proceeds through a one-loop "electromagnetic penguin" diagram, whereas at the exclusive level the process $B \to K^*\gamma$ is expected with the resulting γ emitted with a smooth momentum spectrum centered near $\frac{1}{2}m_b$, so that it seems possible to distinguish it from the prevalent bremsstrahlung γ. Experimentally, the best piece of information comes from ARGUS (Ref. 5), which has reported the bound $B(B \to K^*\gamma) < 4\times 10^{-4}$. Unfortunately the relation between the exclusive $B \to K^*\gamma$ and the inclusive $b \to s\gamma$ is plagued by relevant theoretical uncertainties. As reported by Ellis and Franzini,[31] two published estimates of $B \to K^*\gamma$ yield

$$\Gamma(B \to K^*\gamma)/\Gamma(b \to s\gamma) = 4.5\text{–}7\ \% \ .$$

The ARGUS limit would then imply $B(b \to s\gamma) < (6\text{–}9)\times 10^{-3}$.

The SM predicts $B(b \to s\gamma) \sim 10^{-5}$ before the inclusion of QCD corrections. However, radiative decays constitute an ideal spot for large (and even dramatic) QCD contributions. Indeed, it was shown[32] that QCD corrections are absolutely crucial. They change the Glashow-Iliopoulos-Maiani[33] (GIM) suppression in the amplitude from the typical form of a power law $\sim (m_q^2/m_W^2)\ln(m_q^2/m_W^2)$ (where m_q is the mass of the quark in the loop) to a softer form with a leading logarithmic term $\sim \ln(m_t^2/m_c^2)$.

This leads to an enhancement of the $B(b \to s\gamma)$ by approximately 1 order of magnitude (this enhancement decreases for increasing values of m_t: for instance, for $m_t = 45$ GeV it amounts to a factor ~ 16, while for $m_t = 80$ GeV it is a factor ~ 9). These surprisingly large corrections have been very recently recalculated and confirmed.[34] So, taking into account QCD corrections, the SM prediction becomes $B(b \to s\gamma) \sim 10^{-4}$ [for instance, for $m_t = 80$ GeV, $B(b \to s\gamma) \simeq 3\times 10^{-4}$]. It is well known that, going beyond the standard model, models which introduce a fourth generation[35] or a charged Higgs boson[12,36] have all the potential to enhance by 1 order of magnitude the branching ratios (QCD corrected) expected in the SM, whereas the supersymmetric contributions[37] could produce a higher branching ratio, a few times 10^{-3}.

The aim of this section is to derive the contributions due to the LRS approach to the electroweak theory. The most general form of the electromagnetic vertex is given by

$$T_\mu^{(\gamma)} = \frac{1}{4\pi^2} e \frac{g^2}{M_L^2} \bar{s}[(q^2\gamma_\mu - q_\mu \not{q})\tilde{F}_{1\gamma}(q^2) + i m_b \sigma_{\mu\nu} q^\nu \tilde{F}_{2\gamma}(q^2)]b \ , \qquad (12)$$

where $q = p_b - p_s$. The emission of a real photon, for which the magnetic transition form factor $F_{2\gamma}$ is solely responsible, is still an example of generalized "hard" GIM suppression. The width of this radiative decay, as given by the free quark model in which the light quark in B is considered to be a spectator, has been calculated and gives

$$\Gamma^{\mathrm{LRS}}(b \to s\gamma) = \frac{3\alpha}{2\pi}\Gamma_0 \left| \sum V_{is}^* V_{ib} F_{2\gamma} \right|^2 , \qquad (13)$$

where $F_{2\gamma} = F_{2\gamma}^{LL} + F_{2\gamma}^{RR} + F_{2\gamma}^{LR}$.

If we assume $M_R > 1.6$ TeV, the penguin diagram where W_L is replaced by W_R is strongly suppressed. The leading contribution arises through the W_L-W_R mixing in the loop shown in Fig. 1. Three contributions have been considered:

$$H_a^{LR}(b \to s\gamma) = C M_W^2 \cos\xi \sin\xi \times \sum_j V_{js}^* V_{jb} m_j \left[\frac{1}{M_1^2} F_{2\gamma a}^{LR}(x_{j1}) - \frac{1}{M_2^2} F_{2\gamma a}^{LR}(x_{j2}) \right] \times O_L^T \ , \qquad (14)$$

$$H_b^{LR}(b \to s\gamma) = C M_W^2 \cos\xi \sin\xi \times \sum_j V_{js}^* V_{jb} m_j \left[\frac{1}{M_1^2} F_{2\gamma b}^{LR}(x_{j1}) - \frac{1}{M_2^2} F_{2\gamma b}^{LR}(x_{j2}) \right] \times O_L^T \ , \qquad (15)$$

$$H_c^{LR}(b \to s\gamma) = C M_W^2 \cos\xi \sin\xi \times \sum_j V_{js}^* V_{jb} m_j \left[\frac{1}{M_1^2} F_{2\gamma c}^{LR}(x_{j1}) - \frac{1}{M_2^2} F_{2\gamma c}^{LR}(x_{j2}) \right] \times O_R^T \ , \qquad (16)$$

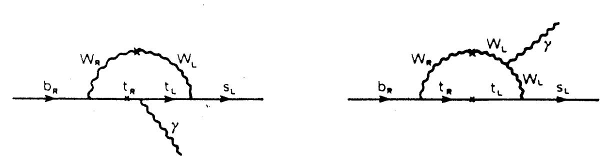

FIG. 1. Typical penguin diagrams with W_L-W_R mixing contributing to the transition $b \to s\gamma$ in LRS models.

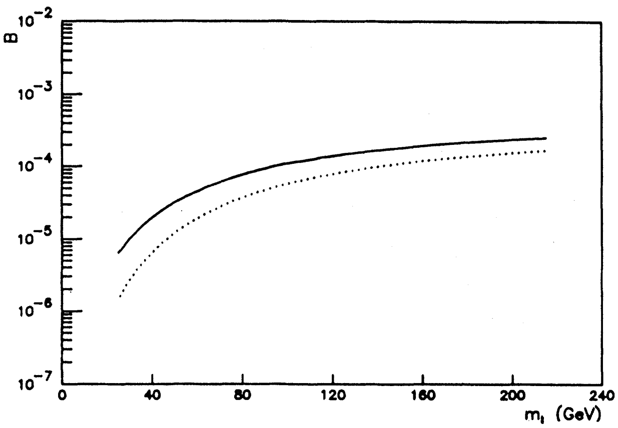

FIG. 2. The branching ratio of the B meson radiative decay in LRS models, for $\xi=4\times10^{-3}$ and $M_R=1.6$ TeV. The dashed line corresponds to the result of the standard model.

where

$$F_{2\gamma a}^{LR}(x)=\frac{1}{1-x}\left[\frac{2(1+x)}{(1-x)}+\frac{4x\ln x}{(1-x)^2}\right],$$

$$F_{2\gamma b}^{LR}(x)=\frac{1}{x-1}\left[-\frac{3(3x-1)}{2(x-1)}+\frac{3x^2\ln x}{(x-1)^2}\right], \qquad (17)$$

$$F_{2\gamma c}^{LR}(x)=\frac{1}{x-1}\left[\frac{(3x-1)}{4(x-1)}-\frac{x^2\ln x}{2(x-1)^2}\right].$$

In the above formulas $C=(e/4\pi^2)(G_F/\sqrt{2})$, $x_{j1}=(m_j/M_1)^2$, $x_{j2}=(m_j/M_2)^2$, and the two left and right operators are given by $O_{L,R}^T=\bar{s}i\sigma_{\mu\nu}\epsilon^\mu q^\nu P_{L,R}b$. Equations (14) and (15) refer to the diagrams of Fig. 1, while Eq. (16) corresponds to an unphysical-Higgs-boson exchange.

The different contributions to the magnetic form factor can be expressed as

$$F_{2\gamma}^{LL}=f_{2\gamma}(x_i),\quad F_{2\gamma}^{RR}=f_{2\gamma}(\beta x_i)$$

and (18)

$$F_{2\gamma}^{LR}=\cos\xi\sin\xi\frac{m_i}{m_b}[f_{2\gamma}^{LR}(x_i)-\beta f_{2\gamma}^{LR}(\beta x_i)],$$

where $x_i=m_i^2/M_1^2$, $\beta=M_1^2/M_2^2$ and the internal loop factors are estimated to be

$$f_{2\gamma}(x)=-2\left[\frac{x(8x^2+5x-7)}{12(x-1)^3}-\frac{x^2(3x-2)}{2(x-1)^4}\ln x\right], \qquad (19)$$

$$f_{2\gamma}^{LR}(x)=\left[\frac{(13-7x)}{4(1-x)^2}+\frac{(8-5x)}{2(1-x)}\ln x\right].$$

If we consider the strictest bound on this mixing, $\xi<4\times10^{-3}$, the contribution from the diagram of Fig. 1 never has a chance to prevail over the SM contribution. However, given the uncertainties on this bound and the fact that the branching ratio depends on ξ^2, we cannot rule out a possible enhancement in LRS models. Clearly, if we were to allow a few percent for ξ, as some authors claim, it would not be difficult to regain the 1 order of magnitude enhancement of other classes of theories. The behavior of $B(b\to s\gamma)$ in LRS models (including the usual contribution from W_L exchange) in terms of m_t is reported in Fig. 2, where we have fixed $\xi=4\times10^{-3}$ and $M_R=1.6$ TeV.

B. $b\to sl^+l^-$ and $b\to s\nu\bar{\nu}$

Another interesting and experimentally clear test of the one-loop effects in the standard model is the rare B decay $B\to Kl^+l^-$ where $l=e$ or μ (Ref. 38), which proceeds with a branching ratio of the order 10^{-6}. CLEO (Ref. 39) puts the bound $B(B\to l^+l^-+X)<1.2\times10^{-3}$ on the inclusive process $b\to sl^+l^-$. It is worth noting that the presence of a fourth generation could increase the branching ratio appreciably to perhaps 1 order of magnitude.[40]

The quark-level process $b\to sl^+l^-$ occurs at the one-loop level by means of the so-called Λ, R, and box classes of Feynman diagrams shown in Fig. 3. These contributions include (i) the diagrams of Figs. 3(a) and 3(b) with either a photon or a Z as the virtual gauge boson which creates the $l\bar{l}$ pairs, (ii) the dominant $W_2^+W_1^-$ exchange diagrams shown in Fig. 3(c), so the total weak amplitude

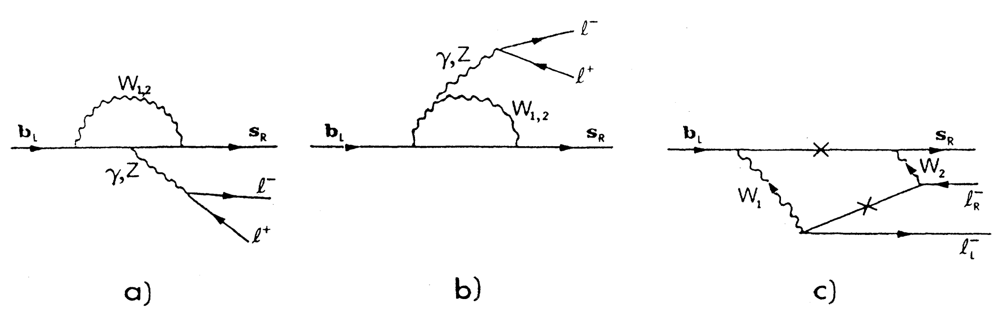

FIG. 3. The dominant γ-exchange, Z-exchange, and box diagrams contributing to the transition $b\to sl^+l^-$ in LRS models. We consider the induced vertices where the resulting γ and Z can be attached either (a) to the internal fermionic lines (Λ graphs) or (b) to the W_1, W_2 bosonic lines (R diagrams). (c) shows the extra box graphs with W_L, W_R exchanges.

can be written as the sum of the three different contributions:

$$A(b\rightarrow sl^+l^-)=A^{\gamma}+A^{Z}+A^{WW} . \tag{20}$$

As was already discussed, the electromagnetic transition vertex has two independent components, a magnetic part $F_{2\gamma}$ ($q^2=0$) and a charge radius part $F_{1\gamma}$ which contributes to the (infinite) renormalizations and vanishes for transitions to a real photon. The complete formula in the Feynman-'t Hooft gauge for the charge radius $F_{1\gamma}$ transition form factor is given, in analogy with the previous case, by $F_{1\gamma}\simeq F_{1\gamma}^{LL}+F_{1\gamma}^{LR}$, where the computation of the LL coupling gives the Inami-Lim result[38] and the contribution $F_{1\gamma}^{LR}$ can be extracted from the graphs of Figs. 3(a) and 3(b). The calculation of the LR contribution yields

$$F_{1\gamma}^{LR}\simeq\frac{e}{4\pi}\frac{G_F}{\sqrt{2}}\cos\xi\sin\xi\sum_i V_{is}^*V_{ib}m_i \times[Q_tF_{1\gamma}^{\Lambda}(x_i)-F_{1\gamma}^{R}(x_i)] \tag{21}$$

with

$$F_{1\gamma}^{\Lambda}(x)=\frac{17+8x-x^2}{9(1-x)^3}+\frac{2+6x}{3(1-x)^5}\ln x ,$$
$$F_{1\gamma}^{R}(x)=\frac{17x^3+8x^2-x}{12x(x-1)^4}-\frac{x^3+3x^2}{2(x-1)^5}\ln x . \tag{22}$$

Let us consider now the Z-exchange contribution to the process $b\rightarrow sl^+l^-$ shown in Figs. 3(a) and 3(b). The dominant contribution to the one-loop renormalized $Zb\bar{s}$ effective vertex (besides the standard LL contribution[38]) comes from the mixing between W_L and W_R. The results from the one-loop diagrams can be summarized as

$$\Gamma_Z=\frac{1}{4\pi^2}\frac{G_F}{\sqrt{2}}\cos\xi\sin\xi\sum_j V_{js}^*V_{jb}m_j\left\{\frac{g}{\cos\theta_W}(1-4Q\sin^2\theta_W)[\Gamma_1^{\Lambda}(x_j)-\beta\Gamma_1^{\Lambda}(\beta x_j)]\bar{s}(\gamma_\mu m_bP_R+i\sigma_{\mu\nu}q^\nu P_L)b\right.$$
$$\left.-g\cos\theta_W\{[\Gamma_1^R(x_j)-\beta\Gamma_1^R(\beta x_j)]\bar{s}\gamma_\mu m_bP_Rb+[\Gamma_2^R(x_j)-\beta\Gamma_2^R(\beta x_j)]\bar{s}\sigma_{\mu\nu}q^\nu P_Lb\}\right\} \tag{23}$$

with

$$\Gamma_1^{\Lambda}(x)=\frac{1-x^2+2x\ln x}{2(1-x)^3},\quad \Gamma_1^R(x)=-3\Gamma_1^{\Lambda}(x) ,$$
$$\Gamma_2^R(x)=\frac{3(1-3x)}{2(x-1)^2}\frac{3x^2\ln x}{(x-1)^3} . \tag{24}$$

Consider now the box diagrams with the exchange of two W bosons. The dominant contribution comes from the W_L^+-W_R^- exchange amplitude in Fig. 3(c). We find that this contribution can be written as

$$A^{WW}(b\rightarrow sl^+l^-)=\frac{g^4}{64\pi^2M_L^2}\sum_{jl}V_{js}^*V_{jb}B(x_j,y_{l\alpha}) \times[(\bar{s}\gamma_\alpha P_Lb\bar{l}\gamma^\alpha P_Ll) +(L\rightarrow R)] , \tag{25}$$

where $x_i=m_i^2/M_L^2$ with $i=u,c,t$ and $y_l=m_{\nu_l,\alpha}^2/M_L^2$ with l the lepton index and α selecting the heavy or light neutrino sector. Finally $B(x_j,y_{l\alpha})$ expresses in a synthetic form the loop momenta integration. In the case of the W_L-W_R contribution we may write

$$B(x_j,y_{l\alpha})=m_{q_j}m_{\nu_l\alpha}[(4+x_jy_{l\alpha}\beta)I_1(x_j,y_{l\alpha},\beta) -(1+\beta)I_2(x_j,y_{l\alpha},\beta)] , \tag{26}$$

where the loop integrals I_1 and I_2 are given by

$$I_1(x_i,x_j;z)=\frac{1}{x_i-x_j}\left[\frac{x_i\ln x_i}{(1-x_i)(1-zx_i)}-\frac{x_j\ln x_j}{(1-x_j)(1-zx_j)}\right]-\frac{z\ln z}{(1-z)(1-zx_i)(1-zx_j)} ,$$
$$I_2(x_i,x_j;z)=\frac{1}{x_i-x_j}\left[\frac{x_i^2\ln x_i}{(1-x_i)(1-zx_i)}-\frac{x_j^2\ln x_j}{(1-x_j)(1-zx_j)}\right]-\frac{\ln z}{(1-z)(1-zx_i)(1-zx_j)} . \tag{27}$$

It is worth noting that an estimate of this box contribution requires a detailed knowledge of the neutrino sector. In the LRS models with Dirac neutrinos, the contribution is at an unobservable level. On the other hand, if the neutrino is a Majorana particle, the box contribution is of the order $\beta\ln(m_N^2/M_R^2)$. Within the stated approach to the neutrino sector, the dominant results from one-loop contributions to semileptonic decays come from the electric form factor $F_{1\gamma}$ only for the small top-quark mass. Furthermore, the eventual existence of a heavy quark addresses the question of the importance of the top-quark mass effect that arises from the W_i exchange in the loop corrections to the $Zb\bar{s}$ vertex. The full dependence on the internal quark mass for the$B(b\rightarrow sl^+l^-)$ of this decay mode is shown in Fig. 4. The process $b\rightarrow s\nu\bar{\nu}$ can be

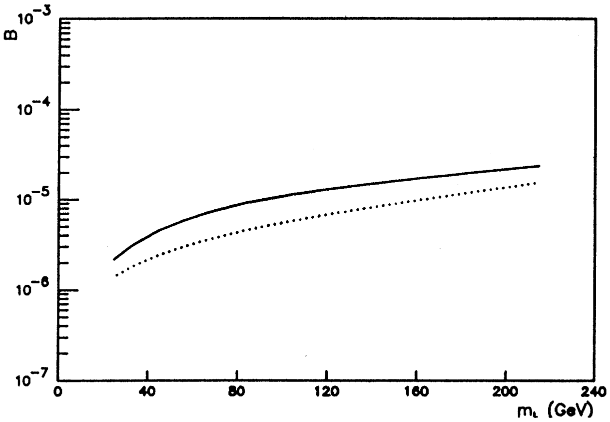

FIG. 4. Branching ratio of the process $b \to se^+e^-$ in LRS models, for $\xi = 4\times10^{-3}$ and $M_R = 1.6$ TeV. The dashed line corresponds to the result of the standard model.

easily extracted from the previous one, taking simply the loop-induced effective Z vertex and the box contribution. Needless to say, experimentally $b \to s\nu\bar{\nu}$ poses a much more formidable challenge than $b \to sl^+l^-$. As for the extension in LRS models, the resultant inclusive branching ratio (Fig. 5) does not present any significant enhancement.

C. $b \to s$"g"

This corresponds to a variety of processes leading to charmless final states which include strange mesons. According to the q^2 carried out by "g" we have on-shell single gluon emission ($q^2=0$), or two quark final states ($q^2>0$ or $q^2<0$) or two gluon final states ($q^2>0$). The $q^2>0$ case gives the dominant contribution in the SM, with "g"$\to gg$ of the same order (if not dominant) with respect to "g"$\to q\bar{q}$. In any case, the expected branching ratios in the SM for $B \to$ (K + charmless particles) cannot exceed (1–2)%.

At the level of the hadronic mode there are several

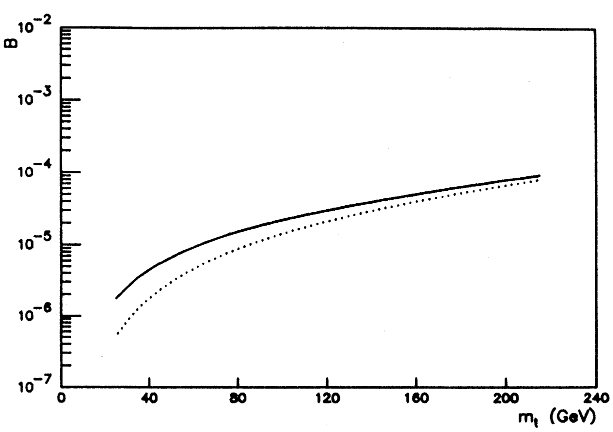

FIG. 5. $B(b \to s\nu\bar{\nu})$ in LRS models, for $\xi = 4\times10^{-3}$ and $M_R = 1.6$ TeV. The dashed line corresponds to the result of the standard model.

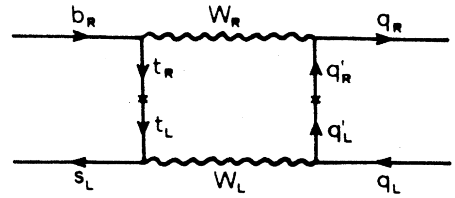

FIG. 6. Typical box diagrams contributing to the transition $b \to sq\bar{q}$ in LRS models.

theoretical and experimental difficulties. Experimentally no bound on this kind of charmless B decay is available. The best one can do is to put together the inclusive K yield reported by CLEO (Ref. 41) and the inclusive charm production at $\Upsilon(4S)$ reported by CLEO (Ref. 42) and ARGUS (Ref. 1). At present, due to the experimental and to possible Monte Carlo uncertainties, a 10–20 % $B \to K$ without charm cannot be ruled out.[43]

The decay can be analyzed in a way similar to the $b \to s\gamma$ decay. The standard model gives rise to a flavor-changing gluonic emission by means of the same diagrams of Fig. 1 where now a gluon replaces the photon. As already stressed, the QCD-induced one-loop b decays to charmless final states involve the penguinlike process $b \to s$"g", where "g" may be lightlike (i.e., on its mass shell, $q^2=0$), timelike (with emission of a $q\bar{q}$ pair with $q = u,d,s$, or of two on-shell gluons), or spacelike (with a real penguin process). The transition $b \to sg$ is a "magnetic" transition involving only the "magnetic" (i.e., F_2^g) form factor since the electric form factor does not contribute. The loop-induced $b \to sg$ coupling (all external momenta are ignored whenever possible and the helicity projection part is taken into account) can be written as ($i = u,c,t$)

$$T_\mu^{(g)} = g_s \frac{2\sqrt{2}G_F}{16\pi^2} V_{js}^* V_{jb} \bar{s} \frac{\lambda^a}{2} \times [(q^2\gamma_\mu - q_\mu \not{q})\tilde{F}_{1g}(q^2) + i m_b \sigma_{\mu\nu} q^\nu \tilde{F}_{2g}(q^2)] b \ , \tag{28}$$

the form factors being extracted from the effective $b \to s\gamma$ coupling with the result

$$F_{1g}^{LR}(x) = F_{1\gamma}^{\Lambda}(x), \quad F_{2g}^{LR}(x) = F_{2\gamma a}^{LR} \ . \tag{29}$$

As expected, it turns out that $b \to sg$ at $q^2=0$ receives a negligible contribution from loops with W_L-W_R mixing of $\sim 10^{-4}$–10^{-5}. Even taking $\xi = 0.06$ this contribution becomes only barely comparable to the SM results. QCD corrections may, however, change this picture. With $\xi = 4\times10^{-3}$ and $m_t = 75$ GeV the LRS models yield $B(b \to sg(q^2=0)) \simeq 10^{-3}$ (including the SM contribution). Now, if the QCD corrections again produce an enhancement of almost 1 order of magnitude, we see that, already with $\xi = 4\times10^{-3}$, the contribution LR would reach the same order of magnitude of that of the SM; then an enhancement is possible for $\xi > 4\times10^{-3}$.

In the $q^2>0$ case, where the penguinlike diagrams are dressed with external $q\bar{q}$ legs, we must consider both the penguin and box contributions. We can readily dispose of penguin contributions. They can be (at most) of the same order of the corresponding diagrams in the SM with a helicity flip (i.e., the term proportional to $\sigma_{\mu\nu}q^{\nu}$). But we know that in the SM the penguins without helicity flip (i.e., those proportional to $q^{\mu}q^{\nu}-q^2g^{\mu\nu}$) dominate (because of the presence of the GIM large logarithmic factors) and, hence, we can safely disregard the LR penguins.

Coming to the box diagrams (see the example in Fig. 6), their contribution to $\Gamma(b \rightarrow sg)$ is represented by the effective Hamiltonian

$$H_{\square}^{LR} = \frac{G_F^2 M_W^2}{4\pi^2} 2\beta \sum_{ij} H_{ij} \sqrt{x_i x_j} \times [(4+\beta x_i x_j) I_1(x_i, x_j; \beta) + (1+\beta) I_2(x_i, x_j; \beta)] O_{LR}^{PS} , \quad (30)$$

where $\beta=(M_{W_L}/M_{W_R})^2$, $O_{LR}^{PS}=\bar{s}_L b_R \bar{q}_L q_R$ while we put $H_{ij}=V_{it}^* V_{is} V_{jq_1}^* V_{jq_2}, x_i=(m_i/M_{W_L})^2$ and the loop integrals have been already defined in Eq. (27). This last contribution satisfies the transition form factors, though it should be stressed that at the inclusive level the effects of the box diagrams never overcome the penguinlike processes for a mass of the right-handed gauge boson in the TeV region.

IV. CONCLUSIONS

A detailed analysis of the rare processes $b \rightarrow s\gamma$, $b \rightarrow sg$, $b \rightarrow sl^+l^-$, and $b \rightarrow s\nu\bar{\nu}$ within the framework of the LRS models has been presented, with an explicit estimate of the dominant one-loop graphs contributing to the above processes. An essential point in the analysis is the choice of the parameters in terms of which the charged right-handed currents are described. On the other hand, the experimental absence of right-handed interactions imposes a suppression of the W_R mediated processes. This is achieved in two ways: increasing the W_R mass and restricting the amount of mixing ξ between W_L and W_R. There is, in principle, a third, model-dependent way of realizing this suppression, through small mixing angles in the right-handed sector. If one considers generic left-right-symmetric models there is a rather large amount of arbitrariness in the interplay among the three aforementioned ways of suppressing right-handed current interactions. For definiteness, in our analysis we kept to the simplest case, namely, the situation of "manifest" left-right symmetry where the mixing angles in the right-handed sector coincide with the CKM angles of the usual left-handed charged currents. Even restricting ourselves to the small class of models with manifest left-right symmetry there is still some arbitrariness on the most stringent lower bounds on M_R and ξ. As we emphasized, the limits $M_R > 1.6$ TeV and $\xi < 4\times10^{-3}$ rest on some assumptions and, thus, there is still room for some surprise.

We have shown that if one takes the above quoted bounds in the analysis of models with manifest left-right symmetry, then the possible enhancements for rare B decays become rather marginal. At most, the LR contributions (i.e., the one-loop diagrams where both W_L and W_R are present instead of only W_L as in the SM) can be of the same order as the ones which are present in the SM case. However, it is obvious that these results are quite sensitive to the values of M_R, ξ, and CKM angles in the right-handed sector. For instance, the diagrams with W_L-W_R mixing lead to rates for rare B decays which depend on ξ^2, and, thus, there is ample room for conspicuous enhancements if one lets ξ vary from 4×10^{-3} (which is derived making use of current-algebra in strangeness-changing hadronic decays) to 6×10^{-2} (which is the direct bound from μ decay). Needless to say, the same holds true if one releases the strong assumption of manifest left-right symmetry. Thus, the full expressions of the W_L-W_R-mediated one-loop contributions to the rates of rare B decays that we have displayed in this paper can be used to impose constraints on left-right-symmetric models. Alternatively, should we find experimentally any departure from the SM predicted rates, our analysis should be used together with complementary tests on LR models to ascertain whether some class of these models may be responsible for such an enhancement.

In conclusion, we think that, analogously to what occurred for other extensions of the SM (in particular the supersymmetric one), rare B physics represents also for LR models an ideal place to establish severe constraints. Conversely, and admittedly, in a more hopeful attitude, rare B decays have some chance to constitute the first ground where departures from the SM predictions are observed. If this is the case, our analysis shows that LR models are still there to provide a valid alternative to the SM in the realm of new physics in the TeV region.

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