First limit on inclusive $B \to X_s \nu \bar{\nu}$ decay
and constraints on new physics

(Erratum)

Yuval Grossman$^a$, Zoltan Ligeti$^b$ and Enrico Nardi$^a$

$^a$Department of Particle Physics
Weizmann Institute of Science, Rehovot 76100, Israel

$^b$California Institute of Technology, Pasadena, CA 91125

(July 16, 1996)

In a recent article [1] we presented the derivation of the first limit on the inclusive $B \to X_s \nu \bar{\nu}$ decay rate. Our work stemmed from the observation that a published ALEPH bound on BR($B \to \tau \bar{\nu}$) [2], inferred from the absence of large missing energy events in $B$ decays, implies also a limit on BR($B \to X_s \nu \bar{\nu}$). We estimated this limit by comparing the missing energy spectra in the two decay modes. Theoretically, $B \to X_s \nu \bar{\nu}$ is a very clean process, which is also sensitive to several possible sources of new physics [1]. Therefore, a dedicated experimental search is important. The result of such an analysis by the ALEPH Collaboration, using the full LEP–I data sample, is to appear soon [3].

As a result of discussions with members of the ALEPH Collaboration [4], we found two errors in our numerical results. The corrected analysis yields a limit weaker than our original result [1] by about a factor of three. The corrections included here are formally beyond the free quark decay result at tree level (of order $\alpha_s$, $\Lambda_{QCD}/m_b$, and higher). However, they affect the final limit significantly, due to the specific details of the experimental analysis.

First, due to a mistake in our Monte Carlo code, the energy of the $X_s$ hadronic final state in the $B$ rest frame was evaluated as a fraction of the $b$ quark mass, rather than as a fraction of the $B$ meson mass. This resulted in a harder missing energy spectrum than the true one. Correcting for this weakens the limit by about 35%.

Second, we neglected the effects of the invariant mass distribution of the final hadrons,
which we (implicitly) assumed not to extend to high mass states. However, the invariant mass of the $X_s$ system can be well above 1 GeV [5]. Because of the large boost into the LEP laboratory frame and the high missing energy range (> 35 GeV) used in the analysis [2] (which is large compared to the average $B$ meson energy at LEP), neglecting this effect yielded a missing energy spectrum in the laboratory frame considerably harder than the correct one.

It is not straightforward to include properly this second effect into the analysis. However, an estimate can be obtained by approximating the invariant mass spectrum of the $X_s$ system with a Gaussian distribution with mean ($\mu$) and variance ($\sigma$) fitted to the averages $\langle M_X^2 \rangle$ and $\langle M_X^4 \rangle$, as given in [5]. Such a fit yields $\mu \simeq 1.35$ GeV and $\sigma \simeq 0.6$ GeV. There is about a 10% uncertainty in these values, due to their dependences on the heavy quark effective theory parameter $\bar{\Lambda}$ ($\simeq m_B - m_b$). Treating the $X_s$ mass distribution as independent of the energy spectrum, and including a theoretical uncertainty of about 15% related to the fitted values of $\mu$ and $\sigma$, we find that the original limit given in [1] is weakened by about a factor of three. While small values of $E_X$ favor small values of $M_X$ (beyond the trivial kinematic constraint $E_X > M_X$), the correlation between $E_X$ and $M_X$ only slightly improve the bound.

In conclusion, the bound on BR($B \to X_s \nu \bar{\nu}$) given in the Abstract and in Eqs. (1.2) and (6.1) of [1] is weaker by about a factor of three, while the limits on the new physics parameters, collected in Table 2, are weaker by about a factor 1.8. Until the ALEPH bound on the $B \to X_s \nu \bar{\nu}$ branching ratio will appear, the limit BR($B \to X_s \nu \bar{\nu}$) < $1.3 \times 10^{-3}$ should be used, instead of that quoted in [1].
REFERENCES

[1] Y. Grossman, Z. Ligeti and E. Nardi, Nucl. Phys. B 465 (1996) 369 [see also the corresponding entry in: R.M. Barnett et al., Particle Data Group, Phys. Rev. D 54 (1996) 1].

[2] D. Buskulic et al., ALEPH Collaboration, Phys. Lett. B 343 (1995) 444.

[3] P. Perrodo of the ALEPH Collaboration, private communication.

[4] We thank Ian Tomalin for useful discussions, and especially Pascal Perrodo for comments that led us to find the errors.

[5] A.F. Falk, M. Luke, and M.J. Savage, Phys. Rev. D 53 (1996) 2491.
First limit on inclusive $B \to X_s \nu \bar{\nu}$ decay 
and constraints on new physics

Yuval Grossman$^a$, Zoltan Ligeti$^b$ and Enrico Nardi$^a$

$^a$Department of Particle Physics 
Weizmann Institute of Science, Rehovot 76100, Israel

$^b$California Institute of Technology, Pasadena, CA 91125

Abstract

The inclusive $B \to X_s \nu \bar{\nu}$ decay rate, on which no experimental bound exists to date, can be constrained by searching for large missing energy events in $B$ decays. Carefully examining the experimental and theoretical aspects of such an analysis, we argue that the published ALEPH limit on $\text{BR}(B \to \tau \bar{\nu})$ implies, conservatively, the bound $\text{BR}(B \to X_s \nu \bar{\nu}) < 3.9 \times 10^{-4}$, which is less than one order of magnitude above the standard model prediction. The LEP collaborations could significantly improve this bound by a dedicated experimental analysis. We study the constraints this new limit imposes on various extensions of the standard model. We derive new bounds on the couplings of third generation fermions in models with leptoquarks, and in supersymmetric models without R-parity. We also constrain models where new gauge bosons are coupled dominantly to the third generation, such as TopColor models and models based on horizontal gauge symmetries. For models which predict an enhanced effective $bsZ$ vertex, the constraint from $B \to X_s \nu \bar{\nu}$ is competitive with the limits from inclusive and exclusive $B \to X_s \ell^+ \ell^-$ decays.
I. INTRODUCTION

Recent progress in experiment and theory has made flavor changing neutral current (FCNC) $B$ decays a stringent test of the Standard Model (SM) and a powerful probe of New Physics (NP). The CLEO Collaboration observed the exclusive decay $B \rightarrow K^{*} \gamma$ as well as the inclusive decay $B \rightarrow X_{s} \gamma$. The UA1 upper limit on the inclusive decay $B \rightarrow X_{s} \mu^{+} \mu^{-}$, and the recent CLEO and CDF upper limits on the exclusive decays $B \rightarrow K^{(*)} \ell^{+} \ell^{-}$, are less than one order of magnitude above the SM predictions. These decays are likely to be observed within the next few years. These experimental results constitute a set of very strong constraints on several possible sources of NP.

The FCNC decay $B \rightarrow X_{s} \nu \bar{\nu}$ is also very sensitive to extensions of the SM, and provides a unique source of constraints on some NP scenarios which predict a large enhancement of this decay mode. In particular, the $B \rightarrow X_{s} \nu_{\tau} \bar{\nu}_{\tau}$ mode is very sensitive to the relatively unexplored couplings of third generation fermions. However, no experimental upper bound on this decay mode has been established to date.

The decay $B \rightarrow X_{s} \nu \bar{\nu}$ can be searched for through the large missing energy associated with the two neutrinos. Using such techniques, the ALEPH, L3, and OPAL collaborations have been able to measure the inclusive $B \rightarrow X_{c} \tau \bar{\nu}$ decay rate. The large missing energy in this case is associated with the two neutrinos in the decay chain $B \rightarrow X_{c} \tau \bar{\nu}$ followed by $\tau \rightarrow \nu X$. From the absence of excess events with very large missing energy, ALEPH also established the 90%CL bound

$$\text{BR}(B \rightarrow \tau \bar{\nu}) < 1.8 \times 10^{-3}.$$  \hspace{1cm} (1.1)

In this paper we point out that a similar analysis of the same data also implies a limit on $\text{BR}(B \rightarrow X_{s} \nu \bar{\nu})$. While a detailed and complete analysis of this set of data can only be performed by the ALEPH Collaboration, we show that using some conservative and simplifying assumptions it is possible to derive the bound

$$\text{BR}(B \rightarrow X_{s} \nu \bar{\nu}) < 3.9 \times 10^{-4},$$  \hspace{1cm} (1.2)
which is less than one order of magnitude above the SM prediction \[11\]

\[
\text{BR}^{\text{SM}}(B \to X_s \nu \bar{\nu}) \approx 5 \times 10^{-5}.
\] (1.3)

We expect that a dedicated analysis of the LEP collaborations will strengthen our bound. If background subtraction can be performed, or extra experimental cuts can reduce the background, then we estimate that using the full LEP–I data sample a 90\% CL bound of order

\[
\text{BR}(B \to X_s \nu \bar{\nu}) < (1 - 2) \times 10^{-4},
\] (1.4)

could be within the reach of the LEP experiments.

In section II we discuss, in a model independent manner, the inclusive $B \to X_s \nu \bar{\nu}$ decay rate and the missing energy spectrum. In section III we describe the theoretical issues involved in relating the limits on large missing energy events to the $B \to X_s \nu \bar{\nu}$ decay rate, and we estimate the corresponding theoretical uncertainties. This section also contains the details of the derivation of our bound (1.2). While the limit we obtain does not allow a direct test of the SM, it still implies stringent constraints on various new physics scenarios. In section IV we analyze the implications of existing measurements of FCNC processes, particularly $B \to X_s \gamma$ and $B \to X_s \mu^+ \mu^-$, for $B \to X_s \nu \bar{\nu}$, in various classes of NP models. Section V contains an extensive discussion of the constraints implied by the bound (1.2) on various extensions of the SM. We derive new constraints on models in which the couplings of the third family fermions differ from those of the first two generations, and we present numerical limits on several NP parameters. Finally, section VI contains a summary of our results and the conclusions.

II. THE $B \to X_s \nu \bar{\nu}$ DECAY RATE

From the theoretical point of view, the decay $B \to X_s \nu \bar{\nu}$ is a very clean process, since both the perturbative $\alpha_s$ and the non-perturbative $1/m_b^2$ corrections are known to be small.
Furthermore, in contrast to the decay $B \to X_s \ell^+ \ell^-$, which suffers from background such as $B \to X_s J/\psi \to X_s \ell^+ \ell^-$, there are no analogous long-distance QCD contributions, since there are no intermediate states that can decay into a neutrino pair. Therefore, the decay $B \to X_s \nu \bar{\nu}$ is well suited to search for and constrain NP effects.

As our aim is to derive constraints on NP scenarios, we discuss the missing energy spectrum in a model independent framework. Limits on NP parameters can then be derived by comparing the experimental bound with the theoretical expressions as derived in specific models. A model independent expression for the missing energy spectrum can be straightforwardly obtained from the general result for muon decay \cite{12}. Under the only assumption of two component left-handed neutrinos (possible neutrino mass effects are at most of order $m_\nu/m_b < 10^{-2}$), the most general form of the four-fermion interaction responsible for $B \to X_q \nu_i \bar{\nu}_j$ reads

$$L = C_L O_L + C_R O_R , \quad (2.1)$$

where

$$O_L = [\bar{q}_L \gamma_\mu b_L] [\bar{\nu}_L^i \gamma^\mu \nu_L^j] , \quad O_R = [\bar{q}_R \gamma_\mu b_R] [\bar{\nu}_L^i \gamma^\mu \nu_L^j] . \quad (2.2)$$

Here $L$ and $R$ denote left- and right-handed components, $q = d, s$, and $i, j = e, \mu, \tau$. In this article we adopt the notation that a generic $B$ meson contains a $b$ quark, rather than a $\bar{b}$ quark. As the flavors of the decay products are not detected, in certain models more than one final state can contribute to the observed decay rate. Then, in principle, both $C_L$ and $C_R$ carry three indices $q, i, j$, which label the quark and neutrino flavors in the final state. Throughout our discussion, we shall only keep track of these indices when they are important, otherwise we will suppress them.

At lowest order, the missing energy spectrum in the $B$ rest-frame is given by \cite{12}

$$\frac{d\Gamma(B \to X_q \nu_i \bar{\nu}_j)}{dx} = \frac{m_b^5}{96\pi^3} \left( |C_L|^2 + |C_R|^2 \right) S(r, x) . \quad (2.3)$$

Here we have not yet summed over the neutrino flavors. The function $S(r, x)$ describes the shape of the missing energy spectrum
The dimensionless variable $x = E_{\text{miss}}/m_b$ can range between $(1 - r)/2 \leq x \leq 1 - \sqrt{r}$, and $r = m_{\tau}^2/m_b^2$. The parameter $\eta = -\text{Re}(C_L C_R^*)/(|C_L|^2 + |C_R|^2)$, which is the analog of the Michel parameter in $\mu$-decays, ranges between $-\frac{1}{2} \leq \eta \leq \frac{1}{2}$.

In the SM, $B \to X_s \nu \bar{\nu}$ proceeds via $W$ box and $Z$ penguin diagrams, therefore only $O_L$ is present and $\eta = 0$. The corresponding coefficient reads $^{13}$

$$C_L^{\text{SM}} = \frac{\sqrt{2} G_F}{\pi \sin^2 \theta_W} \frac{\alpha_s}{4 \pi} V_{tb}^* V_{ts} X_0(x_t),$$

where $x_t = m_{\tau}^2/m_W^2$, and

$$X_0(x) = \frac{x}{8} \left[ \frac{2 + x}{x - 1} + \frac{3x - 6}{(x - 1)^2} \ln x \right].$$

In the limit of large top quark mass, $X_0$ has a quadratic dependence on $m_t$, $X_0(x_t) \sim x_t/8$. Therefore, the main source of uncertainty in the SM prediction for the total decay rate comes from the uncertainty in $m_t$.

The leading $1/m_b^2$ and $\alpha_s$ corrections to the SM result (calculated in the free quark decay model) are known. The $\alpha_s$ correction to the total decay rate is given by replacing $X_0(x_t)$ in Eq. (2.5) by $^{11}$

$$X_0(x_t) \to \left[ X_0(x_t) + \frac{\alpha_s}{4 \pi} X_1(x_t) \right] \left[ 1 - \frac{\alpha_s}{3 \pi} \left( \frac{\pi^2}{4} - \frac{25}{4} \right) \right].$$

Here the second term represents the correction to the matrix element of $O_L$. This term cancels almost completely when the $B \to X_s \nu \bar{\nu}$ branching ratio is normalized to the semileptonic $B \to X_c e \nu \bar{\nu}$ rate (see Eq. (2.9) below). The first term contains the QCD correction to the box and penguin diagrams. We do not display here the explicit form of $X_1(x_t)$, which can be found in $^{11}$. The most important effect of this correction is to reduce the scale dependence of the SM prediction from about $\pm 10\%$ to below $\pm 2\%$ $^{11}$.

The $1/m_b^2$ correction to the contribution of the $O_L$ operator to the missing energy spectrum can be read off from $^{14}$. The result is the following modification of the function $S(r, x)$ in Eq. (2.4)
\[ S(r, x) \rightarrow S(r, x) \]
\[ \frac{1}{\sqrt{(1-x)^2-r}} \left\{ \frac{\lambda_1}{6m_b^2} [2(1-x)^2(19 - 38x + 10x^2) - r(61 - 23r - 122x + 52x^2)] \right\} \]
\[ + \frac{\lambda_2}{2m_b^2} [(1-x)(47 - 126x + 96x^2 - 20x^3) - r(70 - 23r - 125x + 52x^2)] \right\} \]
\[ + \left[ \frac{\lambda_1}{48m_b^2} (5 - r) - \frac{\lambda_2}{16m_b^2} (1 - 5r) \right] (1 - r)^3 \delta \left( \frac{1}{2}(1-r) - x \right) \]
\[ + \frac{\lambda_1}{96m_b^2} (1 - r)^5 \delta' \left( \frac{1}{2}(1-r) - x \right) . \]

(2.8)

Here \( \lambda_1 \) and \( \lambda_2 \) are related to the kinetic energy of the \( b \) quark inside the \( B \) meson and to the mass splitting between the \( B \) and the \( B^* \) mesons, respectively. Experimentally, \( \lambda_2 \approx 0.12 \text{ GeV}^2 \), and following the discussion in [15], we use \( 0 < -\lambda_1 < 0.5 \text{ GeV}^2 \). When integrated over the spectrum, this correction amounts to about 3% suppression of the total decay rate. Even this small correction cancels almost entirely when the \( B \to X_s \nu \bar{\nu} \) branching ratio is normalized to the semileptonic \( B \to X_c e \bar{\nu} \) rate (see Eq. (2.9) below). As we will discuss in the next section, although the \( B \to X_s \nu \bar{\nu} \) branching fraction in the SM is known rather precisely, the theoretical prediction for the missing energy spectrum is more uncertain.

In several NP models \( O_R \) is also present. The structure of the operator product expansion [16, 18] shows that to all orders in the \( \alpha_s \) and \( 1/m_b \) expansions, both the perturbative and non-perturbative corrections to the contribution of \( O_R \) to the missing energy spectrum are identical to those of \( O_L \). This holds as long as the phase space is symmetric in the two leptons (this can be violated only by negligible neutrino mass effects). Thus, the shape of the missing energy spectrum is model independent, up to possible small interference effects between \( O_L \) and \( O_R \), of order \( m_s/m_b \). Once \( C_L \) and \( C_R \) are computed in any particular model, the unknown contributions to the total decay rate are only \( \mathcal{O}(\alpha_s^2; \alpha_s m_s/m_b; \alpha_s \Lambda^2/m_b^2; \Lambda^3/m_b^3; m_s \Lambda^2/m_b^3) \), where \( \Lambda \) denotes some scale of order \( \Lambda_{\text{QCD}} \).

In some NP models \( O_L \) and \( O_R \) are simultaneously present with coefficients of comparable size, giving rise to interference effects proportional to \( \eta \), which modify the shape of the missing energy spectrum. Due to the \( m_s/m_b \) suppression, this effect is always small, except close to the endpoint region \((x \sim 1 - \sqrt{r})\), where the leading term in (2.4) also becomes of
order $\sqrt{r}$. Therefore, close to the endpoint, the corrections to the shape of the spectrum could be relevant. Thus, besides the quark mass ratio, an additional uncertainty in (2.3) is associated with the value of $\eta$. The softest missing energy spectrum, and thus the most conservative bound on the $B \to X_s \nu \bar{\nu}$ branching ratio, is obtained by using a large quark mass ratio $r \simeq 0.002$ (corresponding to $m_s \simeq 0.2$ GeV and $m_b \simeq 4.8$ GeV), and $\eta = -\frac{1}{2}$. The bound on the branching fraction derived using these values of $r$ and $\eta$ holds in any model.

In order to find the constraints on NP, we will express the total decay rate in terms of two “effective” coefficients $\tilde{C}_L$ and $\tilde{C}_R$, which can be computed in terms of the parameters of any NP model and are directly related to the experimental measurement (see (2.12)). To remove the large uncertainty in the total decay rate associated with the $m_b^5$ factor, it is convenient to normalize $\text{BR}(B \to X_s \nu \bar{\nu})$ to the semileptonic rate $\text{BR}(B \to X_c e \bar{\nu})$, since the experimental value of the latter is known quite precisely. The contribution from $B \to X_u e \bar{\nu}$, as well as possible NP effects on the semileptonic decay rate are negligible. In constraining NP, we can also set $m_s = 0$ and neglect both order $\alpha_s$ and $1/m_b^2$ corrections. This is justified, since when averaged over the spectrum these effects are very small, and would affect the numerical bounds on the NP parameters only in a negligible way. For the total $B \to X_q \nu_i \nu_j$ decay rate into all possible $q = d, s$ and $i, j = e, \mu, \tau$ final state flavors, we then obtain

$$
\frac{\text{BR}(B \to X \nu \bar{\nu})}{\text{BR}(B \to X_c e \bar{\nu})} = \frac{\tilde{C}_L^2 + \tilde{C}_R^2}{|V_{cb}|^2 f_{PS}(m_c^2/m_b^2)},
$$

where $f_{PS}(x) = 1 - 8x + 8x^3 - x^4 - 12x^2 \ln x$ is the usual phase-space factor, and we defined

$$
\tilde{C}_L^2 = \frac{1}{8G_F^2} \sum_{q,i,j} |C_{qi}^{Lj}|^2, \quad \tilde{C}_R^2 = \frac{1}{8G_F^2} \sum_{q,i,j} |C_{qi}^{Rj}|^2.
$$

The SM prediction for the branching ratio in Eq. (1.3) is obtained by inserting into Eq. (2.3) the semileptonic rate $\text{BR}(B \to X_c e \bar{\nu}) \approx 10.5\%$ [19], $f_{PS}(m_c^2/m_b^2) \approx 0.5$, together with

$$
\left(\tilde{C}_L^{\text{SM}}\right)^2 \approx 3.8 \times 10^{-7},
$$

which follows from (2.3) by using $|V_{tb}^* V_{ts}|/|V_{cb}| \approx 1$, $|V_{cb}| \approx 0.04$, $\alpha(m_Z) \approx 1/129$, $\sin^2 \theta_W \approx 0.23$, and $m_t \approx 180$ GeV [14]. In a SM analysis it would be more natural to factor out
from the definitions of $\tilde{C}_{L,R}$ the small mixing angles $V_{tb}^* V_{ts}$ as well as $\alpha$, resulting in a dimensionless coefficient of order unity. However, there is no reason to do so for $\tilde{C}_R$, and since in several NP models $\tilde{C}_{L,R}$ are induced at the tree level, they do not contain any small parameters analogous to $\alpha$ or to the CKM angles. In fact, often the only suppression factors in these coefficients come from inverse powers of some large mass scale.

Comparing the limit (1.2) with (2.9) yields the following constraint on possible NP contributions:

$$\tilde{C}_L^2 + \tilde{C}_R^2 < 3.0 \times 10^{-6} \left[ \frac{\text{BR}(B \to X_s \nu \bar{\nu})}{3.9 \times 10^{-4}} \right]. \quad (2.12)$$

In section V we will constrain NP models by comparing this limit with the various theoretical predictions for the coefficients $\tilde{C}_{L,R}$. We will quote the bounds that can be derived on each NP parameter, even in those cases when $O_L$ and $O_R$ are simultaneously present and certain combinations of the NP parameters may be better constrained.

**III. EXPERIMENTAL ANALYSIS**

In this section we discuss the theoretical uncertainties involved in the experimental analysis at LEP, and we derive the bound (1.2). To obtain a limit on $B \to X_s \nu \bar{\nu}$ using the data from ALEPH [7], we have to estimate the relative efficiency of $B \to X_s \nu \bar{\nu}$ and $B \to \tau \bar{\nu}$ to pass the experimental missing energy cut (ALEPH used the cut $E_{\text{miss}} > 35 \text{ GeV}$). We use the $B \to X_s \nu \bar{\nu}$ missing energy spectrum as given in Eqs. (2.3) and (2.4), together with the conservative values $r = 0.002$ and $\eta = -\frac{1}{2}$ which, as discussed in the previous section, lead to the softest missing energy spectrum. As a result, the bound we derive holds in any model.

**A. Theoretical uncertainties**

The largest theoretical uncertainty is related to the reliability of the theoretical calculation of the missing energy spectrum near its endpoint ($x \sim 1 - \sqrt{r}$). While both the $\alpha_s$
and $1/m_b^2$ corrections to the total decay rate are small, this is not true point-by-point for the differential energy spectrum.

The $1/m_b$ expansion for inclusive semileptonic decays of hadrons containing a single $b$ quark based on an operator product expansion is not reliable when the invariant mass of the decay product hadronic system is small. In the endpoint region, near $x = 1 - \sqrt{r}$, there are large corrections to the free quark decay prediction, and the spectrum has to be smeared to get a reliable result. This smearing region has to be chosen large enough so that after smearing, the $1/m_b^2$ corrections in (2.8) produce only small effects on the decay spectrum. Following this criterion, we find that in the present case the smearing region extends about $0.5 - 0.7$ GeV from the maximal value of the missing energy.

The order $\alpha_s$ corrections to the missing energy spectrum can be read off from, and are less problematic. While there are large logarithms for small values of the hadronic invariant mass, the perturbative correction to the hadron energy spectrum is smooth in the limit of vanishing hadronic energy. Thus, the perturbative corrections to the missing energy spectrum are small and reliably calculable in the region relevant for our analysis.

The large non-perturbative corrections near maximal missing energy could represent a problem for the analysis of $B \to X_s \nu \bar{\nu}$, since the experimental cuts select events with large missing energy, thus enhancing the weight of the problematic region. However, the experimental search is at the endpoint region in the laboratory frame. Events with large missing energy in the laboratory frame are not necessarily those with missing energy close to the endpoint in the $B$ rest-frame. We need to take into account the large boost from the $B$ rest-frame to the laboratory frame. We define $P_{E_{\text{cut}}}[\text{GeV}](E_{\text{cut}}[\text{GeV}])$ as the fraction of the events with $E_{\text{miss}} > E_{\text{cut}}$ in the laboratory frame, that come from the missing energy region in the $B$ rest-frame between the endpoint and $E_0 = m_b - m_s - \epsilon$. To simulate the $B$ hadron energy in the laboratory frame, we use in our Monte Carlo code the Peterson fragmentation function, and $m_b \approx 4.8$ GeV. We find the following representative numbers.
The results in (3.1) show some sensitivity to the size of the smearing region, while they suggest that varying the experimental cut has a much smaller effect. We have also checked that these figures are not very sensitive to the choice of the quark masses and of $\eta$.

We conclude that the missing energy spectrum in the laboratory frame of the LEP experiments can be estimated reliably from the spectator model. However, the uncertainty related to the endpoint region of the missing energy spectrum in the $B$ rest-frame should be taken into account. To estimate this uncertainty, we rely on the fact that the size of the endpoint region is chosen such that the integration over it can be trusted. Assigning to the fraction of the events coming from endpoint region the lowest possible value of the missing energy, $m_b - m_s - \epsilon$, we can bound the uncertainty related to using the spectator model. For $\epsilon = 0.7\text{ GeV}$ we find in this way a 5\% reduction in the total number of events that pass the 35 GeV cut. We consider this as a reliable estimate of the theoretical uncertainty of the analysis related to the use of the spectator model.

A final remark is in order. At LEP, $b$ quarks hadronize into a variety of $b$ hadrons. However, by the time it decays, the $b$ quark is contained either in a $B$ meson or in a $\Lambda_b$ baryon. The fraction of $\Lambda_b$ baryons at LEP has been measured to be about 10\% [22]. In the limit when the $b$ quark is treated as very heavy, the missing energy spectrum from $B$ decays should be similar to that from $\Lambda_b$ decays, up to small effects originating from the polarization of the baryons. A possible indication for significant heavy quark symmetry violating effects, is the experimentally measured lifetime ratio between the $B$ and the $\Lambda_b$ hadrons, which appears to be larger than the theoretical prediction (for a recent discussion see, e.g., [23]). However, the resolution of this problem is most likely related either to experimental issues or to the theoretical calculation of the hadronic decay widths, so it is probably irrelevant for the analysis in this paper. The polarization of the $\Lambda_b$ baryons produced at LEP has been measured to be $-30 \pm 30\%$ [24]. A non-vanishing left-handed polarization enhances
the missing energy from $\Lambda_b$ decays, and therefore neglecting it is conservative.

B. Derivation of the bound

After this discussion, we are ready to estimate the bound on $\text{BR}(B \to X_s \nu \bar{\nu})$. To translate the ALEPH bound on $B \to \tau \bar{\nu}$ [7] into a limit on $B \to X_s \nu \bar{\nu}$, we need to compare the theoretical predictions for the fraction of events that pass the missing energy cut for the two decay modes. The resulting bound on $B \to X_s \nu \bar{\nu}$ is stronger than that on $B \to \tau \bar{\nu}$ for the following reasons:

(i) The $B \to \tau \bar{\nu}$ decay is allowed only for the charged $B$ mesons, while all $b$ flavored hadrons can decay through the parton level process $b \to s \nu \bar{\nu}$;

(ii) In order to reject background from semileptonic $B$ decays, only the hadronic $\tau$ decays were used in the ALEPH search for $B \to \tau \bar{\nu}$;

(iii) The missing energy spectrum is somewhat harder in $B \to X_s \nu \bar{\nu}$ than in $B \to \tau \bar{\nu}$ decays, increasing the efficiency of the analysis.

To evaluate the last factor (iii), we need to estimate the missing energy spectrum for the $B \to \tau \bar{\nu}$ decay as well. In estimating the $\tau \to \nu X$ missing energy spectrum, we consider only two-body (20%) and three-body (80%) hadronic $\tau$ decays. Since a significant fraction of hadronic tau decays are four- and five-body decays, the resulting missing energy spectrum for $B \to \tau \bar{\nu}$ is harder than the real one. Therefore we get a conservative upper bound on the $B \to X_s \nu \bar{\nu}$ branching ratio, i.e., weaker than what a detailed analysis would obtain.

We take into account that the $\tau$ polarization decreases the efficiency of the $B \to \tau \bar{\nu}$ analysis by about 20%, as estimated in the ALEPH analysis [7]. We find that in our approximation the efficiency of the $B \to X_s \nu \bar{\nu}$ decay to pass the ALEPH cut $E_{\text{miss}} > 35$ GeV is about 15% larger than that of $B \to \tau \bar{\nu}$ followed by hadronic $\tau$ decay.

*We thank Ian Tomalin for pointing this out to us.
Collecting the results of our previous discussion, we can finally estimate that the bound on $\text{BR}(B \to X_s \nu \bar{\nu})$ is stronger than the ALEPH bound on $\text{BR}(B \to \tau \bar{\nu})$ by an overall factor
\[
R = R_{B^\pm} R_{\text{hadr}} R_{\text{eff}} \approx 0.21. \tag{3.2}
\]
Here $R_{B^\pm} \approx 0.37$ is the ratio of $B^\pm$ mesons to the total number of $b$ hadrons at ALEPH \cite{22} and $R_{\text{hadr}} \approx 0.65$ is the hadronic $\tau$ branching fraction \cite{19}. The factor $R_{\text{eff}}$ accounts for the efficiency of the missing energy cut in the $B \to \tau \bar{\nu}$ decay followed by hadronic $\tau$ decays, relative to that in $B \to X_s \nu \bar{\nu}$ decays. We conservatively use the upper bound $R_{\text{eff}} = 0.90$ which includes the 5\% theoretical uncertainty related to the reliability of the spectator model spectrum. We expect the uncertainty related to the use of the Peterson fragmentation function to cancel to a large extent from the estimate of the relative efficiency of the $B \to X_s \nu \bar{\nu}$ and $B \to \tau \bar{\nu}$ analyses. Using (3.2) and the experimental bound (1.1), we find the limit given in Eq. (1.2).

We would like to emphasize that the derivation of our bound relies on a set of conservative simplifying assumptions, which could be avoided in a dedicated experimental analysis. Our limit (1.2) should be considered as a conservative upper bound, to which we purposely do not (and cannot) assign a confidence level. We hope that a more complete and detailed investigation by the LEP collaborations will be carried out. The reward of such an analysis could be a bound of order $(1 - 2) \times 10^{-4}$, that is only a factor $2 - 4$ above the SM prediction.

**IV. CONSTRAINTS ON $B \to X_s \nu \bar{\nu}$ FROM OTHER PROCESSES**

Models that can give rise to large new contributions to the $B \to X_s \nu \bar{\nu}$ decay often predict an enhancement of other FCNC processes as well. In this section we analyze what constraints the existing experimental data on FCNC processes imply for various NP models. For each specific model, these constraints result in upper limits on the allowed $B \to X_s \nu \bar{\nu}$ decay rate. Bounds on $\text{BR}(B \to X_s \nu \bar{\nu})$ can be derived from the limits on rare processes, such as $K_L \to \mu^+ \mu^-$, $K \to \pi \nu \bar{\nu}$, $\epsilon_K$, $K - \bar{K}$ and $B - \bar{B}$ mixing. The most restrictive
constraints are imposed by the measurements of the radiative decay $B \to X_s \gamma$, and by the limits on inclusive and exclusive $B \to X_s \ell^+ \ell^-$ decays.

The radiative $B \to X_s \gamma$ decay proceeds via photon penguin diagrams, and therefore it is not directly related to $B \to X_s \nu \bar{\nu}$. However, in many models the details of the underlying physics imply relations between the $Z$ and the photon penguins. In all these models the recent CLEO measurement \[2\]

\[
\text{BR}(B \to X_s \gamma) = (2.32 \pm 0.51 \pm 0.32 \pm 0.20) \times 10^{-4},
\]

(4.1)

which is in agreement with the SM, forbids large deviations from the SM prediction for the $B \to X_s \nu \bar{\nu}$ decay rate as well.

On the other hand, a large class of NP models predict (or can accommodate) an enhanced $bsZ$ effective vertex without giving rise to a large enhancement of the $bs\gamma$ effective coupling. Then the constraints from inclusive and exclusive $B \to X_s \ell^+ \ell^-$ decays are important, as these decays, like $B \to X_s \nu \bar{\nu}$, are dominated by $Z$ exchange. In these models, a naive estimate of the ratio of inclusive rates gives $\text{BR}(B \to X_s \nu \bar{\nu})/\text{BR}(B \to X_s \ell^+ \ell^-) \approx 6$. The factor of six enhancement arises due to a factor of approximately two in the ratio between the neutrino and the charged lepton couplings to the $Z$, and a factor of three from the sum over the neutrino flavors. A more precise calculation which includes the photon exchange contribution to $B \to X_s \ell^+ \ell^-$ bounded by the CLEO measurement of $B \to X_s \gamma$, as well as the sizable QCD corrections to $B \to X_s \ell^+ \ell^-$ \[25,26\], can increase the above ratio up to 7. Hence, for this class of models, the UA1 experimental limit on the inclusive $B \to X_s \mu^+ \mu^-$ decay \[3\].

\[\text{†}\]

\[\text{†}\]The UA1 experiment searched for events in the region $3.9 < E_{\mu\mu} < 4.4 \text{ GeV}$. However, the theoretical prediction for the spectrum is uncertain in this endpoint region. Moreover, the theoretical spectrum shown in Ref. \[3\] (and presumably used to relate the bound on the number of events in the experimental window to the quoted bound on the total branching fraction) seems to be in disagreement with that calculated in \[25\]. Despite these uncertainties we use the published bound.
\[
\text{BR}(B \rightarrow X_s \mu^+ \mu^-) < 5 \times 10^{-5},
\]

implies
\[
\text{BR}(B \rightarrow X_s \nu \bar{\nu}) \lesssim 3.5 \times 10^{-4} \quad \Longrightarrow \quad \tilde{C}_L^2 + \tilde{C}_R^2 < 2.7 \times 10^{-6}.
\]

This limit is comparable with the limit (4.2), however it is weaker than the expected LEP sensitivity (1.4). This underlines the importance of a more detailed experimental analysis aimed at searching for \(B \rightarrow X_s \nu \bar{\nu}\), and shows that LEP measurements could compete with future new data from CLEO and CDF in constraining models which predict an enhanced \(bsZ\) coupling. Moreover, as we have emphasized, the neutrino mode is particularly interesting since it is theoretically cleaner than the charged lepton modes: (\(i\)) there are no long-distance QCD effects; (\(ii\)) the short-distance QCD corrections are small; (\(iii\)) there is no photon penguin contribution, and therefore this process can be straightforwardly related to the effective \(bsZ\) vertex. In conclusion, although the constraints resulting from our limit (4.2) are numerically slightly weaker than those implied by \(B \rightarrow X_s \mu^+ \mu^-\) (4.3), we think that they are more reliable.

The exclusive dilepton decays are also sensitive to an enhanced \(bsZ\) vertex, and provide additional constraints. The exclusive decay modes \(B \rightarrow K^{(*)} \ell^+ \ell^-\) have been searched for by the CLEO \(^4\) and the CDF \(^5\) collaborations. For our purposes, the most restrictive limit has been established by CLEO \(^4\)
\[
\text{BR}(B \rightarrow K^* e^+ e^-) < 1.6 \times 10^{-5}.
\]

However, the interpretation of this limit is obscured by the significant model dependence of the exclusive \(B \rightarrow K^{(*)}\) form factors. The estimates for the ratio
\[
\rho \equiv \frac{\text{BR}(B \rightarrow K^* e^+ e^-)}{\text{BR}(B \rightarrow X_s e^+ e^-)},
\]
range between 0.10 and 0.35 \(^27\). Clearly, the value of \(\rho\) is crucial for relating the limits on exclusive decays to a limit on the FCNC transition at the quark level, for which the SM
predicts \( 25 \). \( \frac{\text{BR}(B \to X_s e^+ e^-)}{\text{BR}(B \to X_s \mu^+ \mu^-)} \approx 10 \times 10^{-6}, \quad \text{BR}(B \to X_s \mu^+ \mu^-) \approx 7 \times 10^{-6}. \) (4.6)

It is also questionable whether the above estimates of \( \rho \), based on SM computations, are reliable for the analysis of NP scenarios. In fact, in the SM a significant contribution to \( B \to K^* e^+ e^- \) comes from soft photons, corresponding to small \( e^+ e^- \) invariant mass. However, in the NP models under consideration \( Z \) penguins dominate, and therefore \( \rho \) is expected to be smaller than in the SM.

The only known model independent way to predict the \( B \to K^{(*)} \) form factors (besides lattice calculations) is to relate them to measurable semileptonic decay form factors using heavy quark symmetries \( 28 \). However, until experimental information on the \( B \to \rho \ell \bar{\nu} \) decay spectrum becomes available, the dilepton spectrum can only be predicted in the small window \( 4.0 < m_{\ell^+\ell^-} < 4.4 \text{ GeV} \) from the \( D \to K^* \ell^+\ell^- \nu \) data. (Even after the measurement of the \( B \to \rho \ell \bar{\nu} \) spectrum, the unknown symmetry breaking corrections will amount to an uncertainty of order 20 – 30%.)

Rare \( K \) decays also imply bounds on \( Z \) penguins. However, in relating them to the \( B \to X_s \nu \bar{\nu} \) decay rate large uncertainties arise from poorly known quark mixing angles. The only measured rate is \( K_L \to \mu^+ \mu^- \), which receives large long-distance corrections. Therefore, the short-distance parameters can only be extracted with large uncertainties \( 29 \). The decay \( K^+ \to \pi^+ \nu \bar{\nu} \), however, receives negligible long-distance contributions. While the existing experimental bound is about fifty times the SM prediction, an order of magnitude improvement is expected from Brookhaven in the coming years. In summary, at present, for all the models in our analysis, rare \( K \) decays are less constraining than the rare \( B \) decays discussed above.

Once these considerations are taken into account, it appears that the most reliable con-

\( \dagger \)The spread in the SM predictions in the literature is in part due to the different signs for the \( O_7 - O_9 \) interference term in various papers. (We use here the notation of Ref. \( 26 \).)
constraints on an anomalous effective $bsZ$ vertex, apart from $B \to X_s \nu \bar{\nu}$, come from the limits on the inclusive $B \to X_s \mu^+ \mu^-$ decay rate. Numerically similar bounds (or possibly slightly better – depending on the adopted values of model-dependent hadronic form factors) are provided by the bounds on exclusive $B \to K(\nu) \ell^+ \ell^-$ decay rates.

V. NEW PHYSICS

Following the discussion in the previous section, it is useful to classify the NP models which could enhance the SM prediction for $B \to X_s \nu \bar{\nu}$ into three classes: (A) Highly constrained models; (B) Weakly constrained models; (C) Unconstrained models.

The first class (A) includes models in which the existing bounds on other FCNC processes (mainly $B \to X_s \gamma$) imply that the rate for $B \to X_s \nu \bar{\nu}$ cannot exceed the SM prediction by any factor larger than two. Thus, a bound on $B \to X_s \nu \bar{\nu}$ of the order of the LEP sensitivity ($1.4$) does not imply any new constraint on the underlying NP. For completeness, we briefly explain below why the $B \to X_s \nu \bar{\nu}$ branching fraction must be close to its SM value in the most popular models belonging to this class: the minimal supersymmetric standard model; multi Higgs doublet models; left-right symmetric models.

To class (B) belong all models which predict (or allow for) a large effective $bsZ$ (or $bsZ'$) vertex. We call these models “weakly constrained”, as the limits on inclusive and exclusive $B \to X_s \ell^+ \ell^-$ already constrain them. The limits from (1.2) will not represent any numerical improvement over the existing bounds on the parameters of these models. However, our new limits are more reliable, since the theoretical uncertainties involved in $B \to X_s \nu \bar{\nu}$ are very small. Moreover, if a bound of the order (1.4) can be obtained by the LEP collaborations, then the constraints from $B \to X_s \nu \bar{\nu}$ will become markedly stronger than the present ones from $B \to X_s \ell^+ \ell^-$. A list of interesting NP models belonging to this class includes: models with additional $Q = -\frac{1}{3}$ isosinglet quarks; models with large $bsZ'$ vertex; models with an anomalous effective $tcZ$ vertex; models with a heavy fourth generation; models with anomalous $WWZ$ couplings; a class of extended technicolor models. These models will be
discussed in some detail below.

In the unconstrained models of class (C), the couplings responsible for enhancing $B \rightarrow X_s \nu \bar{\nu}$ are to a large extent independent of those constrained by any other existing experimental bound. Therefore, even a $B \rightarrow X_s \nu \bar{\nu}$ decay rate orders of magnitude above the SM prediction is still consistent with the existing constraints, and the new limit (1.2) represents the most stringent bound on the corresponding NP parameters. Examples of theoretically interesting models belonging to this class are: models with light leptoquarks; supersymmetric models with broken R-parity; TopColor models; some models based on non-Abelian horizontal gauge symmetries.

A. Highly constrained models

1. Minimal Supersymmetric Standard Model (MSSM)

In a large part of the SUSY parameter space, the MSSM (for review and notations see, e.g., [30]) is known to produce large effects on the radiative decay $B \rightarrow X_s \gamma$, as well as on $B \rightarrow X_s \ell^+ \ell^-$ (see [31,32]). The effects on $B \rightarrow X_s \nu \bar{\nu}$ have been studied in [32,33]. It was found that the contributions to the rate can be non-negligible only for tan $\beta$ close to unity, while for increasing values of tan $\beta$ the prediction rapidly converges to the SM value, regardless of the values of the other SUSY parameters. Even for tan $\beta \approx 1$ it seems to be problematic to enhance the rate up to the level observable at LEP, while keeping BR($B \rightarrow X_s \gamma$) within the experimental limits. This is due to the fact that the SUSY contributions to the $bs\gamma$ vertex tend to dominate over those to the $bsZ$ vertex. The SUSY corrections to the $B \rightarrow X_s \gamma$ decay can be kept small while allowing for a large BR($B \rightarrow X_s \nu \bar{\nu}$), only for low values of tan $\beta$, and by invoking large cancellations between the charged Higgs and chargino contributions. It is then conceivable that with a specific choice of several SUSY parameters, a “fine-tuned” MSSM could produce BR($B \rightarrow X_s \nu \bar{\nu}$) close to (1.4). However, we regard such a choice as unnatural.
2. Multi Higgs Doublet Models (MHDM)

MHDM (for review and notation see, e.g., [31,34]) are severely constrained by $B \to X_s \gamma$, $Z \to b\bar{b}$, $B \to X_c \tau \bar{\nu}$, and lepton universality in tau decays. The same is true for the more familiar two Higgs doublet models (2HDM), which represent a subclass of the general MHDM with natural flavor conservation. In a general MHDM the single parameter $\tan \beta$ of the 2HDM is replaced by three complex coupling constants, $X$, $Y$, and $Z$, which describe the Yukawa interactions of the lightest charged scalar with the down-type quarks, up-type quarks, and charged leptons, respectively. The new $Z$ penguin diagrams present in these models are related to new photon penguins, and thus are severely constrained by $B \to X_s \gamma$ and cannot contribute significantly to $B \to X_s \nu \bar{\nu}$. A large enhancement not affecting $B \to X_s \gamma$ and $B \to X_s \mu^+ \mu^-$ could arise for large enough values of the $H^\pm \tau \nu_\tau$ Yukawa coupling via charged scalar box diagrams involving two external $\nu_\tau$’s. These box diagrams are proportional to the combination $m_\tau^2 Z^2 (m_t^2 Y^2 + m_b m_s X^2)$ [34]. However, $Y$ is constrained from $Z \to b\bar{b}$, the coupling $Z$ from lepton universality in tau decays, and the product $XZ$ from $B \to X_c \tau \bar{\nu}$ [34,15,35], implying together that also the contribution of the box diagrams to $B \to X_s \nu \bar{\nu}$ has to be very small.

3. Left–Right Symmetric Models (LRSM)

In these models (for review and notations see, e.g., [36]) new heavy charged $W_R$ gauge bosons coupled to right-handed currents are exchanged in a new set of electroweak penguin diagrams, which could give rise to deviations in the rate for $B \to X_s \nu \bar{\nu}$. However, new photon penguins would enhance $B \to X_s \gamma$ as well. The $B \to X_s \gamma$ decay rate has been calculated in the minimal as well as in some non-minimal versions of the LRSM. Once the limits on $M_{W_R}$ and on the $W_L - W_R$ mixing angle are imposed (from direct $W_R$ searches, $K_L - K_S$ mass difference, etc.), the rate for $B \to X_s \gamma$ cannot differ from the SM prediction by more than about 50% [37]. The CLEO measurement of $B \to X_s \gamma$ further constrains
the parameter space of non-minimal models to regions where the $SU(2)_R$ gauge coupling is small. In this region of parameters, the $Z$ penguins are similarly suppressed, and therefore significant deviations of the $B \to X_s \nu \bar{\nu}$ decay rate from the SM prediction are excluded.

B. Weakly constrained models

1. Mixing of the $b$ with new exotic $Q = -\frac{1}{3}$ quarks

It is well known that the presence of new heavy fermions with non-canonical $SU(2)$ transformations ($L$-handed singlets and/or $R$-handed doublets) mixed with the standard leptons and quarks would give rise to tree level FCNC in $Z$ interactions [38]. In particular, the presence of new $Q = -\frac{1}{3}$ isosinglet quarks, $D_L$ and $D_R$, as they appear for example in the 27 representation of $E_6$ and in several superstring inspired extensions of the SM, would generate a $b_L s_L Z$ vertex. New $SU(2)$ doublets $(U^D)_L$, $(U^D)_R$ mixed with the $R$-handed $d$-type quarks would give rise to the new FCNC operator $b R s R Z$, which is absent in the SM. Both effects could appear simultaneously in the presence of a set of multiplets of mirror fermions $(U^D)_R$, $U_L$, $D_L$.

It is easy to see why such flavor changing vertices are generated. For each $L,R$ chirality state, the vector $\Psi^o_{L,R} = (d^o, D^o)_{L,R}$ of the ordinary $d^o$ and new exotic $D^o$ quarks couples to the $Z$ through the matrix of the isospin charges

$$T^3_L = -\frac{1}{2} \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix}, \quad T^3_R = -\frac{1}{2} \begin{pmatrix} 0 & 0 \\ 0 & 1 \end{pmatrix},$$

which, in contrast to the SM, are not proportional to the identity matrix. Then the isospin part of the neutral current for the mass eigenstates, $\Psi_{L,R} = U^j_{L,R} \Psi^o_{L,R}$, which contains the flavor changing part of the coupling, reads

$$\mathcal{L}_{\text{mix}} = \frac{g}{2 \cos \theta_W} \sum_{i \neq j} \left[ \Psi^i_L \kappa^{ij}_L \gamma^\mu \Psi^j_L + \Psi^i_R \kappa^{ij}_R \gamma^\mu \Psi^j_R \right] Z^\mu .$$

Here and henceforth the hermitian conjugate is understood. The strength of the flavor changing coupling $\kappa^{ij}_{L,R}$ is given by
\[- \frac{1}{2} \kappa_{L,R}^{ij} \equiv \left[ U_{L,R}^\dagger T_{L,R}^3 U_{L,R} \right]_{ij}, \quad (i \neq j). \quad (5.3)\]

As a result, a mixing of the $b$ with exotic $d$-type quarks will induce a $b q Z$ ($q = d, s$) vertex at tree level. We stress that a mixing is allowed only between states which carry the same quantum number of an unbroken gauge group. Therefore, for the corresponding photon coupling we have $U^\dagger Q U \propto I$ (with $Q = -\frac{1}{3} I$). Namely, the electromagnetic couplings remain flavor diagonal. The same holds for the part of the $Z$ coupling proportional to the electromagnetic generator.

The effective interaction (2.1) arising from (5.2) yields the coefficients

\[ C_{L,R} = \frac{g^2}{4 \cos^2 \theta_W M_Z^2} \kappa_{L,R}^{bq} \quad \Rightarrow \quad C_{L,R}^2 = \frac{3}{4} \sum_{q=d,s} |\kappa_{L,R}^{bq}|^2, \quad (5.4) \]

where the factor of 3 arises from the sum over the neutrino flavors. The current bounds on $|\kappa_{L,R}^{bq}|$ are obtained from the limit on $B \rightarrow X_s \mu^+ \mu^-$ (for $q = s$) and from $B_d - \bar{B}_d$ mixing (for $q = d$) \cite{39,40}. Numerically, they are comparable with the bound implied by the limit (1.2)

\[ |\kappa_{L,R}^{bq}| < 2.0 \times 10^{-3}. \quad (5.5) \]

In some mass matrix models, the flavor changing couplings generated by the ordinary-exotic fermion mixing are expected to be of the order of the ratio of the light to heavy mass scales \cite{38} (see however \cite{41} for some interesting exceptions). We see that the limits on $b s Z$ transitions are sensitive to $b$ mixing with heavy particles up to a mass scale $M \sim \mathcal{O}(\sqrt{m_b m_s/\kappa_{bs}})$, namely about 500 GeV.

2. Models with large $bsZ'$ vertex

Quite in general, new fermions have to be present in models based on extended gauge groups (rank $> 4$), since they are needed to ensure the absence of anomalies in the new gauge currents. $E_6$ models are a well known example where new $SU(2)$ singlets $D_L$ and $D_R$ are present, together with new neutral gauge bosons corresponding to the two additional
Cartan generators of the group. We generically denote the new gauge bosons as \( Z' \), which are coupled to the fermions through charges \( Q' \). A mixing with new ordinary quarks (i.e., with conventional \( SU(2) \) quantum number assignments) would not affect the couplings of the \( Z \). However, as the matrix of the \( Q' \) charges is in general not proportional to the identity, \( \kappa'_{L,R}^{ij} = (U_{a}^{\dagger}Q'_{L,R}U_{a})_{i\neq j} \) will not vanish, thus inducing a \( bsZ' \) vertex.

In spite of the \( 1/M_{Z'}^{2} \) suppression, the \( Z' \) mediated FCNC are expected to be as large as the corresponding transitions induced by \( Z \) exchange \(^{[12]} \). This is because while in general the flavor changing mixings affecting the \( Z \) couplings are suppressed as the ratio of the light to heavy fermion masses, no analogous suppression is expected for the mixings between fermions of the same isospin, which affect only the \( Z' \) couplings. The absence of the suppression in the mixing can compensate for the \( M_{Z}^{2}/M_{Z'}^{2} \) suppression of the \( Z' \) relative to the \( Z \) amplitude, implying that the coefficients describing the \( Z \) and the \( Z' \) flavor changing effective interactions can be comparable in size \(^{[12]} \).

The analysis of the NP effects on an enhanced \( bsZ' \) vertex parallels quite closely that of the large \( bsZ \) models. However, in the former case the enhancement of the \( B \to X_{s}\nu\bar{\nu} \) decay mode relative to \( B \to X_{s}\ell^{+}\ell^{-} \) is not as large as in the latter one. Since \( \nu_{L} \) and \( \ell_{L} \) appear in an \( SU(2) \) doublet, they have the same \( Q'_{L} \) charge. Then the ratio of the rates of the two processes is given by

\[
\frac{\text{BR}(B \to X_{s}\nu\bar{\nu})}{\text{BR}(B \to X_{s}\mu^{+}\mu^{-})} \approx \frac{Q'_{\nu L}^{2} + Q'_{\mu L}^{2} + Q'_{\tau L}^{2}}{Q'_{\nu L}^{2} + Q'_{\mu L}^{2} + Q'_{\tau L}^{2}}. \tag{5.6}
\]

In most models, the \( Z' \) couples universally to all three generations, and then the above ratio cannot be larger than 3. (This upper bound corresponds to the case when the right-handed leptons are almost decoupled from the \( Z' \), namely \( Q'_{R} \sim 0 \). This can happen in \( E_{6} \) models in which the \( Z' \) arises as a particular combination of the two additional \( U(1)' \) generators.) As a result, these models are significantly more constrained by the limit on \( B \to X_{s}\mu^{+}\mu^{-} \) \(^{[12]} \) than the models with a large \( bsZ \) vertex. The UA1 limit implies

\[
\text{BR}(B \to X_{s}\nu\bar{\nu}) \lesssim 1.8 \times 10^{-4},
\]

which is only marginally within the reach of the expected LEP sensitivity \(^{[14]} \).
The expressions for the $C_{L,R}$ coefficients are similar to (5.4)

$$C_{L,R} = \frac{g'^2}{M_{Z'}^2} \mathcal{F}_{L,R}(Q') \kappa'^{bq}_{L,R} \implies \tilde{C}_{L,R}^2 = 3 \left[ 2 r_g \cos^2 \theta_W \frac{M_Z^2}{M_{Z'}^2} \mathcal{F}_{L,R}(Q') \right]^2 \sum_{q=d,s} |\kappa'^{bq}_{L,R}|^2.$$  

(5.7)

Here $\kappa'^{ij}_{L,R}$ describe the strengths of the FCNC couplings of the $Z'$ (the analogs of (5.3)), and we defined

$$\mathcal{F}_{L,R}(Q') = Q'(\nu) [Q'(d_{L,R}) - Q'(D_{L,R})], \quad r_g = \frac{g'^2}{g^2}.$$  

(5.8)

For any particular $Z'$ model, $\mathcal{F}_{L,R}$ and $r_g$ are known. Then the limit (1.2) implies constraints on the ratios $\kappa'^{ij}_{L,R}/M_{Z'}^2$, typically about a factor two weaker than the order $10^{-3}$ limits from the present bound on $B \to X_s \mu^+ \mu^-$. Thus, in this class of models the largest allowed effects can hardly enhance the $B \to X_s \nu \bar{\nu}$ decay rate up to the estimated LEP sensitivity (1.4). However, the sensitivity of the $B \to X_s \nu \bar{\nu}$ decay rate to $Z'$ effects can be larger in a class of unconventional $E_6$ models [43]. In these models, the different generations are embedded in three fundamental $27$ representations in a generation dependent way, implying different $Q'_L$ charges for the left-handed leptons of the different families. In such a scenario the muon can be weakly coupled to the $Z'$ without implying the same for $\nu_\tau$ and $\nu_e$, and then the ratio of the $B \to X_s \nu \bar{\nu}$ and $B \to X_s \mu^+ \mu^-$ decay rates can exceed the previously derived limit of $3$. For the unconventional $E_6$ models of Ref. [43] we find that the ratio in (5.6) can be as large as $5$. In this case the constraint from $B \to X_s \mu^+ \mu^-$ (1.2) still allows $\text{BR}(B \to X_s \nu \bar{\nu})$ up to $2.9 \times 10^{-4}$.

3. Large $tcZ$ effective coupling

In several NP models, new sources of FCNC are naturally suppressed, as they are related to ratios between the masses of the fermions involved in the flavor changing transitions and some large mass scale. Due to the large value of the top mass, such a suppression might not be effective for flavor changing transitions involving the top quark [44][45], and theoretically
the presence of a large $tcZ$ vertex is indeed an open possibility (see [46–48], and references therein). This can be the case in models which predict new dynamical interactions of the top quark [49,50], in MHDM without natural flavor conservation [51], or in the presence of mixing with new $Q = \frac{2}{3}$ isosinglet quarks [48,52,53].

To date, the couplings of the top quark have not been directly measured, and a large $tcZ$ vertex is not constrained by the measurement of $B \to X_s \gamma$, neither by $B^0 - \bar{B}^0$ mixing. However, an anomalous $tcZ$ coupling will give new contributions to the effective $bsZ$ vertex [46,48]. Contributions to the $B \to X_s \nu \bar{\nu}$ decay would arise from diagrams involving loops of charged $W^\pm$ bosons and of unphysical $\phi^\pm$ scalars, with an insertion of the $tcZ$ coupling on the fermion line [46,48]. Since these diagrams are not CKM suppressed (they are proportional to $V_{tb}^* V_{cs}$), there is no additional suppression beyond the loop factor.

The case of a tree level $tcZ$ vertex induced by a mixing with new $Q = \frac{2}{3}$ isosinglets was analyzed in [48]. In this particular case the vertex arises at tree level. Therefore, after the underlying theory is fully specified, the computation of the new penguin diagrams can be performed in detail, yielding a finite result [48]. It was found that the new contributions to the $bsZ$ effective vertex are always smaller in absolute value than the SM contributions, and opposite in sign. As a result, they interfere destructively thus lowering the expected rates for $B \to X_s \ell^+ \ell^-$ and $B \to X_s \nu \bar{\nu}$. Thus, an anomalous $tcZ$ vertex induced by mixing with $Q = \frac{2}{3}$ isosinglets cannot be constrained by an upper bound on the branching ratio.

If the $tcZ$ vertex is an effective one, and the underlying theory is not specified, the expression for the loop induced $bsZ$ vertex by itself is formally divergent, and it has to be regulated. We take the result of the computation of the flavor changing penguins from [48] and we substitute the function resulting from the finite loop integration of that specific case, with a regulator $\log (\Lambda^2/m_t^2)$. We further assume that new effects dominate over the SM contributions. This yields for the induced $bqZ$ vertex

$$\Gamma_{bqZ}^{\text{eff}} \approx \frac{g}{\cos \theta_W} \left[ \frac{\alpha}{4\pi \sin^2 \theta_W} \left( V_{tb}^* \kappa_{tc} L V_{cq} \right) \log \frac{\Lambda^2}{m_t^2} \right], \quad (5.9)$$

where $\kappa_{tc} L$ parameterizes the strength of the $tcZ$ coupling. For the dimensionless coefficient
defined in (2.10) we get
\[ \tilde{C}_L^2 = \frac{3}{4} \left( \frac{\alpha}{4\pi \sin^2 \theta_W} |\kappa_{te}^L| \log \frac{\Lambda^2}{m_t^2} \right)^2 \sum_{q=d,s} |V_{tb}^* V_{cq}|^2. \]  
(5.10)

Assuming \( \Lambda \sim 1 \text{ TeV} \), \(|V_{tb}| \sim |V_{cs}| \sim 1, |V_{cd}| \simeq 0.22 \), we get from (1.2)
\[ |\kappa_{te}^L| \log \frac{\Lambda^2}{m_t^2} < 0.55 \quad \implies \quad |\kappa_{te}^L| \lesssim 0.16, \]  
(5.11)
which is comparable with the limit given in [46], once the differences in the normalization and in the overall coefficient in (5.9) are accounted for.

We have chosen to present our bound in the form (5.11) for uniformity of notations, and to allow for comparison with [46]. However, not to mislead the reader, a remark is in order. \( \kappa_{te}^L \) parameterizes an effective vertex, and thus if the heavy physics decouples from low energy, this parameter must be proportional to inverse powers of the NP scale \( \Lambda \). Therefore, even if this is not apparent from the notation we used in (5.9) and from the one used in [46], in the \( \Lambda \to \infty \) limit the \( bsZ \) coupling approaches its SM value, as it should, in order to account for the decoupling of heavy physics.

4. Fourth Generation

Extensions of the SM including a fourth generation are still an open possibility [54], and there are theoretically motivated models in which the additional neutrino is naturally heavy [55], thus escaping the LEP limit on the number of light neutrino species. It has also been shown that SUSY models with four generations are consistent with unification, and by imposing this requirement some conditions on the \( t', b', \tau', \) and \( \nu' \) spectrum have been derived [56]. In these models large CKM mixings with a heavy \( t' \) may enhance the \( Z \) and photon penguins [57].

The constraints implied by the CLEO measurement of \( B \to X_s \gamma \) were presented in [58]. After constraining the new mixing angles by the limits on deviations of the \( 3 \times 3 \) CKM matrix from unitarity, it was found that a new \( t' \) quark in the mass range \( 200 - 400 \text{ GeV} \) is
still consistent with experimental data. The possibility of having a measurable enhancement of the $Z$ mediated FCNC decays, while keeping $B \to X_s \gamma$ close to the SM prediction, relies on the fact that the effective $bs\gamma$ vertex has a logarithmic dependence on $m_{t'}$, while for $bsZ$ this dependence is quadratic. The dependence of the $B \to X_s \nu \bar{\nu}$ rate on $m_{t'}$ is numerically similar to that of $B \to X_s \mu^+ \mu^-$. The different contributions of the box diagrams tend to slightly enhance the neutrino decay over the charged lepton mode. We find that the constraint (4.2) still allows

$$\text{BR}(B \to X_s \nu \bar{\nu}) < 3.7 \times 10^{-4}.$$ (5.12)

In the presence of a heavy $t'$ quark, the SM coefficient (2.5) is modified according to

$$\tilde{C}_L^2 = \left(C_L^{\text{SM}}\right)^2 \left[1 + \frac{V_{t's}^* V_{ts}}{V_{tb}^* V_{ts}} \frac{X_0(x_{t'})}{X_0(x_t)}\right]^2,$$ (5.13)

where, for simplicity, we assumed $|V_{t's}| \gg |V_{t'd}|$ and $|V_{ts}| \gg |V_{td}|$. From the limit on $B \to X_s \nu \bar{\nu}$ (L.2) we get the bound

$$\left|1 + \frac{V_{t's}^* V_{ts}}{V_{tb}^* V_{ts}} \frac{X_0(x_{t'})}{X_0(x_t)}\right| < 2.8,$$ (5.14)

which is numerically similar to the limit implied by $B \to X_s \ell^+ \ell^-$. 

5. Anomalous $WWZ$ couplings

The $SU(2)_L \times U(1)_Y$ gauge symmetry of the SM fixes the dimension-4 operators that describe vector-boson self-couplings. This symmetry should be respected by any low energy effective theory, independently of possible NP at energies above the electroweak scale. NP may still signal itself in low energy experiments through higher dimensional operators, suppressed by inverse powers of the NP scale $\Lambda$. While dimension-6 operators modify the $WW\gamma$ and $WWZ$ vertices identically, this is no longer true for dimension-8, and higher dimensional operators [59]. Thus, as long as dimension-6 operators dominate the possible deviations from the SM, the experimental measurement of $B \to X_s \gamma$ implies that also the $bsZ$ coupling must be close to its SM value.
In general, dimension-6 operators, which are suppressed by two inverse powers of the NP scale, are expected to dominate over dimension-8 operators, suppressed by $\Lambda^{-4}$. However, in any perturbative underlying theory the dimension-6 operators can only arise from loop diagrams, and therefore they have an additional loop suppression factor of $1/16\pi^2$. In contrast, dimension-8 operators can also arise at tree level \cite{59}. Thus, if the scale of new physics is below about 2 TeV, then dimension-8 operators can dominate over the dimension-6 ones. In this case, the measurement of $B \to X_s \gamma$ implies no direct constraints on the $bsZ$ vertex.

The effects of anomalous $WWZ$ couplings on rare $K$ and $B$ decays were studied in \cite{60–63}. In general, seven coupling constants parameterize the $WWZ$ interaction \cite{64}. In low energy processes, when the external momenta can be neglected, only $g_1^Z$ and $g_5^Z$ contribute \cite{61,62}. A recent CDF measurement \cite{65} implies that $\Delta g_1^Z = g_1^Z - 1$ has to be small. Therefore, for simplicity we assume $g_1^Z = 1$, and we study only the effects of $g_5^Z$. The $Z$ penguin \cite{2.6} contribution is modified according to \cite{51–53}:

\begin{equation}
X_0(x_t) \to X_0(x_t) + g_5^Z \cos^2 W(x_t), \quad W(x) = \frac{3x(1 - x + x \ln x)}{4(1 - x)^2}.
\end{equation}

The strongest published bound on $g_5^Z$ is obtained from the measurement of $K_L \to \mu^+ \mu^-$ \cite{61}. As we have discussed, because of the large long-distance contribution to this process, such a bound is not very reliable. We prefer to quote the bound resulting from BR($B \to X_s \mu^+ \mu^-$).

For $m_t = 180$ GeV we found that the limit on this decay mode implies $-8.3 < g_5^Z < 3.7$. This constraint allows for BR($B \to X_s \nu \bar{\nu}$) up to the value in Eq. (4.3). Our bound (1.2) implies

\begin{equation}
-8.6 < g_5^Z < 4.1.
\end{equation}

\footnote{We use the sign convention of [63] for $g_5^Z$, which differs from that of [61,62].}
6. Extended Technicolor (ETC)

FCNC processes impose very strong constraints on technicolor models. Either “traditional” ETC models with a minimal set of interactions necessary for third family quark mass generation, or ETC models which incorporate a techni-GIM mechanism, yield a $B \to X_s \gamma$ rate at most slightly larger than the SM rate [66]. While at leading order the $bs\gamma$ coupling is not affected, these classes of ETC models typically induce large flavor changing $Z$ boson couplings of the form $\bar{s}_L \gamma_{\mu} b_L Z^\mu$. Techni-GIM models yield an even larger enhancement of 4-fermion interactions of the form $O_L$ in (2.2).

It was subsequently noted [67] that ETC models with a techni-GIM mechanism also predict the $B \to X_s \mu^+ \mu^-$ decay rate about a factor of 30 above the SM, violating the experimental bound (4.2), unless significant cancellations occur between various contributions. However, “traditional” ETC models predict only about a factor of 4 enhancement of $B \to X_s \ell^+ \ell^-$ over the SM. Therefore, these models are not yet excluded, and yield a similar enhancement for the $B \to X_s \nu \bar{\nu}$ decay rate, which could be within the reach of the expected LEP sensitivity (1.4).

C. Unconstrained models

1. Light leptoquarks

Leptoquarks (LQ) couple directly leptons to quarks. Such particles appear in several extensions of the SM. A comprehensive analysis of the experimental constraints on the LQ couplings has been given in [68], and is summarized in Table 15 of this reference. After summing over all the possible neutrino flavors in the final state, it turns out that the existing limits on the LQ Yukawa couplings allow for $\text{BR}(B \to X_q \nu \bar{\nu})$ up to $O(10\%)$. For the relevant couplings, our new limit (1.2) imposes much stronger constraints than the existing ones. Several types of LQ are possible, and we adopt here the notations of [68]. LQ can be scalar ($S$) or vector ($V$) particles, and can belong to different $SU(2)_L$ representations.
The $SU(2)_L$ singlets, doublets, and triplets are labeled with the lower index 0, 1/2, and 1, respectively.

The following LQ can mediate the $B \to X_s \nu \bar{\nu}$ decay:

$$S_0, \quad \tilde{S}_{1/2}, \quad S_1, \quad V_{1/2}, \quad V_1.$$ \hfill (5.17)

Written in components, the relevant scalar and vector terms in the interaction Lagrangian are \(^6\)

$$L_{LQ} = -\lambda_{iq} S_0 \bar{q}_c^i \nu_L^i S_0 + \lambda_{iq} \tilde{S}_{1/2} \bar{q}_L^i \tilde{\nu}_L^{\dagger} \tilde{S}_{1/2} - \lambda_{iq} \bar{q}_L^i \nu_L^i S_1^{\dagger} + \lambda_{iq} \bar{q}_L^i \nu_L^i V_{1/2}^{\mu\dagger} + \lambda_{iq} \bar{q}_L^i \nu_L^i V_{1}^{\mu\dagger}, \hfill (5.18)$$

where $q = d, s, b$, and it is understood that only the charge $\frac{1}{3}$ component of the $SU(2)_L$ multiplets appears in \((5.18)\). For simplicity, we assume that all the $\lambda_{iq}$ couplings are real.

Integrating out the LQ fields and Fiertz transforming, yields for the coefficients of the effective four-fermion interaction \((2.1)\) induced by \((5.18)\)

$$C_{qij}^{L,R} = \eta_{LQ} \frac{\lambda_{iq} \lambda_{j3}^{\frac{1}{3}}}{m_{LQ}^2}, \quad (q = d, s), \hfill (5.19)$$

where $\eta_{LQ} = 1/2$ (1) for scalar (vector) LQ. The coupling $C_L$ is generated through $S_0$, $S_1$, and $V_1$ exchange, while $C_R$ appears from $\tilde{S}_{1/2}$ and $V_{1/2}$. As different types of LQ can exist, both $C_L$ and $C_R$ can be simultaneously present. If in addition the LQ carry some generation index, cancellations between different generations of LQ are also possible, and then the limit \((1.2)\) constrains only the total LQ-mediated rate. For simplicity, we will restrict ourselves to the case when only one type of LQ is present, and it does not carry any generation index.

For scalar LQ the limit \((1.2)\) implies the following new bounds

$$\lambda_{iq} \lambda_{j3}^{\frac{1}{3}} < 1.1 \times 10^{-3} \left( \frac{m_{LQ}}{100 \text{ GeV}} \right)^2, \hfill (5.20)$$

while for vector LQ

$$\lambda_{iq} \lambda_{j3}^{\frac{1}{3}} < 5.7 \times 10^{-4} \left( \frac{m_{LQ}}{100 \text{ GeV}} \right)^2. \hfill (5.21)$$
These bounds are much stronger than the existing limits [68].

Other LQ couplings involving the light fermions of the first and second generations are constrained by the existing experimental data to be much smaller than the limits (5.20) and (5.21) [69]. However, if we have to learn a lesson from the hierarchy in the fermion Yukawa couplings, then it seems natural to expect a large hierarchy in the LQ couplings to the different generations as well. This is the case, for example, in models that explain the quark and lepton mass hierarchies as originating from horizontal symmetries. If LQ exist, in these models they couple more strongly to the third generation fermions [70]. Since three third generation fields and only one from the second generation participate in the process \( B \to X_s \nu_\tau \bar{\nu}_\tau \), any improvement in the search for the \( B \to X_s \nu \bar{\nu} \) decay would represent an important test of these models.

2. SUSY with broken R-parity

In SUSY models it is usually assumed that R-parity is a good symmetry. However, this is not necessarily the case, and one can construct SUSY models with broken R-parity. We concentrate on the MSSM without R-parity [71]. Extra trilinear terms are allowed in the superpotential, and some of them can give rise to a large enhancement of the \( B \to X_q \nu \bar{\nu} \) decay rate. Denoting by \( L^i_L, Q^i_L \) and \( d^i_R \) the chiral superfields containing respectively the left-handed lepton and quark doublets, and the right-handed down-type quark singlets of the \( i \)-th generation, these terms read

\[
W_R = \lambda'_{ijk} L^i_L Q^j_L d^k_R, \quad (5.22)
\]

where, for simplicity, we assume the \( \lambda'_{ijk} \) couplings to be real. Omitting terms involving the \( u_L \) and \( \ell_L \) fermions which are not relevant for the tree level \( b \to q \nu_i \bar{\nu}_j \) transition, the Yukawa interactions of the R-parity breaking Lagrangian generated by (5.22) are [72]

\[
\mathcal{L}_R = \lambda'_{ijk} \left[ \bar{d}_L^j \left( d^k_R \nu^i_L \right) + \bar{d}_R^k \left( \nu^c_i d^j_L \right) \right]. \quad (5.23)
\]
The exchange of $\tilde{d}_R$ and $\tilde{d}_L$ squarks gives rise to the effective four-fermion interaction (2.1) responsible for $B \rightarrow X_q \nu_i \bar{\nu}_j$, with the coefficients

$$C_{qij}^L = \sum_k \frac{\lambda'_{iqk} \lambda'_{j3k}}{2 m_{\tilde{d}_R}^2}, \quad C_{qij}^R = \sum_k \frac{\lambda'_{ikq} \lambda'_{jk3}}{2 m_{\tilde{d}_L}^2}, \quad (q = d, s).$$

(5.24)

We note that, in contrast to the LQ case, here both the $O_L$ and $O_R$ operators are necessarily present, due to the simultaneous appearance of the $\tilde{d}_L$ and $\tilde{d}_R$ scalar superpartners of the $d_{L,R}$ quarks. For simplicity, we neglect possible cancellations among the different $\tilde{d}^k$ ($k = 1, 2, 3$) exchange amplitudes. Then the bound (1.2) implies the following limits on the product of R-parity violating couplings

$$\lambda'_{iqk} \lambda'_{j3k} < 1.1 \times 10^{-3} \left( \frac{m_{\tilde{d}_R}}{100 \text{ GeV}} \right)^2, \quad \lambda'_{ikq} \lambda'_{jk3} < 1.1 \times 10^{-3} \left( \frac{m_{\tilde{d}_L}}{100 \text{ GeV}} \right)^2.$$

(5.25)

These limits represent the strongest constraints on the product of couplings involving the third generation, such as $\lambda'_{3q3} \lambda'_{333}$, $\lambda'_{33q} \lambda'_{333}$. As in the LQ models, these couplings are expected to be particularly large in models that relate their size (relative to the couplings involving the lighter generations) to the fermion mass hierarchy [73].

3. TopColor models

TopColor models [49] attempt to explain the large value of the top mass through the dynamical formation of a $t \bar{t}$ condensate. In these models, the basic assumption is that new dynamics strong enough to form chiral condensates is effective for the third generation, which therefore is treated differently from the first two. As a consequence, TopColor models have peculiar implications for the phenomenology involving the third generation, such as top and bottom production [74] and the $Z \rightarrow b \bar{b}$ decay rate [75]. In particular, they are expected to yield large effects in FCNC $B$ decays involving third generation leptons [76, 77].

Several models can be constructed along these lines, and we concentrate on the one studied in [76]. In this model the gauge symmetry breaking structure is

$$SU(3)_1 \times U(1)_1 \times SU(3)_2 \times U(1)_2 \rightarrow SU(3)_{\text{QCD}} \times U(1)_Y.$$  

(5.26)
Here $SU(3)_1 \times U(1)_1$ couples only to the third generation while $SU(3)_2 \times U(1)_2$ couples only to the first and second generations. The quantum numbers under these groups coincide with those under the usual $SU(3)_{QCD} \times U(1)_Y$. The breaking into the SM group \( (5.26) \) is induced by a $\langle t\bar{t} \rangle$ condensate, which is generated at the 1 TeV scale when the $SU(3)_1$ coupling becomes strong. The initial symmetry is larger than the electroweak gauge group, and this implies the existence of new massive gauge bosons corresponding to the additional broken generators: a color octet $B^a_\mu$ (topgluons) and a singlet $Z'_\mu$. We concentrate on the $Z'$ boson since it can mediate $B \rightarrow X_s \nu \bar{\nu}$ decays. The couplings of $Z'$ to the fermions are given by \[ \mathcal{L}_{\text{TopC}} = g_1 f_i(\theta') \left( \frac{1}{6} Q^i_L \gamma_\mu Q^i_L + \frac{2}{3} \bar{u}^i_R \gamma_\mu u^i_R - \frac{1}{3} \bar{u}^i_R \gamma_\mu d^i_R - \frac{1}{2} \bar{L}^i_L \gamma_\mu L^i_L - \bar{e}^i_R \gamma_\mu \ell^i_R \right) Z'^\mu, \] \( (5.27) \) with $Q_L = (u,d)_L$, $L_L = (\ell,\nu)_L$, $g_1 \simeq 0.35$ is the $U(1)_Y$ coupling constant, and $i = 1,2,3$ is a generation index.

The most important difference between this model and the usual $Z'$ models arises from the $f_i(\theta')$ factor, which enhances the strength of the third generation couplings with respect to the first and second generations. One has $f_{1,2}(\theta') = \tan \theta'$ for the first two generations and $f_3(\theta') = -\cot \theta'$ for the third generation. In general $\cot \theta' \gg 1$ is expected, in order to ensure that the condensate forms in the top direction \[ [49] \]. After integrating out the $Z'$ boson and rotating the $d_{L,R}$ quarks into the mass basis by the unitary matrices $U_{L,R}$, the interaction Lagrangian \( (5.27) \) gives rise to the following coefficients for the effective four-fermion interaction \( (2.1) \)

\[ C^q_{ii} = \frac{1}{12} \frac{g_1^2}{M^2_{Z'}} \kappa'^{bq} f_i \quad \Rightarrow \quad \tilde{C}^2_L = \left( \frac{\sin^2 \theta_W \ M^2_{Z'}}{6 \ M^2_{Z'}} \right)^2 \sum_{i=1,2,3 \atop q=d,s} f_i^2 \left| \kappa'^{bq}_L \right| ^2, \]

\[ C^q_{ii} = -\frac{1}{6} \frac{g_1^2}{M^2_{Z'}} \kappa'^{bq} f_i \quad \Rightarrow \quad \tilde{C}^2_R = \left( \frac{\sin^2 \theta_W \ M^2_{Z'}}{3 \ M^2_{Z'}} \right)^2 \sum_{i=1,2,3 \atop q=d,s} f_i^2 \left| \kappa'^{bq}_R \right| ^2. \] \( (5.28) \)

Here $\kappa'^{bq} = \sum_j (U^*_{bj} \ U_{jq}) f_j \approx -(U^*_{k3} \ U_{3q}) \cot \theta'$ gives the flavor changing $Z'$ couplings to the quarks (the $L$ and $R$ chirality labels for $\kappa'^{bq}_{L,R}$ and for $U_{L,R}$ are understood).
A too large value of \( \cot \theta' \) would lead to a spontaneous breaking of the chiral symmetry for the tau lepton. This implies the bound \([77,78]\)

\[
\cot^2 \theta' < \frac{8\pi^2}{g_1^2} \quad \Rightarrow \quad |\cot \theta'| \lesssim 25.
\] (5.29)

Inserting the values of the various hypercharges given in (5.27) into Eq. (5.6), we find

\[
\frac{\text{BR}(B \to X_s \nu \bar{\nu})}{\text{BR}(B \to X_s \mu^+ \mu^-)} = \frac{f_1^2 + f_2^2 + f_3^2}{5f_2^2} \approx \frac{\cot^4 \theta'}{5},
\] (5.30)

where we assumed that NP effects dominate the decay rates, and thus we neglected the SM contribution. As is apparent from (5.30), due to their different dependence on \( \theta' \), a measurement of both \( B \to X_s \mu^+ \mu^- \) and \( B \to X_s \nu \bar{\nu} \) would allow us to separately determine the value of this parameter. The current limit on \( B \to X_s \nu \bar{\nu} \) in this model is obtained by combining Eqs. (5.29) and (5.30) with the experimental upper bound on \( B \to X_s \mu^+ \mu^- \) (4.3).

The bounds (4.3) gives the new stringent limits

\[
|\cot \theta' \kappa_{bq}^L| < 5.4 \left( \frac{M_{Z'}}{1 \text{ TeV}} \right)^2, \quad |\cot \theta' \kappa_{bq}^R| < 2.7 \left( \frac{M_{Z'}}{1 \text{ TeV}} \right)^2.
\] (5.31)

4. **Horizontal gauge symmetries**

Attempts to explain the hierarchical pattern of fermion masses and mixings by some underlying dynamical interaction, are often based on broken horizontal gauge symmetries \([79,80,55]\). In the non-Abelian case, the fermions of different generations are assigned to some irreducible representation of the horizontal gauge group. Flavor changing transitions occur, suppressed by the masses of the heavy horizontal gauge bosons.

The fermion mass pattern is generated through the so-called universal see-saw mechanism \([51]\), in which the fermion masses are suppressed from their natural scale \( G_F^{-1/2} \), to the observed values by inverse powers of some large mass scale, with no need to fine tune the Yukawa couplings \([79,80]\). Various mass scales are associated with different stages of the horizontal gauge symmetry breaking in such a way that the heavier fermions of the third
generation couple to the lightest horizontal gauge bosons. As a consequence, rare FCNC transitions of the third family are “naturally” enhanced.

To illustrate the main features of these models, we consider as a simple example a model based on the horizontal gauge symmetry group $SU(3)_H \times U(1)_H$ [79]. The standard $q$ quarks are assigned to the $(3,1)_H$ representation of $SU(3)_H \times U(1)_H$, while new isosinglet $Q$ quarks are in $(\bar{3},-1)_H$. The breaking of the horizontal symmetry is achieved through the non-vanishing VEVs $\xi_\alpha$, of a set of SM singlet horizontal scalars belonging to the $(\bar{3},2)_H$ or $(6,2)_H$, which also give large masses to the heavy isosinglet quarks. The mass terms for standard quarks are provided by the VEV of the SM Higgs, $v = (2\sqrt{2}G_F)^{-1/2} \approx 175$ GeV, and by an additional VEV $\eta$ of a real scalar isosinglet belonging to $(1,0)_H$. Then, for the fermion mass terms we have

$$L_M \sim \bar{Q}_L^i M(\xi)_{ij} Q_R^j v + \bar{Q}_L^i \bar{q}_R^i \eta,$$

where $i,j = 1,2,3$ are flavor indices. We note that due to the assumed horizontal gauge symmetry, the second and third terms in (5.32) involving the light fermions are the same for the three generations.

By assumption $v, \eta \ll \xi_\alpha$, and thus the quark masses are generated via a see-saw like mechanism, yielding for the mass matrices $m_q \approx v \eta M(\xi)^{-1}$. All information on the quark mass hierarchies and mixings are contained in $M(\xi)$, and depend on the hierarchical structure of the VEVs $\xi_\alpha$. It is assumed that $\xi_1 \gg \xi_2 \gg \xi_3$ in order to reproduce the generation hierarchy pattern. For example, together with additional smaller induced VEVs contributing to $M$, they can be chosen to form a Fritzsch structure [82], yielding

$$\xi_3 : \xi_2 : \xi_1 \sim 1 : \sqrt{m_t/m_u} : \sqrt{m_c/m_u} \sim 1 : 190 : 3300.$$  

(5.33)

The horizontal VEVs hierarchy induces the breaking chain

$$SU(3)_H \times U(1)_H \xrightarrow{\xi_1} SU(2)_H \times U(1)'_H \xrightarrow{\xi_2} U(1)''_H \xrightarrow{\xi_3} I.$$  

(5.34)

In the first stage, four flavor changing gauge bosons carrying family “charge” plus a combination of the three neutral generators acquire a large mass of order $\xi_1$, leaving a residual
SU(2)\textsubscript{H} × U(1)\textsubscript{H} symmetry unbroken, which acts only on the second and third generations. The four “charged” bosons couple the first family fermions directly (namely, not through small mixing angles) to the fermions of the second and third generations, and in particular they induce \( \Delta S = 2 \) effective operators. The requirement that these operators will not generate unacceptably large contributions to the \( \bar{K}^0 - K^0 \) mass difference implies for the scale of the first breaking \[ \xi_1 \gtrsim 3 \times 10^{3} \text{ TeV}. \] (5.35)

In the second stage, two charged bosons and a second neutral combination acquire a mass of order \( \xi_2 \). These bosons give direct contributions to the transition \( B \rightarrow X_s \nu_\tau \bar{\nu}_\mu \), and as a consequence our limit (1.2) implies a lower bound on \( \xi_2 \) (see Eq. (5.38) below). The remaining local symmetry \( U(1)\textsubscript{H}'' \) that acts only on the third generation is broken at a scale \( \xi_3 \) much smaller than \( \xi_1 \) and \( \xi_2 \). For example, if we assume the hierarchy (5.33), the bound (5.33) implies the rather weak constraint \( \xi_3 > 0.9 \) TeV.

We see that this scenario has the natural implication of a new \( X^0_\mu \) boson corresponding to \( U(1)\textsubscript{H}'' \) in (5.34), coupled only to the third generation, while due to the hierarchy (5.33) the effects of the heavier horizontal bosons are much more suppressed. Due to the mixing between the quark mass eigenstates, \( X^0_\mu \) will also mediate flavor changing transitions. The neutral current interaction Lagrangian describing these transitions is analogous to (5.2), and the terms relevant to the process we are interested in are

\[
\mathcal{L}_H = g_H \left[ \kappa_{L}^{bg} \bar{b}_L \gamma^\mu q_L + \kappa_{R}^{bg} \bar{b}_R \gamma^\mu q_R + \bar{\nu}_{\tau L} \gamma^\mu \nu_{\tau R} \right] X^0_\mu, \tag{5.36}
\]

where \( \kappa_{L,R} = U_{L,R} \dagger P_3 U_{L,R} \) with \( P_3 = \text{diag}(0, 0, 1) \), and the matrices \( U_{L,R} \) rotate the \( d_{L,R}-\text{type} \) quarks mass eigenstates. Eq. (5.36) yields the coefficients

\[
C_{L,R} = \frac{1}{\xi_3^4} \kappa_{L,R}^{bg} \quad \Rightarrow \quad C^2_{L,R} = \frac{\nu^4}{\xi_3^4} \sum_{q=d,s} |\kappa_{L,R}^{bg}|^2. \tag{5.37}
\]

If the FCNC mixings \( \kappa_{L,R}^{bg} \) are very small, then in spite of the larger mass suppression the direct transition \( B \rightarrow X_s \nu_\tau \bar{\nu}_\mu \) mediated by the “charged” \( SU(2)\textsubscript{H} \) bosons of mass

34
\[ M_{X^\pm} = g_H \xi_2 \text{ (see (5.34))} \] can give contributions to \( B \to X_s \nu \bar{\nu} \) which are competitive to those mediated by \( X^0_\mu \). The rate for the \( X^\pm_\mu \) mediated decay \( B \to X_s \nu_\tau \bar{\nu}_\mu \) becomes as large as the rate for \( B \to X_s \nu_\tau \bar{\nu}_\tau \) mediated by \( X^0_\mu \) when \( |\kappa_{L,R}^{bq}| \sim \xi_3^2/\xi_2^2 \). Due to the different final states, the two amplitudes do not interfere, and the coefficients \( C_{L,R} \) for the second case follow easily from (5.37) by the substitution \( \kappa_{L,R}^{bq}/\xi_2^3 \to 1/\xi_2^2 \). Our limit (1.2) implies the two bounds

\[ |\kappa_{L,R}^{bq}| < 6 \times 10^{-4} \left( \frac{\xi_3}{100 \text{ GeV}} \right)^2, \quad \xi_2 > 6 \text{ TeV}. \] (5.38)

A different scenario which also predicts a new \( U(1) \) bosons coupled only to the third family fermions was presented in [50]. Similar to TopColor models, in this scenario the large value of the top mass is explained by a dynamical scheme which implies a large isospin breaking for the \( t \) and \( b \) masses. However, that does not feed back directly into the \( W \) and \( Z \) masses. The \( U(1) \) gauge boson coupled to the third generation is the remnant of the breaking of some large non-Abelian semi-simple gauge group which embeds the HyperColor interaction responsible for the formation of the isospin breaking condensate. The phenomenology of such a gauge boson, including the consequences on LEP physics due to its mixing with the \( Z \), have been extensively studied [83]. Recently it was also suggested that the \( B \to X_s \nu \bar{\nu} \) decay rate could be largely enhanced because of the new FCNC contributions, and could even approach the rate of semileptonic \( B \) decay [84]. Such a mechanism for increasing the \( B \) width was proposed as a possible solution to the claimed discrepancy between the observed experimental value of the semileptonic \( B \) branching ratio and the theoretical predictions. The limit (1.2) obviously rules out this possibility, as well as any other mechanism attempting to achieve a sizable increase of the \( B \) width through an enhancement of decay modes associated with large missing energy.

The interaction Lagrangian for the \( X_\mu \) boson of this model, written in the fermion mass basis, is identical to the Lagrangian in Eq. (5.36). However, there is now one additional constraint for the \( X_\mu \) mass [84]

\[ \frac{g_H^2}{M_X^2} = \frac{G_F}{2\sqrt{2}}. \] (5.39)
Using $M_X = g_R \xi_3$, it is then straightforward to derive from (5.37) the expression for the $\tilde{C}_{L,R}$ coefficients

$$\tilde{C}_{L,R}^2 = \frac{1}{64} \sum_{q=d,s} |\kappa_{L,R}^{bq}|^2.$$  (5.40)

The NP parameters $\kappa_{L,R}^{bq}$ are denoted by $\lambda_{l3}^L$ and $-\lambda_{l3}^R (i = 1, 2)$ in Ref. [84]. It was speculated that $|\kappa_{L}^{bs}|^2 + |\kappa_{R}^{bs}|^2 \approx 30 |V_{cb}|^2$ could account for the possible discrepancy in the $B$ semileptonic branching ratio [84]. Our limit (1.2) implies a bound two orders of magnitude smaller, $|\kappa_{L}^{bs}|^2 + |\kappa_{R}^{bs}|^2 < 2 \times 10^{-4}$. For $q = d$, the limit $|\kappa_{L}^{bd} - \kappa_{R}^{bd}| < 2 \times 10^{-3}$ can be derived from $B_d - \bar{B}_d$ mixing [84]. However, there is no similarly strong limit on the individual $|\kappa_{L,R}^{bd}|$ parameters. From (1.2) we obtain

$$|\kappa_{L,R}^{bq}| < 1.4 \times 10^{-2} \quad (q = d, s).$$  (5.41)

**VI. SUMMARY AND CONCLUSIONS**

In this paper we discussed the derivation of the first bound on the inclusive $B \to X_s \nu \bar{\nu}$ decay rate, using the large missing energy tag in $B$ decays at LEP. We studied in detail the theoretical ingredients needed to carry out such an analysis, and we found that the overall theoretical uncertainty is small. Therefore, this decay mode is well suited to search for physics beyond the SM. We translated the ALEPH bound on the $B \to \tau \bar{\nu}$ branching ratio, which resulted from a search for $B$ decays with large missing energy, into a limit on the $B \to X_s \nu \bar{\nu}$ branching ratio. To derive a numerical limit, we had to make a number of conservative and simplifying assumptions. Thus, the resulting bound is weaker than what a dedicated experimental analysis will be able to achieve. Our conservative upper bound is

$$\text{BR}(B \to X_s \nu \bar{\nu}) < 3.9 \times 10^{-4},$$  (6.1)

which is less than one order of magnitude above the SM prediction. We estimated that using the full LEP–I data sample, the LEP collaborations may be able to set a limit of order
(1 − 2) \times 10^{-4}. Due to the theoretical interest in the $B \to X_s \nu \bar{\nu}$ decay mode, we think it is important that the LEP collaborations will perform a dedicated analysis of this process.

We studied a variety of new physics models. After discussing the constraints from existing experimental data, we divided the NP models into three classes, according to increasing allowed values of the the $B \to X_s \nu \bar{\nu}$ branching ratio.

To class (A) belong those models in which the $B \to X_s \nu \bar{\nu}$ branching ratio is already constrained to be below the expected sensitivity of the LEP experiments.

Class (B) contains the models which allow for an enhancement of BR($B \to X_s \nu \bar{\nu}$) up to values that will be observable at LEP. These models are listed in Table II, together with the maximal $B \to X_s \nu \bar{\nu}$ branching ratio they allow for, after the existing constraints are taken into account. These models naturally evade the constraints imposed by $B \to X_s \gamma$. The limits on the relevant NP parameters which we obtain from the bound on the $B \to X_s \nu \bar{\nu}$ decay rate, are numerically close to the limits provided by the bounds on the inclusive and exclusive $B \to X_s \ell^+ \ell^-$ decays. However, our bounds are more reliable, as the $B \to X_s \nu \bar{\nu}$ decay is theoretically cleaner. We expect that in the near future, new results emerging from CLEO, CDF and (hopefully) from LEP can compete in further constraining these models.

The most interesting models for our investigation belong to class (C). For the models in this class, the bound on the $B \to X_s \nu \bar{\nu}$ branching ratio implies new constraints that are not matched by other existing experimental data. A generic feature of these models is that they yield a natural enhancement of the FCNC processes involving third generation fermions, without conflicting with other constraints. We derived new limits on the couplings of the third generation fermions in models with leptoquarks, and in supersymmetric models with broken R-parity. Our bounds also imply stringent constraints on models in which new gauge bosons are coupled dominantly to the third generation, such as TopColor models, models based on non-Abelian horizontal gauge symmetries, and other models attempting to explain dynamically the large value of the top mass. The new bounds on the parameters of these models, implied by the limit on $B \to X_s \nu \bar{\nu}$, are summarized in Table II.

In future $B$ factories much larger data samples will become available. While – to our
knowledge – no detailed study concerning the possibility of measuring the $B \to X_s \nu \bar{\nu}$ decay rate at $B$ factories has been carried out, we hope that it will be possible to perform such a search. A precise measurement of the $B \to X_s \nu \bar{\nu}$ decay rate would provide a very reliable means of directly determining the $|V_{ts}|$ element of the CKM matrix. This would allow for an important new test of the unitarity of the CKM matrix. If deviations from the SM predictions will be detected in rare $B$ decays, it will also be important to measure as many decay modes as possible. The pattern of deviations from the SM predictions in different decay rates would help us to distinguish between various possible NP scenarios.

ACKNOWLEDGMENTS

We are grateful to Ian Tomalin of ALEPH for very helpful correspondence about the experimental analysis. We also thank Yossi Nir and Mark Wise for discussions and comments on the manuscript. ZL was supported in part by the U.S. Dept. of Energy under Grant no. DE-FG03-92-ER 40701.
REFERENCES

[1] R. Ammar et al., CLEO Collaboration, Phys. Rev. Lett. 71 (1993) 674.
[2] M.S. Alam et al., CLEO Collaboration, Phys. Rev. Lett. 74 (1995) 2885.
[3] C. Albajar et al., UA1 Collaboration, Phys. Lett. B 262 (1991) 163.
[4] R. Balest et al., CLEO Collaboration, CLEO-CONF-94-4. Submitted to Int. Conf. on High Energy Physics, Glasgow, Scotland, Jul 20-27, 1994.
[5] C. Anway-Wiese et al., CDF Collaboration, Fermilab-Conf-95/201-E. LP’95: International Symposium on Lepton-Photon Interactions (IHEP).
[6] D. Buskulic et al., ALEPH Collaboration, Phys. Lett. B 298 (1993) 479.
[7] D. Buskulic et al., ALEPH Collaboration, Phys. Lett. B 343 (1995) 444.
[8] M. Acciarri et al., L3 Collaboration, Phys. Lett. B 332 (1994) 201.
[9] F. Behner, talk given at the EPS–HEP Conference, Brussels (1995).
[10] A.F. Falk, Z. Ligeti, M. Neubert, and Y. Nir, Phys. Lett. B 326 (1994) 145.
[11] G. Buchalla and A.J. Buras, Nucl. Phys. B 400 (1993) 225.
[12] See e.g., R.E. Marshak, Riazuddin and C.P. Ryan, Theory of weak interactions in particle physics (John Wiley & Sons, Inc., 1969).
[13] T. Inami and C.S. Lim, Prog. Theor. Phys. 65 (1981) 297; (E) ibid. 1772.
[14] A.F. Falk, M. Luke, and M.J. Savage, Phys. Rev. D 53 (1996) 2491.
[15] Y. Grossman and Z. Ligeti, Phys. Lett. B 332 (1994) 373.
[16] J. Chay, H. Georgi, and B. Grinstein, Phys. Lett. B 247 (1990) 399.
[17] A.V. Manohar and M.B. Wise, Phys. Rev. D 49 (1994) 1310.
[18] B. Blok, L. Koyrakh, M. Shifman and A.I. Vainshtein, Phys. Rev. D 49 (1995) 3356.
[19] Particle Data Group, Phys. Rev. D 50 Part. I (1994).
[20] A. Czarnecki, M. Ježabek and J.H. Kühr, Act. Phys. Pol. B20 (1989) 961.
[21] C. Peterson et al., Phys. Rev. D 27 (1983) 105; D. Buskulic et al., ALEPH Collaboration, Z. Phys. C 62 (1994) 179.
[22] D. Buskulic et al., ALEPH Collaboration, Phys. Lett. B 322 (1994) 441.
[23] I.I. Bigi, UND-HEP-95-BIG02 [hep-ph/9508408].
[24] D. Buskulic et al., ALEPH Collaboration, Phys. Lett. B 357 (1995) 685.
[25] B. Grinstein, M.J. Savage and M.B. Wise, Nucl. Phys. B 319 (1989) 271.
[26] A.J. Buras and M. Münz, Phys. Rev. D 52 (1995) 186.
[27] N.G. Deshpande and J. Trampetic, Phys. Rev. Lett. 60 (1988) 2583; C.A. Domingues, N. Paver and Riazuddin, Z. Phys. C 48 (1990) 55; W. Jaus and D. Wyler, Phys. Rev. D 41 (1990) 3405; A. Ali and T. Mannel Phys. Lett. B 264 (1991) 447, (E) ibid. 274 (1992) 526.
[28] N. Isgur and M.B. Wise, Phys. Rev. D 42 (1990) 2388; for a more recent discussion see Ref. 76.
[29] P. Ko, Phys. Rev. D 45 (1992) 174; G. Buchalla and A.J. Buras, Nucl. Phys. B 412 (1994) 106.
[30] H.E. Haber and G.L. Kane, Phys. Rep. 117 (1985) 75; H.P. Nilles, Phys. Rep. 110 (1984) 1; M.F. Sohnius, Phys. Rep. 128 (1985) 41.
[31] J.F. Gunion, H.E. Haber, G.L. Kane and S. Dawson, The Higgs Hunter’s Guide (Addison-Wesley Publishing Co., Reading MA, 1990); and references therein.
[32] S. Bertolini, F. Borzumati, A. Masiero and G. Ridolfi, Nucl. Phys. B 353 (1991) 591.
[33] A. Ali, G.F. Giudice and T. Mannel, Z. Phys. C 67 (1995) 417.
[34] Y. Grossman, Nucl. Phys. B 426 (1994) 355.
[35] Y. Grossman, H.E. Haber and Y. Nir, Phys. Lett. B 357 (1995) 630.
[36] P. Langacker and S.U. Sankar, Phys. Rev. D 40 (1989) 1569, and references therein; M. Leurer, Ph.D. Thesis, Weizmann Institute (1985).
[37] P. Cho and M. Misiak, Phys. Rev. D 49 (1994) 5894.
[38] See, for example, P. Langacker and D. London, Phys. Rev. D 38 (1988) 886; ibid. 38 (1988) 907; E. Nardi, E. Roulet and D. Tommasini, Nucl. Phys. B 386 (1992) 239.
[39] Y. Nir and D.J. Silverman, Phys. Rev. D 42 (1990) 1477; D.J. Silverman, Phys. Rev. D 45 (1992) 1800.
[40] V. Barger, M.S. Berger and R.J.N. Phillips, Phys. Rev. D 52 (1995) 1663.
[41] E. Nardi, E. Roulet and D. Tommasini, Phys. Lett. B 327 (1994) 319; ibid. 344 (1995) 225.
[42] E. Nardi, Phys. Rev. D 48 (1993) 1240; J. Bernabéu, E. Nardi and D. Tommasini, Nucl. Phys. B 409 (1993) 69.
[43] E. Nardi, Phys. Rev. D 48 (1993) 3277; ibid. 49 (1994) 4394; E. Nardi and T.G. Rizzo, Phys. Rev. D 50 (1994) 203.
[44] T.P. Cheng and M. Sher, Phys. Rev. D 35 (1987) 3484; M. Sher and Y. Yuan, ibid. 44 (1991) 1461; A. Antaramian, L.J. Hall and A. Rasin, Phys. Rev. Lett. 69 (1992) 1871; M. Luke and M. Savage, Phys. Lett. B 307 (1993) 387.
[45] L.J. Hall and S. Weinberg, Phys. Rev. D 48 (1993) 979.
[46] T. Han, R.D. Peccei and X. Zhang, Nucl. Phys. B 454 (1995) 527; R.D. Peccei, S. Peris and X. Zhang, Nucl. Phys. B 349 (1991) 305.
[47] D. Atwood, L. Reina and A. Soni, Phys. Rev. D 53 (1996) 1199; Phys. Rev. Lett. 75 (1995) 3800.
[48] E. Nardi, Phys. Lett. B 365 (1996) 327.
[49] C.T. Hill, Phys. Lett. B 345 (1995) 483.
[50] B. Holdom, Phys. Lett. B 336 (1994) 85.
[51] B. Grzadkowski, J.F. Gunion and P. Krawczyk, Phys. Lett. B 268 (1991) 106; G. Eilam, J.L. Hewett and A. Soni, Phys. Rev. D 44 (1991) 1473; M. Luke and M.J. Savage, Phys. Lett. B 307 (1993) 387.
[52] W. Buchmüller and M. Gronau, Phys. Lett. B 220 (1989) 641.
[53] G.C. Branco, P.A. Parada and M.N. Rebelo, Phys. Rev. D 52 (1995) 4217.
[54] N. Evans, Phys. Lett. B 340 (1994) 81.
[55] Z. Berezhiani and E. Nardi, Phys. Lett. B 355 (1995) 199; Phys. Rev. D 52 (1995) 3087.
[56] J.F. Gunion, D.W. McKay and H. Pois, Phys. Lett. B 334 (1994) 339; Phys. Rev. D 53 (1996) 1616.
[57] J.L. Hewett, Phys. Lett. B 193 (1987) 327; W.S. Hou, A. Soni and H. Steger, Phys. Lett. B 192 (1987) 441.
[58] J.L. Hewett, SLAC-PUB-6521, Presented at SLAC Summer Inst. on Particle Physics, SLAC, Jul 6 – Aug 6, 1993 [hep-ph/9406302], and references therein.
[59] See e.g.: A. De Rujula, M.B. Gavela, P. Hernandez, and E. Masso, Nucl. Phys. B 384 (1992) 3; C. Arzt, M.B. Einhorn, and J. Wudka, Nucl. Phys. B 433 (1995) 41, and references therein.
[60] J.A. Grifols, S. Peris and J. Solà, Int. J. Mod. Phys. A 3 (1988) 225; X.-G. He and
B.H.J. McKeller, Phys. Rev. D 51 (1995) 6484.
[61] X.-G. He, Phys. Lett. B 319 (1993) 327;
[62] G. Baillie, Z. Phys. C 61 (1994) 667.
[63] S. Dawson and G. Valencia, Phys. Rev. D 49 (1994) 2188.
[64] K. Hagiwara, R.D. Peccei, D. Zeppenfeld and K. Hikasa, Nucl. Phys. B 282 (1987) 253.
[65] CDF Collaboration, F. Abe et al., Phys. Rev. Lett. 75 (1995) 1017.
[66] L. Randall and R. Sundrum, Phys. Lett. B 312 (1993) 148, and references therein.
[67] B. Grinstein, Y. Nir and J.M. Soares, Phys. Rev. D 48 (1993) 3960.
[68] S. Davidson, D. Bailey and B.A. Campbell, Z. Phys. C 61 (1994) 613.
[69] M. Leurer, Phys. Rev. D 49 (1994) 333; ibid. 50 (1994) 536.
[70] M. Leurer, Y. Nir and N. Seiberg, Nucl. Phys. B 398 (1993) 319; ibid. 420 (1994) 468;
Y. Grossman and Y. Nir, Nucl. Phys. B 448 (1995) 30.
[71] C.S. Aulakh and R.N. Mohapatra, Phys. Lett. B 119 (1982) 136; L.J. Hall and M.
Suzuki, Nucl. Phys. B 231 (1984) 419; S. Dawson, Nucl. Phys. B 261 (1985) 297; G.
Ross and J. Valle, Phys. Lett. B 151 (1985) 375; J. Ellis et al., Phys. Lett. B 150
(1985) 142.
[72] V. Barger, G.F. Giudice and T. Han, Phys. Rev. D 40 (1989) 2987.
[73] T. Banks, Y. Grossman, E. Nardi and Y. Nir, Phys. Rev. D 52 (1995) 5319.
[74] C.T. Hill and S.J. Parke, Phys. Rev. D 49 (1994) 4454.
[75] C.T. Hill and X. Zhang, Phys. Rev. D 51 (1995) 3563.
[76] G. Burdman, Phys. Rev. D 52 (1995) 6400.
[77] G. Buchalla, G. Burdman, C.T. Hill and D. Kominis, Fermilab-Pub-95/322-T [hep-
ph/9510376]; to be published in Phys. Rev. D.
[78] D. Kominis, Phys. Lett. B 358 (1995) 312.
[79] Z.G. Berezhiani and Dzh.L. Chkareuli, Pis'ma Zh. Eksp. Teor. Fiz. 41, (1985) 473 [JETP
Lett. 41 (1985) 577]; Yad. Fiz. 52 (1990) 601, [Sov. J. Nucl Phys. 52 (1990) 383].
[80] J.L. Chkareuli, Pis'ma ZhETF 32 (1980) 684 [JETP Lett. 32 (1980) 671]; Z.G. Berezh-
iani and J.L. Chkareuli, Pis'ma ZhETF 35 (1982) 494 [JETP Lett. 35 (1982) 612]; Yad.
Fiz. 37 (1983) 1043 [Sov. J. Nucl. Phys. 37 (1983) 618].
[81] Z.G. Berezhiani, Phys. Lett. B 129 (1983) 99; D. Chang and R.N. Mohapatra, Phys.
Rev. Lett. 58 (1987) 1600; S. Rajpoot, Phys. Lett. B 191 (1987) 122.
[82] H. Fritzsch, Nucl. Phys. B 155 (1979) 189.
[83] B. Holdom, Phys. Lett. B 339 (1994) 114; ibid. 351 (1995) 279; X. Zhang and B.-L.
Young, Phys. Rev. D 51 (1995) 6584.
[84] B. Holdom and M.V. Ramana, Phys. Lett. B 365 (1996) 309.
TABLES

| Model                        | Allowed branching ratio |
|------------------------------|-------------------------|
| SM                           | $0.5 \times 10^{-4}$    |
| FCNC $Z$                     | $3.5 \times 10^{-4}$    |
| FCNC $Z'$                    | $1.8 \times 10^{-4}$    |
| Unconventional $E_6$         | $2.9 \times 10^{-4}$    |
| Anomalous $tcZ$ vertex       | $3.5 \times 10^{-4}$    |
| Fourth generation            | $3.7 \times 10^{-4}$    |
| Anomalous $WWZ$              | $3.5 \times 10^{-4}$    |
| ETC                          | $3.5 \times 10^{-4}$    |

TABLE I. Summary of the models belonging to class (B), which allow an enhancement of the $B \to X_s \nu \bar{\nu}$ decay rate up to a level observable at LEP. The standard model (SM) prediction for BR($B \to X_s \nu \bar{\nu}$) is given for comparison as the first entry. For each model listed in the first column, the second column gives the maximal $B \to X_s \nu \bar{\nu}$ branching ratio allowed by the existing experimental data.
| Model                                           | New bounds                                                                 |
|------------------------------------------------|----------------------------------------------------------------------------|
| LQ: $S_0, \tilde{S}_{1/2}, S_1$                 | $\lambda_{iq} \lambda_{j3} < 1.1 \times 10^{-3} \left( \frac{m_{LQ}}{100 \text{ GeV}} \right)^2$ |
| $V_{1/2}, V_1$                                  | $\lambda_{iq} \lambda_{j3} < 5.7 \times 10^{-4} \left( \frac{m_{LQ}}{100 \text{ GeV}} \right)^2$ |
| SUSY without R-parity                           | $\lambda'_{iqk} \lambda'_{j3k} < 1.1 \times 10^{-3} \left( \frac{m_{\tilde{d}_R}}{100 \text{ GeV}} \right)^2$ |
| $\lambda'_{ikq} \lambda'_{jk3} < 1.1 \times 10^{-3} \left( \frac{m_{\tilde{d}_L}}{100 \text{ GeV}} \right)^2$ |
| TopColor                                        | $|\cot \theta' \kappa'_{L}^{bq}| < 5.4 \left( \frac{M_{Z'}}{1 \text{ TeV}} \right)^2$ |
| $|\cot \theta' \kappa'_{R}^{bq}| < 2.7 \left( \frac{M_{Z'}}{1 \text{ TeV}} \right)^2$ |
| Horizontal gauge symmetry                       | $|\kappa_{L,R}^{bq}| < 6 \times 10^{-4} \left( \frac{\xi_3}{100 \text{ GeV}} \right)^2$ |
| $\xi_2 > 6 \text{ TeV}$                        | $|\kappa_{L,R}^{bq}| < 1.4 \times 10^{-2}$ |

**TABLE II.** Summary of the models belonging to class (C), for which essentially no constraints on the allowed $B \to X_s \nu \bar{\nu}$ branching fraction existed to date. For each model listed in the first column, the second column gives the bounds on the relevant model parameters implied by the new limit $\text{BR}(B \to X_s \nu \bar{\nu}) < 3.9 \times 10^{-4}$. The indices $i, j, k = e, \mu, \tau$ and $q = d, s$ correspond respectively to the neutrino and quark flavors in the final state.