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Rare $B \to K^{(*)}\nu\overline{\nu}$ decays at B factories

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Abstract

We compute the branching fraction of the decays $B \to K\nu\overline{\nu}$ and $B \to K^*\nu\overline{\nu}$ in the Standard Model. We also comment on the experimental difficulties and procedures to detect such modes at B factories. © 1997 Published by Elsevier Science B.V.

Among the various flavour changing neutral current-induced $b-$quark decays [1], the transition

$$b \to s\nu\overline{\nu} \tag{1}$$

plays a peculiar role, both from a theoretical and an experimental point of view.

Within the Standard Model (SM) the process (1) is governed by the effective hamiltonian

$$\mathcal{H}_{\rm eff} = \frac{G_F}{\sqrt{2}}\frac{\alpha}{2\pi\sin^2(\theta_W)}\, V_{ts}V_{tb}^*\, X(x_t)\, \bar{b}\gamma^\mu \times (1-\gamma_5)s\, \bar{\nu}\gamma_\mu(1-\gamma_5)\nu \equiv c_L^{\rm SM}\, \mathcal{O}_L \tag{2}$$

obtained from Z^0 penguin and box diagrams where the dominant contribution corresponds to a top quark intermediate state. In (2) G_F is the Fermi constant, α the fine structure coupling constant (at the Z^0 scale), θ_W the Weinberg angle and V_{ij} are Cabibbo-Kobayashi-Maskawa (CKM) matrix elements; $x_t = (m_{\rm top}/M_W)^2$. $\mathcal{O}_L$ represents the left-left four-fermion operator $\mathcal{O}_L \equiv \bar{b}\gamma^\mu(1-\gamma_5)s\, \bar{\nu}\gamma_\mu(1-\gamma_5)\nu$. The $\mathcal{O}(\alpha_s)$ contribution deriving from two-loop diagrams is taken into account in the function X:

$$X(x) = X_0(x) + \frac{\alpha_s}{4\pi}X_1(x)\ , \tag{3}$$

where [2]

$$X_0(x) = \frac{x}{8}\left[\frac{x+2}{x-1} + \frac{3x-6}{(x-1)^2}\ln x\right] \tag{4}$$

and [3,4]

$$X_1(x) = \frac{4x^3-5x^2-23x}{3(x-1)^2} - \frac{x^4+x^3-11x^2+x}{(x-1)^3}\ln x + \frac{x^4-x^3-4x^2-8x}{2(x-1)^3}\ln^2 x + \frac{x^3-4x}{(x-1)^2}L_2(1-x) + 8x\frac{\partial X_0(x)}{\partial x}\ln x_\mu \tag{5}$$

($L_2(1-x) = \int_1^x dt\, \ln t/(1-t)$, and $x_\mu = \mu^2/M_W^2$, with $\mu = \mathcal{O}(m_{\rm top})$). Such a correction, using $m_{\rm top} = 175 \pm 9$ GeV [5] and $\alpha_s(m_b) = 0.23$, is around 3 %.

The presence of a single operator governing the transition (1) is a welcomed property, since the theoretical uncertainty is only related, within SM, to the value of one Wilson coefficient $c_L^{\rm SM}$. In other cases, for

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example in $b \to s\ell^+\ell^-$, the effective hamiltonian consists of several terms coherently acting to determine branching ratios, invariant mass distributions, lepton charge and polarization asymmetries, etc., and the uncertainty of a set of coefficients appearing in interfering terms must be taken into account. Moreover, possible New Physics (NP) effects contributing to (1) can only modify the SM value of the coefficient c_L, or introduce one new right-right operator

$$\mathcal{H}_{\rm eff} \equiv c_L\, \mathcal{O}_L + c_R\, \mathcal{O}_R \qquad (6)$$

($\mathcal{O}_R \equiv \bar{b}\gamma^\mu(1+\gamma_5)s\; \bar{\nu}\gamma_\mu(1+\gamma_5)\nu$), with c_R only receiving contribution from phenomena beyond SM.

The process (1) is theoretically appealing also because of the absence of long-distance contributions, which are usually related to the presence of four-quark operators in the effective hamiltonian and, e.g., heavily affect the process $b \to s\ell^+\ell^-$ [6]. In this respect, the transition to neutrinos represents a clean process even in comparison with the $b \to s\gamma$ decay, where long-distance contributions are expected to be present, although small [7].

As for inclusive decays, the analysis of $B \to X_s\nu\bar{\nu}$ in the framework of the expansion in the inverse heavy quark mass shows that the $\mathcal{O}(m_b^{-2})$ preasymptotic corrections to the partonic spectrum are negligible for all values of the squared momentum transferred to the neutrino pair, but for a narrow region near the endpoint [8].

Moreover, the exclusive $B \to K^{(*)}$ transitions induced by (2) can be related to the semileptonic Cabibbo suppressed $B \to \pi(\rho)\ell\nu$ decays, on which first results are now available [9]. The idea is to set up a procedure for a determination of the ratio of the CKM matrix elements appearing in $b \to s$ and $b \to u$ transition $|V_{ts}|/|V_{ub}|$ in a way safe of hadronic uncertainties [10,11].

From the experimental point of view, the analysis of the B meson decays induced by (2) has to be considered together with the study of the purely leptonic decay mode $B^- \to \tau^-\bar{\nu}_\tau$, where two neutrinos are also produced in the final state. This analogy has been exploited [12,13] to establish a bound on the inclusive $B \to X_s\nu\bar{\nu}$ branching ratio using the upper limit obtained by the ALEPH Collaboration at CERN [13,14]:

$$\mathcal{B}(B^- \to \tau^-\bar{\nu}_\tau) \le 1.6 \times 10^{-3} \quad (\text{at } 90\%\ \text{CL}) \quad . \qquad (7)$$

The resulting bound [13]

$$\mathcal{B}(B \to X_s \sum_i \nu_i\bar{\nu}_i) \le 7.7 \times 10^{-4} \qquad (8)$$

must be compared to the SM prediction, which can be derived considering the ratio

$$\frac{\mathcal{B}(B \to X_s \sum_i \nu_i\bar{\nu}_i)}{\mathcal{B}(B \to X_c\ell\bar{\nu}_\ell)} = 3\, \frac{\alpha^2}{4\pi^2 \sin^4(\theta_W)} \left|\frac{V_{ts}}{V_{cb}}\right|^2 \frac{X(x_t)^2}{\eta_0\, f(m_c/m_b)}\, \overline{\eta} \qquad (9)$$

where the theoretical uncertainty related to the m_b^5 factor disappears. In Eq. (9) the factor 3 accounts for the sum over the three neutrino species. Using the phase space factor $f(m_c/m_b) \simeq 0.44$, the QCD correction factors $\eta_0 \simeq 0.87$ and $\overline{\eta} = 1 + (2\alpha_s(m_b)/3\pi)\left(\frac{25}{4} - \pi^2\right) \simeq 0.83$ [4], and the experimental measurement $\mathcal{B}(B \to X_c\ell\bar{\nu}_\ell) = (10.23 \pm 0.30) \times 10^{-2}$ (at $\Upsilon(4S)$) $\left[(10.95 \pm 0.32) \times 10^{-2}\right]$ (at Z^0) [15], one gets the SM prediction for the rate of $B \to X_s\nu\bar{\nu}$:

$$\mathcal{B}(B \to X_s \sum_i \nu_i\bar{\nu}_i) = (4.52 \pm 0.17) \times 10^{-5} \times \left[(4.84 \pm 0.14) \times 10^{-5}\right] \left(\frac{|V_{ts}|}{|V_{cb}|}\right)^2 \quad . \qquad (10)$$

The two predictions (10) refer to the measurements of the semileptonic branching ratio, performed at $\Upsilon(4S)$ and at LEP, respectively. The comparison of (10) with the bound (8) gives an insight into the necessary improvement of the experimental facilities. Such an improvement will be reached in the dedicated B-physics experiments planned for the next future, such as CLEO III at the Cornell e^+e^- storage ring [16], and BaBar [17] and Belle [18] asymmetric B-Factories now under construction at SLAC and KEK laboratories, respectively, which are expected to access B meson decay modes with branching ratios less than 10^{-5}. Therefore, the study of $b \to s\nu\bar{\nu}$, together with the search for $b \to s\ell^+\ell^-$ and $b \to s$ *gluon* processes, with a refinement of the measurement of $B \to X_s\gamma$ [1] and with the investigation of $B \to X_d\gamma$ will

[1] The decay $b \to s\gamma$ is the only flavour changing neutral current $b \to s$ transition observed so far [19]; within the errors, the experimental results, both in the inclusive and exclusive modes, are in agreement with the SM predictions, and already constrain a number of its extensions [20].

allow to exploit a complete program to test the SM properties at the loop level and constrain various new physics scenarios [21].

The first attempt to experimentally access the decay (1) will be through the exclusive modes, which will be better investigated at B factories. Among such modes, the channels $B \to K^{(*)}\nu\bar{\nu}$ are the prime candidates to look for, and therefore it is necessary to characterize them, by computing the fraction of the inclusive branching ratio represented by such exclusive modes, and the quantities that possibly will be used in the experimental analyses, for example the distribution of missing energy in the final state. This kind of analyses will be challenging at the future experimental facilities; however, it is possible that they will be successfully carried out, mainly at the asymmetric B-factories.

Let us first consider the channel $B \to K\nu\bar{\nu}$. In order to compute it, we need the matrix element of the effective hamiltonian (2) between the states of the initial B particle and the final particles $K, \nu, \bar{\nu}$. The hadronic transition $B \to K$ mediated by the vector $\bar{s}\gamma_\mu b$ current can be parameterized in terms of form factors, according to the notation in [22]

$$< K(p')|\bar{s}\gamma_\mu b|B(p) >= (p+p')_\mu F_1(q^2) + \frac{M_B^2 - M_K^2}{q^2} q_\mu \left[F_0(q^2) - F_1(q^2)\right] , \quad (11)$$

where $q = p - p'$ is the momentum transfer to the lepton pair, and the condition $F_1(0) = F_0(0)$ has been imposed to remove the singularity at $q^2 = 0$ in (11). The matrix element in Eq. (11) accounts for the long-distance interaction of quarks and gluons in the mesons, and must be computed by a non-perturbative approach. The result of three-point function QCD sum rules for F_1 [23] [2], in the range of squared momentum transfer $0 \le q^2 \le 15$ GeV^2 (where the method can be meaningfully applied) can be fitted by a polar q^2 dependence

$$F_1(q^2) = \frac{F_1(0)}{1 - q^2/M_P^2} , \quad (12)$$

[2] Other determinations of the form factors have been obtained by light-cone sum rules [24], quark models [25], lattice QCD [26] and χHQET [27].

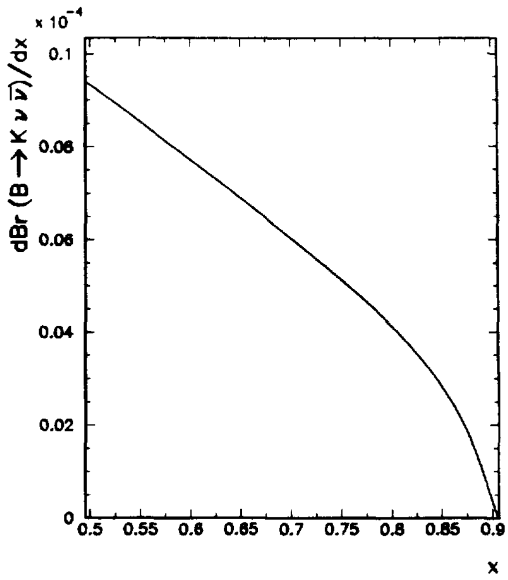

Fig. 1. Missing energy distribution in the decay $B \to K\nu\bar{\nu}$. The sum over the three neutrino species is understood.

with $F_1(0) = 0.25 \pm 0.03$ and $M_P = 5$ GeV.

The missing energy distribution in the decay $B \to K\nu\bar{\nu}$ can be computed using such F_1 (extrapolated up to $q^2_{\max} = (M_B - M_K)^2$); defining $E_{\rm miss}$ the energy of the neutrino pair in the B rest frame, and adopting the dimensionless variable $x = E_{\rm miss}/M_B$, we obtain, in the kinematically allowed range of x

$$\frac{1-r}{2} \le x \le 1 - \sqrt{r} \quad (13)$$

($r = M_K^2/M_B^2$), the result:

$$\frac{d\Gamma(B \to K\nu\bar{\nu})}{dx} = 3\,\frac{|c_L + c_R|^2\,|F_1(q^2)|^2}{48\pi^3 M_B}\sqrt{\lambda^3(q^2, M_B^2, M_K^2)} , \quad (14)$$

where $q^2 = M_B^2(2x-1) + M_K^2$. Eq. (14) shows that a possible NP interaction modifying the effective hamiltonian from (2) to (6) only changes the normalization of the spectrum.

In Fig. 1 the missing energy spectrum, Eq. (14), is plotted using (12) and fixing the Wilson coefficients c_L and c_R to the SM values $c_L = |c_L^{\rm SM}| = 2.7\ \left(|V_{ts}|/0.04\right) \times 10^{-9}$ GeV^{-2} and $c_R = 0$. The predicted branching ratio, using $\tau(B^-) = (1.65 \pm 0.04) \times 10^{-12}$ s, is

$$\mathcal{B}(B^- \to K^- \sum_i \nu_i \bar{\nu}_i)$$
$$= (2.4 \pm 0.6) \left(\frac{|V_{ts}|}{0.04}\right)^2 \times 10^{-6}, \qquad (15)$$

which corresponds to the ratio

$$R_K = \frac{\mathcal{B}(B \to K\nu\bar{\nu})}{\mathcal{B}(B \to X_s\nu\bar{\nu})} = (5.2 \pm 1.3) \times 10^{-2}$$
$$\times \left[(4.9 \pm 1.2) \times 10^{-2}\right]. \qquad (16)$$

For the decay $B \to K^*\nu\bar{\nu}$, the hadronic matrix element can be parameterized in terms of form factors as follows:

$$< K^*(p', \epsilon)|\bar{s}\gamma_\mu(1-\gamma_5)b|B(p) >$$
$$= \epsilon_{\mu\nu\alpha\beta}\epsilon^{*\nu}p^\alpha p'^\beta \frac{2V(q^2)}{M_B + M_{K^*}}$$
$$- i\left[\epsilon^*_\mu(M_B + M_{K^*})A_1(q^2)\right.$$
$$-(\epsilon^* \cdot q)(p+p')_\mu \frac{A_2(q^2)}{(M_B + M_{K^*})}$$
$$\left. - (\epsilon^* \cdot q)\frac{2M_{K^*}}{q^2}\big(A_3(q^2) - A_0(q^2)\big)q_\mu\right], \qquad (17)$$

where $A_0(0) = A_3(0)$, and

$$A_3(q^2) = \frac{1}{2M_{K^*}}\left[(M_B + M_{K^*})\ A_1(q^2)\right.$$
$$\left. + (M_{K^*} - M_B)\ A_2(q^2)\right].$$

The three-point QCD sum rule results for V, A_1 and A_2 (the only relevant form factors for a decay into massless neutrinos) are [23]

$$V(q^2) = \frac{V(0)}{1 - q^2/M_P^2}, \qquad (18)$$

with $V(0) = 0.47 \pm 0.03$ and $M_P = 5$ GeV, and

$$A_i(q^2) = A_i(0)(1 + \beta_i q^2), \qquad (19)$$

with $A_1(0) = 0.37 \pm 0.03$, $\beta_1 = -0.023$ GeV^{-2}, $A_2(0) = 0.40 \pm 0.03$, $\beta_2 = 0.034$ GeV^{-2}.

From these form factors it is easy to derive the missing energy distribution corresponding to the longitudinally and transversely polarized K^*:

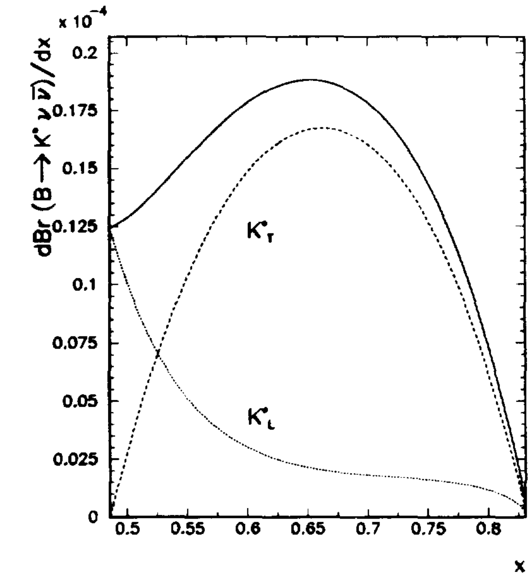

Fig. 2. Missing energy distribution in the decay $B \to K^*\nu\bar{\nu}$. Notations as in Fig. 1.

$$\frac{d\Gamma_L}{dx} = 3\frac{|c_L - c_R|^2}{24\pi^3}\frac{|\mathbf{p}'|}{M_{K^*}^2}$$
$$\times\left[(M_B + M_{K^*})(M_B E' - M_{K^*}^2)A_1(q^2)\right.$$
$$\left. - \frac{2M_B^2}{M_B + M_{K^*}}|\mathbf{p}'|^2 A_2(q^2)\right]^2, \qquad (20)$$

and

$$\frac{d\Gamma_\pm}{dx} = 3\frac{|\mathbf{p}'|q^2}{24\pi^3}\left|(c_L + c_R)\frac{2M_B|\mathbf{p}'|}{M_B + M_{K^*}}V(q^2)\right.$$
$$\left.\mp(c_L - c_R)(M_B + M_{K^*})A_1(q^2)\right|^2 \qquad (21)$$

where $\mathbf{p}'$ and E' are the K^* three-momentum and energy in the B meson rest frame. The missing energy distributions are plotted in Fig. 2, using the SM values of c_L and c_R as for $B \to K\nu\bar{\nu}$. Integrating over the full spectrum we obtain the prediction

$$\mathcal{B}(B^- \to K^{-*} \sum_i \nu_i \bar{\nu}_i)$$
$$= (5.1 \pm 0.8) \left(\frac{|V_{ts}|}{0.04}\right)^2 \times 10^{-6}, \qquad (22)$$

which corresponds to the ratio

$$R_{K^*} = \frac{\mathcal{B}(B \to K^* \nu\bar{\nu})}{\mathcal{B}(B \to X_s \nu\bar{\nu})} = (1.2 \pm 0.2) \times 10^{-1} \quad [(1.1 \pm 0.2) \times 10^{-1}] \quad . \tag{23}$$

The prediction (22) must be compared to the upper bound obtained by DELPHI Collaboration at LEP [28]: $\mathcal{B}(B_d^0 \to K^*(892)^0 \nu\bar{\nu}) < 1.0 \times 10^{-3}$ (at 90 % CL).

The results (15) and (22) are smaller than the estimates obtained in [29] assuming a heavy strange quark. Moreover, (15) and (22) show that the branching ratio of $B \to K^{(*)} \sum_i \nu_i \bar{\nu}_i$ is larger by a factor of five than the analogous $b \to s$ channels $B \to K^{(*)} \ell^+ \ell^-$ (excluding the long distance contribution from the conversion of $J/\psi, \psi'$ resonances). Finally, (16) and (23) imply that the exclusive $B \to K \nu\bar{\nu}$ and $B \to K^* \nu\bar{\nu}$ decays represent a fraction of $\simeq 20\%$ of the inclusive rate, with some correspondence with the $B \to K^* \gamma$ channel [30]:

$$\tilde{R}_{K^*} = \frac{\mathcal{B}(B \to K^* \gamma)}{\mathcal{B}(B \to X_s \gamma)} = 0.186 \pm 0.060 \, .$$

In the case of the radiative transition it was suggested [31] that a large contribution to the inclusive width should be due to the channel $B \to K_1 \gamma$, K_1 being a 1^+ orbital excitation of K^*. It could be interesting to investigate whether this is also the case of $B \to X_s \nu\bar{\nu}$, searching for a $K(n\pi)$ final state in the mass range $1.2 - 1.4$ GeV.

The experimental search for $B \to K^{(*)} \nu\bar{\nu}$ decays can be performed by looking for events with large missing energy, together with an opposite side fully reconstructed B meson. In the case of K^* at an asymmetric B factory, the vertex constraint obtained from the K^* decay products, together with a separation of the two B decay vertices, would provide an efficient background rejection. The hermiticity features of the detectors and the performances of the $K^{(*)}$ identification, together with the integrated luminosity corresponding to the full period of data taking, will be fundamental for a successful detection of such interesting decays. It is possible that the requirement of a complete reconstruction of the opposite-side B meson can be relaxed [32], and that a sufficient background rejection can be obtained by only imposing some energy and momentum conservation conditions. To check such a procedure, a Monte Carlo generator using a reliable set of form factors, as the set reported here, is required to simulate events in the various experimental environments.

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