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Radiative corrections to $b \to c\tau\bar{\nu}_\tau$

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Abstract

Analytical calculation is presented of the QCD radiative corrections to the rate of the process $b \to c\tau\bar{\nu}_\tau$ and to the τ lepton longitudinal polarization in $\tau\bar{\nu}_\tau$ rest frame. The results are given in the form of one dimensional infrared finite integrals over the invariant mass of the leptons. We argue that this form may be optimal for phenomelogical applications due to a possible breakdown of semilocal hadron-parton duality in decays of heavy flavours.

The semileptonic decay rate of B mesons is one of the key ingredients in the determination of weak mixing angles. Transitions of the b quark to the charmed as well as to the up quark have been analysed in great detail, exploiting either the inclusive decay rate or exclusive channels. The analysis of the inclusive rate is, however, affected by the uncertainty in the b mass and by bound state corrections. This problem is only partly circumvented by fixing $m_b - m_c$ through the difference of bottom and charmed meson masses and by relating the bound state effects to phenomenological constants that can be determined from other observables in the context of the heavy quark effective theory (HQET). The decay rate is, furthermore, affected by perturbative QCD corrections, which have been calculated analytically up to order α_s for arbitrary b and c masses and massless leptons [1, 2, 3]. Numerical results for perturbative QCD corrections to the partial decay rate $b \to \tau\bar{\nu}X$ have been recently obtained in Ref.[4]. The bound state corrections to this decay chanel up to order $1/m_b^2$ are also known [4, 5, 6] in the context of the HQET. The comparison between decays into light (μ, e) and heavy (τ) leptons may furthermore help to test the theoretical approach and allow to fix some of the free parameters. Including $b \to u$ transitions, four kinematically different leptonic decay modes are thus available for the comparison.

In this paper analytical results for the decay rate $b \to c\tau\bar{\nu}_\tau$ are presented in the form of an one dimensional integral over the invariant mass squared w^2 of the leptonic system. This formulation allows, at least in principle, the separation of the region of relatively small w^2, where the inclusive parton model description based on local parton-hadron duality should work, from the region of large w^2 where only one or few resonances are produced and the duality between the parton and hadron description may be doubtful[1]. In the region of large w^2 i.e. close to the Shifman-Voloshin limit[8] $\mathrm{w}^2 = \mathrm{w}^2_{max}$ the theoretical description of lepton spectra based on summation over exclusive channels is particularly simple and reliable. The HQET approach for exclusive decays[9, 10] on the other hand becomes quite involved if not impractical in the region of small w^2 which is dominated by multiparticle final states.

The calculation of the lowest order rate as well as corrections can be related in a straightforward way to the corresponding calculations for the decay into a virtual W boson with the subsequent integration over the mass of the $l\bar{\nu}$ system. The differential decay rate is proportional to

$$\mathcal{H}^{\alpha\beta}\,\mathcal{L}_{\alpha\beta}\,\mathrm{dPS}(b \to c\tau\bar{\nu})$$

$\mathcal{H}_{\alpha\beta}$ depends on quark and gluon fields and

$$\begin{aligned}\mathcal{L}_{\alpha\beta}(\tau;\nu) &\sim \sum_s \left[\,\bar{u}_\tau\gamma_\alpha(1-\gamma_5)v_\nu\,\right]\left[\,\bar{u}_\tau\gamma_\beta(1-\gamma_5)v_\nu\,\right]^\dagger \\ &\sim \nu_\alpha\tau_\beta + \tau_\alpha\nu_\beta - \nu\cdot\tau g_{\alpha\beta} - \mathrm{i}\varepsilon_{\alpha\beta\gamma\delta}\nu^\gamma\tau^\delta \end{aligned} \quad (1)$$

[1]In a recent preprint [7] the breakdown of the local parton-hadron duality has been invoked as the origin of problems with the semileptonic decay rate of D mesons in the framework of HQET. Let us remark that for large w^2 the kinetic energy of the hadronic system in B decays can be similar to that in D decays. Thus the semileptonic branching ratios for b decays may be also affected for the effective mass of the hadronic system close to the resonance region.

where τ^α and ν^α are the four-momenta of τ and $\bar{\nu}_\tau$. The phase space for the decay of b into $c\tau\bar{\nu}$ is, in the standard way, decomposed into a sequence of two-particle final states

$$\mathrm{dPS}(b \to c\tau\bar{\nu}) \sim \mathrm{dw}^2\, \mathrm{dPS}(b \to c\mathrm{w})\, \mathrm{dPS}(\mathrm{w} \to \tau\bar{\nu}) \tag{2}$$

where $\mathrm{w}^\alpha = \tau^\alpha + \nu^\alpha$. Then, it is straightforward to show that

$$\begin{aligned} \int \mathrm{dPS}(\mathrm{w} \to \tau\bar{\nu})\mathcal{L}_{\alpha\beta} &\sim \mathcal{A}\left(m_\tau^2/\mathrm{w}^2\right)\, T^{(0)}_{\alpha\beta} + \mathcal{B}\left(m_\tau^2/\mathrm{w}^2\right)\, T^{(1)}_{\alpha\beta} = \\ &\left(1 - m_\tau^2/\mathrm{w}^2\right)^2 \left[\left(1 + 2m_\tau^2/\mathrm{w}^2\right)\mathrm{w}_\alpha \mathrm{w}_\beta - \left(\mathrm{w}^2 + \tfrac{1}{2}m_\tau^2\right) g_{\alpha\beta}\right] \end{aligned} \tag{3}$$

where

$$\begin{aligned} T^{(0)}_{\alpha\beta} &= \mathrm{w}_\alpha \mathrm{w}_\beta \\ T^{(1)}_{\alpha\beta} &= \mathrm{w}_\alpha \mathrm{w}_\beta - \mathrm{w}^2 g_{\alpha\beta} \end{aligned} \tag{4}$$

It follows that the decay rate can be split accordingly into two incoherent terms which are related to weak decays of a heavy quark Q into another quark q and a real spin one or a spin zero boson. The relative weight of spin one versus spin zero contributions is governed by their respective spectral functions and it can be derived from eq.(3). For massless leptons $\mathcal{A}(0) = 0$ and $\mathcal{B}(0) = 1$, and therefore only the spin one (transversal) component ($\sim T^{(1)}_{\alpha\beta}$) contributes. The result is given in [2]. For fixed w^2 this contribution to the rate can be obtained from the formula for $t \to bW$, the top quark decay into b quark and a real W boson [2][2]. Multiplying this formula by $\mathcal{B}\left(m_\tau^2/\mathrm{w}^2\right)$ one obtains the contribution of the spin one component for $m_\tau \neq 0$. The other (longitudinal) contribution ($\sim T^{(1)}_{\alpha\beta}$) can be in an analogous way related to the (yet unobserved) decay $t \to bH^+$ where H^+ denotes a charged Higgs boson. Let Q and q denote the four-momenta of the quarks in the initial and in the final state. For a two-body decay mode the momentum of the W boson is $W = Q - q$. In Born approximation

$$\begin{aligned} W_\mu \bar{u}(q)\,\gamma^\mu(1-\gamma_5)\,u(Q) &= \bar{u}(q)\,[\,(\hat{Q} - \hat{q})(1-\gamma_5)\,]\,u(Q) \\ &= \bar{u}(q)\,[\,(M-m) + (M+m)\gamma_5\,]\,u(Q) \end{aligned} \tag{5}$$

where the equations of motion

$$(\hat{Q} - M)\,u(Q) = 0 \qquad \text{and} \qquad \bar{u}(q)\,(\hat{q} - m) = 0 \tag{6}$$

have been used. The last line in (5) can be interpreted as the amplitude of the decay $Q \to qH$ where H is a spin zero boson whose coupling to the weak quark current is given by

$$g = (M - m) + (M - m)\gamma_5 \tag{7}$$

[2]Note that the rates of up and down type quark decays into their respective isospin partners are of course identical, in contrast to the shapes of the lepton spectra.

Although not applicable for individual Feynman diagrams, the same relation holds true for the longitudinal contribution ($\sim W_\mu$) to the decay amplitude when $\mathcal{O}(\alpha_s)$ QCD corrections are included [11]. Therefore this contribution to the rate $b \to c\tau\bar{\nu}_\tau$ can be extracted from a formula describing $t \to bH^+$ which has been given in [12]; cf. Model I therein.
Let us define now the following dimensionless quantities

$$\rho = m_c^2/m_b^2 \qquad \eta = m_\tau^2/m_b^2 \qquad \text{and} \qquad t = \mathrm{w}^2/m_b^2$$

and the kinematic functions

$$\begin{aligned}
p_0(t) &= (1-t+\rho)/2 \\
p_3(t) &= \sqrt{p_0^2-\rho} \\
p_\pm(t) &= p_0 \pm p_3 = 1 - w_\mp(t) \\
Y_p(t) &= \tfrac{1}{2}\ln\left(p_+/p_-\right) = \ln\left(p_+/\sqrt{\rho}\right) \\
Y_w(t) &= \tfrac{1}{2}\ln\left(w_+/w_-\right) = \ln\left(w_+/\sqrt{t}\right)
\end{aligned} \tag{8}$$

where $\mathrm{w}^\alpha = \tau^\alpha + \nu^\alpha$ and w^2 denotes the effective mass of $\tau\bar{\nu}_\tau$. The partial rate of the decay $b \to c\tau\bar{\nu}_\tau$ is given by

$$\Gamma(b \to c\tau\bar{\nu}_\tau) = \int_\eta^{(1-\sqrt{\rho})^2} \frac{\mathrm{d}\Gamma}{\mathrm{d}t}\,\mathrm{d}t \tag{9}$$

with the differential rate

$$\frac{\mathrm{d}\Gamma}{\mathrm{d}t} = \Gamma_{bc}\left(1-\frac{\eta}{t}\right)^2\left\{\left(1+\frac{\eta}{2t}\right)\left[\mathcal{F}_0(t) - \frac{2\alpha_s}{3\pi}\mathcal{F}_1(t)\right] + \frac{3\eta}{2t}\left[\mathcal{F}_0^s(t) - \frac{2\alpha_s}{3\pi}\mathcal{F}_1^s(t)\right]\right\} \tag{10}$$

$$\Gamma_{bc} = \frac{G_F^2 m_b^5 |V_{cb}|^2}{192\pi^3} \tag{11}$$

$$\mathcal{F}_0(t) = 4p_3\left[(1-\rho)^2 + t(1+\rho) - 2t^2\right] \tag{12}$$

$$\mathcal{F}_0^s(t) = 4p_3\left[(1-\rho)^2 - t(1+\rho)\right] \tag{13}$$

$$\mathcal{F}_1(t) = \mathcal{A}_1\Psi + \mathcal{A}_2 Y_w + \mathcal{A}_3 Y_p + \mathcal{A}_4 p_3 \ln\rho + \mathcal{A}_5 p_3 \tag{14}$$

$$\mathcal{F}_1^s(t) = \mathcal{B}_1\Psi + \mathcal{B}_2 Y_w + \mathcal{B}_3 Y_p + \mathcal{B}_4 p_3 \ln\rho + \mathcal{B}_5 p_3 \tag{15}$$

$$\begin{aligned}
\Psi = {} & 8\ln(2p_3) - 2\ln t + \left[2\mathrm{Li}_2(w_-) - 2\mathrm{Li}_2(w_+) + 4\mathrm{Li}_2(2p_3/p_+)\right. \\
& \left. - 4Y_p\ln(2p_3/p_+) - \ln p_-\ln w_+ + \ln p_+\ln w_-\right] 2p_0/p_3
\end{aligned} \tag{16}$$

$$\begin{aligned}
\mathcal{A}_1 &= \mathcal{F}_0(t) \\
\mathcal{A}_2 &= -8(1-\rho)\left[1+t-4t^2-\rho(2-t)+\rho^2\right]
\end{aligned}$$

$$\begin{aligned}
\mathcal{A}_3 &= -2\left[3+6t-21t^2+12t^3-\rho(1+12t+5t^2)+\rho^2(11+2t)-\rho^3\right] \\
\mathcal{A}_4 &= -6\left[1+3t-4t^2-\rho(4-t)+3\rho^2\right] \\
\mathcal{A}_5 &= -2\left[5+9t-6t^2-\rho(22-9t)+5\rho^2\right] \qquad (17) \\
\mathcal{B}_1 &= \mathcal{F}_0^s(t) \\
\mathcal{B}_2 &= -8(1-\rho)\left[(1-\rho)^2-t(1+\rho)\right] \\
\mathcal{B}_3 &= -4(1-\rho)^4/t-2(-1+3\rho+15\rho^2-5\rho^3)+8(1+\rho)t-6(1+\rho)t^2 \\
\mathcal{B}_4 &= -4(1-\rho)^3/t-2(1-\rho)(1-11\rho)+6(1+3\rho)t \\
\mathcal{B}_5 &= -6(1-3\rho)(3-\rho)+18t(1+\rho) \qquad (18)
\end{aligned}$$

The integral in eq.(9) can be easily performed for the Born contribution. It reads:

$$\begin{aligned}
\Gamma_0(b\to c\tau\bar{\nu}_\tau) = \Gamma_{bc}\Big\{ & 24\left[\eta^2(1-\rho^2)\,\mathcal{Y}_w+\rho^2(1-\eta^2)\,\mathcal{Y}_p\right]+ \\
& 2\mathcal{P}_3\left(1-7\eta-7\eta^2+\eta^3-7\rho+12\eta\rho-7\eta^2\rho-7\rho^2-7\eta\rho^2+\rho^3\right)\Big\} \qquad (19)
\end{aligned}$$

where

$$\mathcal{P}_3 = p_3(\eta), \qquad \mathcal{Y}_p = Y_p(\eta) \qquad \text{and} \qquad \mathcal{Y}_w = Y_w(\eta).$$

In principle the first order QCD correction can be also expressed in terms of polylogarithms. In particular for $\eta = 0$ the formula (12) of [3] is obtained. However, the complete result is lenghty. From the practical point of view it is much simpler to evaluate the integral in eq.(9) numerically. Moreover, as it has been explained, for t close to the Shifman-Voloshin limit the exclusive description is preferable, thus, for applications, it may be better to perform this integral only over a part of the available phase space. In recent articles [15, 16] the size of $\mathcal{O}(\alpha_s^2)$ corrections has been estimated using the scheme of Brodsky, Lepage and Mackenzie [17] for fixing the scale μ of α_s. It has turned out that this scale is rather low. This suggests that next-to-leading QCD corrections are large. The most serious problems arise in the region close to the no-recoil $t = t_{max}$ point. It is well known, cf. [18, 19], that at the boundaries of the available phase space logarithmic divergences may appear as remnants of infrared divergences. This phenomenon arises when real radiation becomes supressed relative to virtual corrections. This is exactly what happens for t in the neighbour of t_{max} and once again one is led to the conclusion that the inclusive parton-like description may break down there.
In the massless limit $\rho \to 0$ which corresponds to $b \to u\tau\bar{\nu}_\tau$ transition the functions in eq.(15) simplify considerably

$$\begin{aligned}
\mathcal{F}_0(t) &= 2(1-t)^2(1+2t) \qquad (20) \\
\mathcal{F}_0^s(t) &= 2(1-t)^2 \qquad (21) \\
\mathcal{F}_1(t) &= \mathcal{F}_0(t)\left[\tfrac{2}{3}\pi^2+4\mathrm{Li}_2(t)+2\ln t\ln(1-t)\right]-(1-t)(5+9t-6t^2) \\
&\quad +4t(1-t-2t^2)\ln t+2(1-t)^2(5+4t)\ln(1-t) \qquad (22)
\end{aligned}$$

$$\mathcal{F}_1^s(t) \;=\; \mathcal{F}_0^s(t)\,\left[\tfrac{2}{3}\pi^2 + 4\mathrm{Li}_2(t) - \tfrac{9}{2} + \ln(1-t)\left(2\ln t - \tfrac{2}{t} + 5\right)\right] + 4(1-t)t\ln t \tag{23}$$

After integration over t one derives the following expression for the total partial rate of $b \to u\tau\bar{\nu}_\tau$

$$\begin{aligned}
\frac{1}{\Gamma_{bu}}\Gamma(b \to u\tau\bar{\nu}_\tau) &= 1 - 8\eta + 8\eta^3 - \eta^4 - 12\eta^2\ln\eta \\
&-\frac{2\alpha_s}{3\pi}\Big\{(-1+\eta)(75 - 539\eta - 476\eta^2 + 18\eta^3)/12 \\
&\quad +(3 - 24\eta - 36\eta^2 + 16\eta^3 - 2\eta^4)\pi^2/3 \\
&\quad +72\eta^2\left[\zeta(3) - \mathrm{Li}_3(\eta)\right] + 2(1 - 8\eta + 36\eta^2 + 16\eta^3 - 2\eta^4)\mathrm{Li}_2(\eta) \\
&\quad +(1-\eta^2)(31 - 320\eta + 31\eta^2)\ln(1-\eta)/6 \\
&\quad +\Big[2\eta + 15\eta^2 - 94\eta^3/3 + 31\eta^4/6 - 8\eta^2\pi^2 + 24\eta^2\mathrm{Li}_2(\eta) \\
&\qquad +2(1-\eta^2)(1 - 8\eta + \eta^2)\ln(1-\eta)\Big]\ln(\eta)\Big\}
\end{aligned} \tag{24}$$

$$\Gamma_{bu} = \frac{G_F^2 m_b^5 |V_{ub}|^2}{192\pi^3} \tag{25}$$

which agrees with eqs.(2.21) and (3.6) in a recent preprint [13].

It has been argued [14] that the τ polarization in $b \to c\tau\bar{\nu}_\tau$ is particularly sensitive to deviations from the Standard Model. The longitudinal component of τ polarization can either be measured with reference to its direction of flight in the b rest frame or, alternatively, relative to its direction of flight in the $\tau\bar{\nu}$ rest frame. To evaluate analytically QCD corrections to the longitudinal τ polarization in the b rest frame is a demanding task. Previous experience from a similar calculation for the decay of polarized top quarks indicates that these corrections are typically quite small through most of the kinematical range. To substantiate this claim, the longitudinal τ polarization in the $\tau\bar{\nu}$ rest frame is evaluated including QCD corrections.
For V-A weak current the leptonic tensor

$$\begin{aligned}
\mathcal{L}_{\alpha\beta}(\tau, s; \nu) &\sim \left[\bar{u}_\tau\gamma_\alpha(1-\gamma_5)v_\nu\right]\left[\bar{u}_\tau\gamma_\beta(1-\gamma_5)v_\nu\right]^\dagger \\
&\sim \nu_\alpha\mathcal{T}_\beta + \mathcal{T}_\alpha\nu_\beta - \nu\cdot\mathcal{T}\,g_{\alpha\beta} - \mathrm{i}\varepsilon_{\alpha\beta\gamma\delta}\nu^\gamma\mathcal{T}^\delta
\end{aligned} \tag{26}$$

where

$$\mathcal{T}^\alpha = \tfrac{1}{2}\left(\tau^\alpha - m_\tau s^\alpha\right) \tag{27}$$

and s^α is the τ polarization fourvector. The helicity states in the $\tau\bar{\nu}$ rest frame are obtained for

$$m_\tau s^\alpha = \pm\left(\tau^\alpha - \tfrac{2\eta}{t-\eta}\nu^\alpha\right) \tag{28}$$

It is evident that the net polarization can be calculated in the same way as the total rate, decomposing again the spin dependent term into longitudinal and transversal parts.

For the positive helicity of τ one derives the following expression for the differential rate

$$\frac{1}{\Gamma_{bc}}\frac{\mathrm{d}\Gamma^{(+)}}{\mathrm{d}t} = \frac{\eta}{2t}\left(1-\frac{\eta}{t}\right)^2 \left\{\mathcal{F}_0(t) + 3\mathcal{F}_0^s(t) - \frac{2\alpha_s}{3\pi}\left[\mathcal{F}_1(t) + 3\mathcal{F}_1^s(t)\right]\right\} \tag{29}$$

For the negative helicity one has

$$\frac{\mathrm{d}\Gamma^{(-)}}{\mathrm{d}t} = \frac{\mathrm{d}\Gamma}{\mathrm{d}t} - \frac{\mathrm{d}\Gamma^{(+)}}{\mathrm{d}t} \tag{30}$$

and the net helicity in the $\tau\nu_\tau$ rest frame is equal to

$$\mathrm{P} = -1 + \frac{2\int_\eta^{(1-\sqrt{\rho})^2}\frac{\mathrm{d}\Gamma^{(+)}}{\mathrm{d}t}\,\mathrm{d}t}{\int_\eta^{(1-\sqrt{\rho})^2}\frac{\mathrm{d}\Gamma}{\mathrm{d}t}\,\mathrm{d}t} \tag{31}$$

In the subsequent discussion the mass difference $m_b - m_c$ will be fixed through the HQET relation

$$m_b - m_c = [(3m_{B^*} + m_B) - (3m_{D^*} + m_D)]/4 + \ldots$$

The most important corrections to the above relation arise from the kinetic energy of heavy quarks in B and D mesons. On physical grounds one expects this contribution to increase the difference between m_b and m_c. We take $m_b - m_c = 3.4$ GeV in numerical calculations and m_b is varied between 4.5 and 5.0 GeV. In order to estimate effects of QCD corrections on measurable quantities we neglects all the problems with the parton-hadron duality and the scale ambiguity. The following results have been obtained:

$$\begin{aligned}
&\Gamma(b \to c\tau\bar{\nu}_\tau) = \Gamma_0(b \to c\tau\bar{\nu}_\tau)\,[\,1 + (-0.450 \pm 0.016)\,\alpha_s\,]\\
&\Gamma(b \to ce\bar{\nu}_e) = \Gamma_0(b \to ce\bar{\nu}_e)\,[\,1 + (-0.545 \pm 0.025)\,\alpha_s\,]\\
&R = BR(b \to \tau X)/BR(b \to eX) = R_0\,[\,1 + (0.094 \mp 0.009)\,\alpha_s\,]\\
&\mathrm{P} = -(0.293 \mp 0.002)\,[\,1 + (0.140 \mp 0.015)\,\alpha_s\,]\\
&<\,t\,>(b \to \tau X) = (0.34 \pm 0.03)\,[\,1 + (0.016 \pm 0.003)\,\alpha_s\,]\\
&<\,t\,>(b \to eX) = (0.20 \pm 0.02)\,[\,1 + (0.035 \pm 0.007)\,\alpha_s\,]
\end{aligned}$$

where $<\,t\,>(\ldots)$ denote the average values of t for the corresponding decay chanels. It is evident that the $\mathcal{O}(\alpha_s)$ corrections practically cancel in the ratio R of the branching ratios as well as in the result for the polarization P. The moments $<\,t\,>(b \to \tau X)$ and $<\,t\,>(b \to eX)$ are also insensitive to α_s corrections. On the other hand all these quantities are sensitive to the quark masses and therefore they may be used for fixing m_b and m_c.

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