Eur. Phys. J. C 45, 151–158 (2006)
Digital Object Identifier (DOI) 10.1140/epjc/s2005-02407-6

THE EUROPEAN
PHYSICAL JOURNAL C

Study of FCNC mediated Z boson effect in the semileptonic rare baryonic decays $\Lambda_b \to \Lambda l^+ l^-$

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Received: 2 June 2005 /
Published online: 26 October 2005 –

Abstract. We study the effect of the FCNC mediated Z boson in the rare semileptonic baryonic decays $\Lambda_b \to \Lambda l^+ l^-$. We consider the model where the standard model fermion sector is extended by an extra vector-like down quark, as a consequence of which it allows for CP-violating Z mediated flavor changing neutral current at the tree level. We find that due to this non-universal Zbs coupling, the branching ratios of the rare semileptonic Λ_b decays are enhanced reasonably from their corresponding standard model values and the zero point of the forward–backward asymmetry for $\Lambda_b \to \Lambda \mu^+ \mu^-$ is shifted to the left.

PACS. 13.30.Ce, 12.60.-i, 11.30.Hv

1 Introduction

In recent years B-physics (studies relating to particles containing a bottom quark or antiquark) has been a very active area of research, both experimentally as well as theoretically. Giant machines, named B-factories, have been constructed and are being in operation now to study the dynamics of these heavy particles. This in turn will help us to verify our understanding based on the predictions of the standard model (SM), which has been very successful so far. Predictions based on the SM remain almost unchallenged even today, except possibly for the neutrino having a mass. On the other hand, the SM contains many parameters which are unknown to a satisfactory level of accuracy. Also the SM does not provide any explanation as to why there are only three generation of fermions, the mass hierarchy among the fermions etc. There are many variants of possible extensions to the SM existing in the literature and it is widely believed that physics beyond the SM might be discovered soon.

Unfortunately, we have not been able to see any indication of physics beyond the SM in the currently running B-factories (SLAC and KEK). Nevertheless, there appears to be some kind of deviation in some $b \to s$ penguin induced transitions (like the deviation in the measurement of $\sin 2\beta$ in $B_d \to \phi K_S$ and also in some related processes [1], the polarization anomaly in $B \to \phi K^*$ [2] and the deviation of the branching ratios from the SM expectation in some rare B decays, etc.). But it is too early to claim or rule out the existence of new physics in the b-sector. On the other hand, B-factories are expected to continue accumulating data for some more years and thereafter the task will be taken over by the new second generation B-experiments such as LHC-b and BTeV. It is therefore perceived very strongly that, in the future, we will be able to identify the existence of new physics (NP), if really there is any. It should be mentioned here that one of the main objectives behind the pursuit of B-factory experiments is to look for physics beyond the SM.

One of the important ways to look for new physics in the b-sector is the analysis of rare B decay modes, which are induced by the flavor changing neutral current (FCNC) transitions. The FCNC transitions generally arise at the loop level in the SM, and thus provide an excellent testing ground for new physics. Therefore, it is very important to study the FCNC processes, both theoretically and experimentally, as these decays can provide a sensitive test for the investigation of the gauge structure of the SM at the loop level. Concerning the semileptonic B decays, $B \to X_s l^+ l^-$ ($X_s = K, K^*,\ l = e, \mu, \tau$) are a class of decays having both theoretical and experimental importance. At the quark level, these decays proceed through the FCNC transition $b \to s$, which occur only through loops in the SM. For the very same reason, the study of the FCNC decays can provide a sensitive test for the investigation of the gauge structure of the SM at the loop level. At the same time, these decays constitute a quite suitable tool of looking for new physics. New physics effects manifest themselves in these rare B decays in two different ways: either through a new contribution to the Wilson coefficients or through the new structure in the effective Hamiltonian, which are both absent in the SM.

It is well known that the theoretical analysis of the inclusive decay is easy, but their experimental detection

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is difficult. For exclusive decays the situation is opposite, i.e., these decays can be easily studied in the experiments but theoretically they have drawbacks, and predictions are model dependent. This is due to the fact that in calculating the branching ratios and other observables for exclusive decays we face the problem of computing the hadronic form factors.

Therefore, the exclusive processes induced by the quark level transition $b \to s l^+ l^-$ have received a considerable attention in the literature because of its richness to study the FCNC. Moreover, the dileptons present in these processes allow us to formulate many observables which can serve as a testing ground to decipher the presence of new physics. With the accumulation of data, day by day in the b-sector, we are in an increasingly better position to experimentally study the semileptonic decay induced by the $b \to s$ transition. In this context a rich and extensive study exists [3] as far as the rare decay process $B \to K l^+ l^-$ and its vector counterpart $B \to K^* l^+ l^-$ are concerned, in the framework of the SM and in many extensions of it. However, the study of baryonic rare semileptonic decay modes, also induced by the same quark level transition, i.e., $b \to s l^+ l^-$, are also as important as its mesonic counterparts and deserve serious attention, both theoretically and experimentally. Since at the quark level they are induced by the same mechanism, we can independently test our understanding of the quark–hadron dynamics and also study some CP-violation parameters with the help of baryonic rare decays, apart from corroborating the findings of the mesonic sector.

In this work, we would like to analyze the rare baryonic decay mode $\Lambda_b \to \Lambda l^+ l^-$. We consider the effect of the non-universal Z boson which induces FCNC interaction at the tree level. It is well known that FCNC coupling of the Z boson can be generated at the tree level in various exotic scenarios. Two popular examples discussed in the literature are the models with an extra $U(1)$ symmetry [4] and those with the addition of a non-sequential generation of quarks [5]. In the case of extra $U(1)$ symmetry the FCNC couplings of the Z boson are induced by Z–Z' mixing, provided the SM quarks have family non-universal charges under the new $U(1)$ group. In the second case, adding a different number of up- and down-type quarks, the pseudo-CKM matrix needed to diagonalize the charged currents is no longer unitary and this leads to tree level FCNC couplings. Here we will follow the second approach to analyze the semileptonic rare Λ_b decays. These decays are studied in the SM [6], in the supersymmetric model with and without R-parity [7] and in a model independent way by Aliev et al. [8]. To have a complete understanding of the nature of the new physics, if it indeed exists, it would be worthwhile to analyze these rare decays in as many new physics models as possible. In this paper, we would like to see the effect of the non-universal Z boson in the decay width and forward–backward asymmetry of the lepton pairs, when the final Λ baryon is unpolarized. We also study its effect on the polarization of Λ baryon.

This paper is organized as follows. In Sect. 2 we briefly describe the decay parameters of the semileptonic rare decays in the standard model. In Sect. 3 the effect of the FCNC mediated Z boson has been considered. The numerical results are presented in Sect. 4, and Sect. 5 contains the conclusion.

2 Standard model contribution

The decay process $\Lambda_b \to \Lambda l^+ l^-$ is described by the quark level transition $b \to s l^+ l^-$. Thus, the effective Hamiltonian describing this process can be given as follows [9]:

$$\mathcal{H}_{\text{eff}} = \frac{G_{\text{F}}\alpha}{\sqrt{2}\,\pi} V_{tb} V_{ts}^* \times \Big[C_9^{\text{eff}} (\bar{s}\gamma_\mu L b)(\bar{l}\gamma^\mu l) + C_{10} (\bar{s}\gamma_\mu L b)(\bar{l}\gamma^\mu \gamma_5\, l) - 2C_7^{\text{eff}} m_b \left(\bar{s}\mathrm{i}\sigma_{\mu\nu} \frac{q^\mu}{q^2} R b \right) (\bar{l}\gamma^\mu l) \Big], \quad (1)$$

where q is the momentum transferred to the lepton pair, given by $q = p_- + p_+$; p_- and p_+ are the momenta of the leptons l^- and l^+ respectively. $R, L = (1 \pm \gamma_5)/2$, and the C_i's are the Wilson coefficients evaluated at the b quark mass scale. The values of these coefficients in NLL order are [10]

$$C_7^{\text{eff}} = -0.308, \quad C_9 = 4.154, \quad C_{10} = -4.261. \quad (2)$$

The coefficient C_9^{eff} has a perturbative part and a resonance part which comes from the long distance effects due to the conversion of the real $c\bar{c}$ into the lepton pair l^+l^-. Therefore, one can write it as

$$C_9^{\text{eff}} = C_9 + Y(s) + C_9^{\text{res}}, \quad (3)$$

where $s = q^2$ and the function $Y(s)$ denotes the perturbative part coming from one loop matrix elements of the four quark operators and is given by [11]

$$\begin{aligned} Y(s) = {} & g(m_c, s)(3C_1 + C_2 + 3C_3 + C_4 + 3C_5 + C_6) \\ & - \frac{1}{2} g(0, s)(C_3 + 3C_4) \\ & - \frac{1}{2} g(m_b, s)(4C_3 + 4C_4 + 3C_5 + C_6) \\ & + \frac{2}{9}(3C_3 + C_4 + 3C_5 + C_6), \end{aligned} \quad (4)$$

where

$$\begin{aligned} g(m_i, s) = {} & -\frac{8}{9} \ln(m_i/m_b^{\text{pole}}) + \frac{8}{27} + \frac{4}{9} y_i \\ & - \frac{2}{9}(2 + y_i)\sqrt{|1 - y_i|} \\ & \times \left\{ \Theta(1 - y_i) \left[\ln\left(\frac{1 + \sqrt{1 - y_i}}{1 - \sqrt{1 - y_i}} \right) - \mathrm{i}\pi \right] \right. \\ & \left. + \Theta(y_i - 1) 2 \arctan \frac{1}{\sqrt{y_i - 1}} \right\}, \end{aligned} \quad (5)$$

with $y_i = 4m_i^2/s$. The values of the coefficients C_i in NLL order are taken from [10] as $C_1 = -0.151$, $C_2 =$

1.059, $C_3 = 0.012$, $C_4 = -0.034$, $C_5 = 0.010$ and $C_6 = -0.040$.

The long distance resonance effect is given as [12]

$$C_9^{\rm res} = \frac{3\pi}{\alpha^2}(3C_1 + C_2 + 3C_3 + C_4 + 3C_5 + C_6) \times \sum_{V_i=\psi(1S),\ldots,\psi(6S)} \kappa_{V_i} \frac{m_{V_i}\Gamma(V_i \to l^+l^-)}{m_{V_i}^2 - s - i m_{V_i}\Gamma_{V_i}}. \quad (6)$$

The phenomenological parameter κ is taken to be 2.3 so as to reproduce the correct branching ratio of $\mathcal{B}(B \to J/\psi K^* \to K^* l^+ l^-) = \mathcal{B}(B \to J/\psi K^*)\mathcal{B}(J/\psi \to l^+l^-)$.

After having an idea of the effective Hamiltonian and the relevant Wilson coefficients, we now proceed to evaluate the transition matrix elements for the process $\Lambda_b(p_{\Lambda_b}) \to \Lambda(p_\Lambda)\, l^+(p_+) l^-(p_-)$. For this purpose, we need to know the matrix elements of the various hadronic currents between the initial Λ_b and the final Λ baryon, which are parametrized in terms of various form factors by

$$\begin{aligned}
&\langle \Lambda|\bar{s}\gamma_\mu b|\Lambda_b\rangle = \bar{u}_\Lambda\left[f_1\gamma_\mu + {\rm i} f_2\sigma_{\mu\nu}q^\nu + f_3 q_\mu\right]u_{\Lambda_b},\\
&\langle \Lambda|\bar{s}\gamma_\mu\gamma_5 b|\Lambda_b\rangle\\
&= \bar{u}_\Lambda\left[g_1\gamma_\mu\gamma_5 + {\rm i} g_2\sigma_{\mu\nu}\gamma_5 q^\nu + g_3\gamma_5 q_\mu\right]u_{\Lambda_b},\\
&\langle \Lambda|\bar{s}{\rm i}\sigma_{\mu\nu}q^\nu b|\Lambda_b\rangle = \bar{u}_\Lambda\left[f_1^T\gamma_\mu + {\rm i} f_2^T\sigma_{\mu\nu}q^\nu + f_3^T q_\mu\right]u_{\Lambda_b},\\
&\langle \Lambda|\bar{s}{\rm i}\sigma_{\mu\nu}\gamma_5 q^\nu b|\Lambda_b\rangle\\
&= \bar{u}_\Lambda\left[g_1^T\gamma_\mu\gamma_5 + {\rm i} g_2^T\sigma_{\mu\nu}\gamma_5 q^\nu + g_3^T\gamma_5 q_\mu\right]u_{\Lambda_b}, \quad (7)
\end{aligned}$$

where $q = p_{\Lambda_b} - p_\Lambda = p_+ + p_-$ is the momentum transfer, and f_i and g_i are the various form factors which are functions of q^2. The number of independent form factors is greatly reduced in the heavy quark symmetry limit. In this limit, the matrix elements of all the hadronic currents, irrespective of their Dirac structure, can be given in terms of only two independent form factors [13] by

$$\langle \Lambda(p_\Lambda)|\bar{s}\Gamma b|\Lambda_b(p_{\Lambda_b})\rangle = \bar{u}_\Lambda[F_1(q^2) + \not{v}F_2(q^2)]\Gamma u_{\Lambda_b}, \quad (8)$$

where Γ is the product of Dirac matrices, and $v^\mu = p^\mu_{\Lambda_b}/m_{\Lambda_b}$ is the four velocity of Λ_b. These two sets of form factors are related to each other by

$$\begin{aligned}
&g_1 = f_1 = f_2^T = g_2^T = F_1 + \sqrt{r}F_2,\\
&g_2 = f_2 = g_3 = f_3 = \frac{F_2}{m_{\Lambda_b}},\\
&g_3^T = \frac{F_2}{m_{\Lambda_b}}(m_{\Lambda_b} + m_\Lambda), \qquad f_3^T = -\frac{F_2}{m_{\Lambda_b}}(m_{\Lambda_b} - m_\Lambda)\\
&f_1^T = g_1^T = \frac{F_2}{m_{\Lambda_b}}q^2, \quad (9)
\end{aligned}$$

where $r = m_\Lambda^2/m_{\Lambda_b}^2$. Thus, using these form factors, the transition amplitude can be written as

$$\begin{aligned}
\mathcal{M}(\Lambda_b \to \Lambda l^+l^-) &= \frac{G_{\rm F}\alpha}{\sqrt{2}\pi}V_{tb}V_{ts}^*\\
&\times\left[\bar{l}\gamma_\mu l\left\{\bar{u}_\Lambda\left(\gamma^\mu(A_1P_R + B_1P_L)\right.\right.\right.\\
&\left.\left.+ {\rm i}\sigma^{\mu\nu}q_\nu(A_2P_R + B_2P_L)\right)u_{\Lambda_b}\right\}\\
&+ \bar{l}\gamma_\mu\gamma_5 l\\
&\times\left\{\bar{u}_\Lambda\left(\gamma^\mu(D_1P_R + E_1P_L) + {\rm i}\sigma^{\mu\nu}q_\nu(D_2P_R + E_2P_L)\right.\right.\\
&\left.\left.\left. + q^\mu(D_3P_R + E_3P_L)\right)u_{\Lambda_b}\right\}\right], \quad (10)
\end{aligned}$$

where the various parameters A_i, B_i and D_j, E_j ($i = 1,2$ and $j = 1,2,3$) are defined by

$$\begin{aligned}
A_i &= \frac{1}{2}C_9^{\rm eff}(f_i - g_i) - \frac{C_7 m_b}{q^2}(f_i^T + g_i^T),\\
B_i &= \frac{1}{2}C_9^{\rm eff}(f_i + g_i) - \frac{C_7 m_b}{q^2}(f_i^T - g_i^T),\\
D_j &= \frac{1}{2}C_{10}(f_j - g_j), \qquad E_j = \frac{1}{2}C_{10}(f_j + g_j). \quad (11)
\end{aligned}$$

Let us first consider the case when the final Λ baryon is unpolarized. The physical observables in this case are the differential decay rate and the forward–backward rate asymmetries. From the transition amplitude (10), one can obtain the double differential decay rate as

$$\frac{{\rm d}^2\Gamma}{{\rm d}\hat{s}{\rm d}z} = \frac{G_{\rm F}^2\alpha^2}{2^{12}\pi^5}|V_{tb}V_{ts}^*|^2 m_{\Lambda_b} v_l \lambda^{1/2}(1,r,\hat{s})\mathcal{K}(s,z), \quad (12)$$

where $\hat{s} = s/m_{\Lambda_b}^2$, $z = \cos\theta$, the angle between p_{Λ_b} and p_+ in the center of mass frame of l^+l^- pair, $v_l = \sqrt{1 - 4m_l^2/s}$ and $\lambda(a,b,c) = \sqrt{a^2 + b^2 + c^2 - 2(ab + bc + ca)}$ is the usual triangle function. The function $\mathcal{K}(s,z)$ is given as

$$\mathcal{K}(s,z) = \mathcal{K}_0(s) + z\mathcal{K}_1(s) + z^2\,\mathcal{K}_2(s), \quad (13)$$

with

$$\begin{aligned}
&\mathcal{K}_0(s)\\
&= 32m_l^2 m_{\Lambda_b}^4\hat{s}(1 + r - \hat{s})\left(|D_3|^2 + |E_3|^2\right)\\
&+64m_l^2 m_{\Lambda_b}^3(1 - r - \hat{s}){\rm Re}(D_1^*E_3 + D_3E_1^*)\\
&+64m_{\Lambda_b}^2\sqrt{r}(6m_l^2 - \hat{s}m_{\Lambda_b}^2){\rm Re}(D_1^*E_1)\\
&+64m_l^2 m_{\Lambda_b}^3\sqrt{r}\\
&\times\left(2m_{\Lambda_b}\hat{s}{\rm Re}(D_3^*E_3) + (1 - r + \hat{s}){\rm Re}(D_1^*D_3 + E_1^*E_3)\right)\\
&+32m_{\Lambda_b}^2(2m_l^2 + m_{\Lambda_b}^2\hat{s})\\
&\times\left((1 - r + \hat{s})m_{\Lambda_b}\sqrt{r}{\rm Re}(A_1^*A_2 + B_1^*B_2)\right.\\
&-m_{\Lambda_b}(1 - r - \hat{s}){\rm Re}(A_1^*B_2 + A_2^*B_1)\\
&\left.-2\sqrt{r}\left[{\rm Re}(A_1^*B_1) + m_{\Lambda_b}^2\hat{s}{\rm Re}(A_2^*B_2)\right]\right)\\
&+8m_{\Lambda_b}^2\left(4m_l^2(1 + r - \hat{s}) + m_{\Lambda_b}^2\left[(1 - r)^2 - \hat{s}^2\right]\right)\\
&\times\left(|A_1|^2 + |B_1|^2\right)\\
&+8m_{\Lambda_b}^4\\
&\times\left(4m_l^2[\lambda + (1 + r - \hat{s})\hat{s}] + m_{\Lambda_b}^2\hat{s}\left[(1 - r)^2 - \hat{s}^2\right]\right)\\
&\times\left(|A_2|^2 + |B_2|^2\right)\\
&-8m_{\Lambda_b}^2\left(4m_l^2(1 + r - \hat{s}) - m_{\Lambda_b}^2\left[(1 - r)^2 - \hat{s}^2\right]\right)
\end{aligned}$$

$$\begin{aligned}
&\times \left(|D_1|^2+|E_1|^2\right)\\
&+8m_{\Lambda_b}^5\hat{s}v_l^2\\
&\times \left(-8m_{\Lambda_b}\hat{s}\sqrt{r}\mathrm{Re}(D_2^*E_2)\right.\\
&+4(1-r+\hat{s})\sqrt{r}\mathrm{Re}(D_1^*D_2+E_1^*E_2)\\
&-4(1-r-\hat{s})\mathrm{Re}(D_1^*E_2+D_2^*E_1)\\
&\left.+m_{\Lambda_b}\left[(1-r)^2-\hat{s}^2\right]\left[|D_2|^2+|E_2|^2\right]\right), \qquad (14)
\end{aligned}$$

$$\begin{aligned}
&\mathcal{K}_1(s)\\
&=-16m_{\Lambda_b}^4\hat{s}v_l\sqrt{\lambda}\{2\mathrm{Re}(A_1^*D_1)-2\mathrm{Re}(B_1^*E_1)\\
&\quad+2m_{\Lambda_b}\mathrm{Re}(B_1^*D_2-B_2^*D_1+A_2^*E_1-A_1^*E_2)\}\\
&\quad+32m_{\Lambda_b}^5\hat{s}v_l\sqrt{\lambda}\{m_{\Lambda_b}(1-r)\mathrm{Re}(A_2^*D_2-B_2^*E_2)\\
&\quad+\sqrt{r}\mathrm{Re}(A_2^*D_1+A_1^*D_2-B_2^*E_1-B_1^*E_2)\}, \qquad (15)
\end{aligned}$$

and

$$\begin{aligned}
&\mathcal{K}_2(s)\\
&=8m_{\Lambda_b}^6v_l^2\lambda\hat{s}\left(\left(|A_2|^2+|B_2|^2+|D_2|^2+|E_2|^2\right)\right.\\
&\quad-8m_{\Lambda_b}^4v_l^2\,\lambda\,\left(|A_1|^2+|B_1|^2+|D_1|^2+|E_1|^2\right). \qquad (16)
\end{aligned}$$

The dilepton mass spectrum can be obtained from (12) by integrating out the angular dependent parameter z, which yields

$$\left(\frac{\mathrm{d}\Gamma}{\mathrm{d}s}\right)_0=\frac{G_F^2\alpha^2}{2^{11}\pi^5m_{\Lambda_b}}|V_{tb}V_{ts}^*|^2v_l\sqrt{\lambda}\left[\mathcal{K}_0(s)+\frac{1}{3}\mathcal{K}_2(s)\right], \qquad (17)$$

where λ is the short hand notation for $\lambda(1,r,\hat{s})$. The limits for s are

$$4m_l^2\le s\le(m_{\Lambda_b}-m_\Lambda)^2. \qquad (18)$$

Another observable is the lepton forward–backward asymmetry (A_{FB}), which is also a very powerful tool for looking new physics. The position of the zero value of A_{FB} is very sensitive to the presence of new physics.

The normalized forward–backward asymmetry is defined by

$$A_{\mathrm{FB}}(s)=\frac{\displaystyle\int_0^1\frac{\mathrm{d}\Gamma}{\mathrm{d}\hat{s}\mathrm{d}z}\mathrm{d}z-\int_{-1}^0\frac{\mathrm{d}\Gamma}{\mathrm{d}\hat{s}\mathrm{d}z}\mathrm{d}z}{\displaystyle\int_0^1\frac{\mathrm{d}\Gamma}{\mathrm{d}\hat{s}\mathrm{d}z}\mathrm{d}z+\int_{-1}^0\frac{\mathrm{d}\Gamma}{\mathrm{d}\hat{s}\mathrm{d}z}\mathrm{d}z}. \qquad (19)$$

Thus one obtains from (12)

$$A_{\mathrm{FB}}(s)=\frac{\mathcal{K}_1(s)}{\mathcal{K}_0(s)+\mathcal{K}_2(s)/3}. \qquad (20)$$

Now let us consider the case when the final Λ baryon is polarized. To study its spin polarization, we write the spin vector of Λ in terms of a unit vector $\hat{\eta}$ along the direction of Λ spin in its rest frame as

$$s_\mu=\left(\frac{\boldsymbol{p}_\Lambda\cdot\hat{\eta}}{m_\Lambda},\ \hat{\eta}+\frac{\boldsymbol{p}_\Lambda\cdot\hat{\eta}}{m_\Lambda(E_\Lambda+m_\Lambda)}\boldsymbol{p}_\Lambda\right). \qquad (21)$$

We also consider three orthogonal unit vectors along the longitudinal, transverse and normal components of the Λ polarization in the Λ_b rest frame as

$$\hat{e}_L=\frac{\boldsymbol{p}_\Lambda}{|\boldsymbol{p}_\Lambda|},\quad \hat{e}_T=\frac{\boldsymbol{p}_+\times\boldsymbol{p}_\Lambda}{|\boldsymbol{p}_+\times\boldsymbol{p}_\Lambda|},\quad \hat{e}_N=\hat{e}_T\times\hat{e}_L, \qquad (22)$$

where $\boldsymbol{p}_\Lambda$ and $\boldsymbol{p}_+$ are three momenta of the Λ and l^+ in the CM frame of the l^+l^- system. Thus, using these spin vectors one can obtain the differential decay rate for any spin direction $\hat{\eta}$ along the Λ baryon as

$$\frac{\mathrm{d}\Gamma(\hat{\eta})}{\mathrm{d}s}=\frac{1}{2}\left(\frac{\mathrm{d}\Gamma}{\mathrm{d}s}\right)_0\left[1+(P_L\,\hat{e}_L+P_N\,\hat{e}_N+P_T\hat{e}_T)\cdot\hat{\eta}\right], \qquad (23)$$

where P_L, P_N and P_T are functions of s, which give the longitudinal, normal and transverse polarization and $(\mathrm{d}\Gamma/\mathrm{d}s)_0$ is the unpolarized decay width. The polarization components P_i $(i=L,N,T)$ can be obtained from

$$P_i(s)=\frac{\dfrac{\mathrm{d}\Gamma}{\mathrm{d}s}(\hat{\eta}=\hat{e}_i)-\dfrac{\mathrm{d}\Gamma}{\mathrm{d}s}(\hat{\eta}=-\hat{e}_i)}{\dfrac{\mathrm{d}\Gamma}{\mathrm{d}s}(\hat{\eta}=\hat{e}_i)+\dfrac{\mathrm{d}\Gamma}{\mathrm{d}s}(\hat{\eta}=-\hat{e}_i)}. \qquad (24)$$

Thus, one can obtain the polarization components:

$$\begin{aligned}
&P_L(s)\\
&=\frac{16m_{\Lambda_b}^2\sqrt{\lambda}}{\mathcal{K}_0(s)+\mathcal{K}_2(s)/3}\\
&\quad\times\left[8m_l^2m_{\Lambda_b}\right.\\
&\quad\times\left(\mathrm{Re}(D_1^*E_3-D_3^*E_1)+\sqrt{r}\mathrm{Re}(D_1^*D_3-E_1^*E_3)\right)\\
&\quad-4m_l^2m_{\Lambda_b}^2\hat{s}\left(|D_3|^2-|E_3|^2\right)\\
&\quad-4m_{\Lambda_b}(2m_l^2+m_{\Lambda_b}^2\hat{s})\mathrm{Re}(A_1^*B_2-A_2^*B_1)\\
&\quad-\frac{4}{3}m_{\Lambda_b}^3\hat{s}v_l^2\,(3\mathrm{Re}(D_1^*E_2-D_2^*E_1)\\
&\quad+\sqrt{r}\mathrm{Re}(D_1^*D_2-E_1^*E_2))\\
&\quad-\frac{4}{3}m_{\Lambda_b}\sqrt{r}(6m_l^2+m_{\Lambda_b}^2\hat{s}v_l^2)\mathrm{Re}(A_1^*A_2-B_1^*B_2)\\
&\quad-\frac{2}{3}m_{\Lambda_b}^4\hat{s}(2-2r+\hat{s})v_l^2\left(|D_2|^2-|E_2|^2\right)\\
&\quad+\left(4m_l^2+m_{\Lambda_b}^2(1-r+\hat{s})\right)\left(|A_1|^2-|B_1|^2\right)\\
&\quad-\left(4m_l^2-m_{\Lambda_b}^2(1-r+\hat{s})\right)\left(|D_1|^2-|E_1|^2\right)\\
&\quad-\frac{1}{3}m_{\Lambda_b}^2(1-r-\hat{s})v_l^2\left(|A_1|^2-|B_1|^2+|D_1|^2-|E_1|^2\right)\\
&\quad-\frac{1}{3}m_{\Lambda_b}^2\left[12m_l^2(1-r)\right.\\
&\quad\left.+m_{\Lambda_b}^2\hat{s}\left(3(1-r+\hat{s})+v_l^2(1-r-\hat{s})\right)\right]\\
&\quad\left.\times\left(|A_2|^2-|B_2|^2\right)\right], \qquad (25)
\end{aligned}$$

$$\begin{aligned}
&P_N(s) \\
&= \frac{8\pi m_{\Lambda_b}^3 v_l \sqrt{\hat{s}}}{\mathcal{K}_0(s)+\mathcal{K}_2(s)/3} \\
&\quad \times \left[-2m_{\Lambda_b}(1-r+\hat{s})\sqrt{r}\mathrm{Re}(A_1^* D_1 + B_1^* E_1)\right. \\
&\quad +4m_{\Lambda_b}^2 \hat{s}\sqrt{r}\mathrm{Re}(A_1^* E_2 + A_2^* E_1 + B_1^* D_2 + B_2^* D_1) \\
&\quad -2m_{\Lambda_b}^3 \hat{s}\sqrt{r}(1-r+\hat{s})\mathrm{Re}(A_2^* D_2 + B_2^* E_2) \\
&\quad +2m_{\Lambda_b}(1-r-\hat{s}) \\
&\quad \times \left(\mathrm{Re}(A_1^* E_1 + B_1^* D_1) + m_{\Lambda_b}^2 \hat{s}\mathrm{Re}(A_2^* E_2 + B_2^* D_2)\right) \\
&\quad -m_{\Lambda_b}^2\left((1-r)^2 - \hat{s}^2\right) \\
&\quad \left.\times \mathrm{Re}(A_1^* D_2 + A_2^* D_1 + B_1^* E_2 + B_2^* E_1)\right], \qquad (26)
\end{aligned}$$

$$\begin{aligned}
&P_T(s) \\
&= -\frac{8\pi m_{\Lambda_b}^3 v_l \sqrt{\hat{s}\lambda}}{\mathcal{K}_0(s)+\mathcal{K}_2(s)/3} \\
&\quad \times \left[m_{\Lambda_b}^2(1-r+\hat{s})\right. \\
&\quad \times \left(\mathrm{Im}(A_2^* D_1 - A_1^* D_2) - \mathrm{Im}(B_2^* E_1 - B_1^* E_2)\right) \\
&\quad +2m_{\Lambda_b}\left(\mathrm{Im}(A_1^* E_1 - B_1^* D_1)\right. \\
&\quad \left.\left. - m_{\Lambda_b}^2 \hat{s}\mathrm{Im}(A_2^* E_2 - B_2^* D_2)\right)\right]. \qquad (27)
\end{aligned}$$

It should be noted that the longitudinal (P_L) and normal (P_N) polarizations are P-odd and T-even whereas the transverse polarization (P_T) is P-even and T-odd.

3 Contribution from FCNC mediated Z boson

We now consider the effect of the FCNC mediated Z boson on the branching ratios and forward–backward asymmetries of the rare semileptonic Λ_b decays. It is a simple model beyond the standard model with an enlarged matter sector due to an additional vector-like down quark D_4. The presence of an additional down quark implies a 4×4 matrix $V_{i\alpha}$ ($i = u, c, t, 4,\ \alpha = d, s, b, b'$), diagonalizing the down quark mass matrix. For our purpose the relevant information for the low energy physics is encoded in the extended mixing matrix. The charged currents are unchanged except that the V_{CKM} is now the 3×4 upper sub-matrix of V. However, the distinctive feature of this model is that the FCNC interaction enters the neutral current Lagrangian of the left handed down quarks as

$$\begin{aligned}
\mathcal{L}_Z &= \frac{g}{2\cos\theta_{\mathrm{W}}} \qquad (28) \\
&\quad \times \left[\bar{u}_{Li}\gamma^\mu u_{Li} - \bar{d}_{L\alpha}U_{\alpha\beta}\gamma^\mu d_{L\beta} - 2\sin^2\theta_{\mathrm{W}} J_{em}^\mu\right] Z_\mu,
\end{aligned}$$

with

$$U_{\alpha\beta} = \sum_{i=u,c,t} V_{\alpha i}^\dagger V_{i\beta} = \delta_{\alpha\beta} - V_{4\alpha}^* V_{4\beta}, \qquad (29)$$

where U is the neutral current mixing matrix for the down sector, which is given above. As V is not unitary, $U \neq \mathbf{1}$. In particular the non-diagonal elements do not vanish. We have

$$U_{\alpha\beta} = -V_{4\alpha}^* V_{4\beta} \neq 0 \quad \text{for} \quad \alpha \neq \beta. \qquad (30)$$

Since the various $U_{\alpha\beta}$ are non-vanishing, they would signal new physics and the presence of FCNC at the tree level and this can substantially modify the predictions of SM for the FCNC processes.

Thus, in this model the effective Hamiltonian for $b \to sl^+l^-$ is given as

$$\mathcal{H}_{\mathrm{eff}} = \frac{G_{\mathrm{F}}}{\sqrt{2}} U_{sb}[\bar{s}\gamma^\mu(1-\gamma_5)b]\left[\bar{l}(C_V^l \gamma_\mu - C_A^l \gamma_\mu\gamma_5)l\right], \qquad (31)$$

where C_V^l and C_A^l are the vector and axial vector Zl^+l^- couplings, which are given as

$$C_V^l = -\frac{1}{2} + 2\sin^2\theta_{\mathrm{W}}, \qquad C_A^l = -\frac{1}{2}. \qquad (32)$$

Since the structure of the effective Hamiltonian (31) in this model is same as that of the SM, i.e., of the $\sim (V-A)(V-A)$ form, its effect on the various decay parameters can be encoded by replacing the SM Wilson coeffients $(C_9^{\mathrm{eff}})^{\mathrm{SM}}$ and $(C_{10})^{\mathrm{SM}}$ by

$$\begin{aligned}
C_9^{\mathrm{eff}} &= (C_9^{\mathrm{eff}})^{\mathrm{SM}} + \frac{2\pi}{\alpha}\frac{U_{sb}}{V_{tb}V_{ts}^*}, \\
C_{10}^{\mathrm{eff}} &= (C_{10})^{\mathrm{SM}} - \frac{2\pi}{\alpha}\frac{U_{sb}}{V_{tb}V_{ts}^*}. \qquad (33)
\end{aligned}$$

It should be noted that U_{sb} is in general complex; hence it induces the weak phase difference (θ) between the SM and new physics contributions. Since the value of the Wilson coefficients C_9 and C_{10} are opposite to each other as seen from (2), and the new physics contributions to C_9 and C_{10} are opposite to each other, one will get constructive or destructive interference of SM and NP amplitudes for $\theta = \pi$ or zero (where θ denotes the relative weak phase between the SM and NP contribution in the above equation). However, we consider the weak phase difference to be π in our numerical analysis, so as to get constructive interference between the SM and NP amplitudes. The value of $|U_{sb}|$ is found to be

$$|U_{sb}| \simeq 10^{-3}, \qquad (34)$$

which has been extracted from the recent data on $\mathcal{B}(B \to X_S l^+l^-)$ [14].

4 Numerical analysis

For numerical evaluation we use the various particle masses and lifetimes of Λ_b baryon from [15]. The quark masses (in GeV) used are $m_b = 4.6$, $m_c = 1.5$, the CKM matrix elements are taken as $|V_{tb}V_{ts}^*| = 0.041$, $\alpha = 1/128$ and the weak mixing angle $\sin^2\theta_{\mathrm{W}} = 0.23$. For the form factors we use the values calculated in the QCD sum rule [6, 7]

Table 1. Values of the hadronic form factors in the QCD sum rule approach for the $\Lambda_b \to \Lambda$ transition

Form factors	$F(0)$	a_F	b_F
F_1	0.462	−0.0182	-1.76×10^{-4}
F_2	−0.077	−0.0685	1.46×10^{-3}

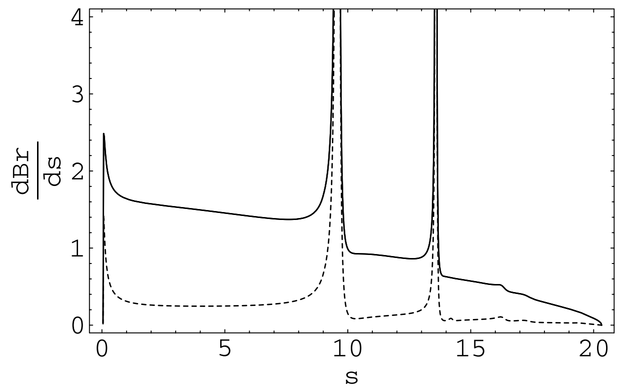

Fig. 1. The differential branching ratio dBr/ds (in units of $10^6\,\mathrm{GeV}^{-2}$) versus s (in GeV^2) for the process $\Lambda_b \to \Lambda\, \mu^+\mu^-$. The solid line denotes the branching ratio including the non-universal Z boson effect whereas the dashed line represents the SM contribution

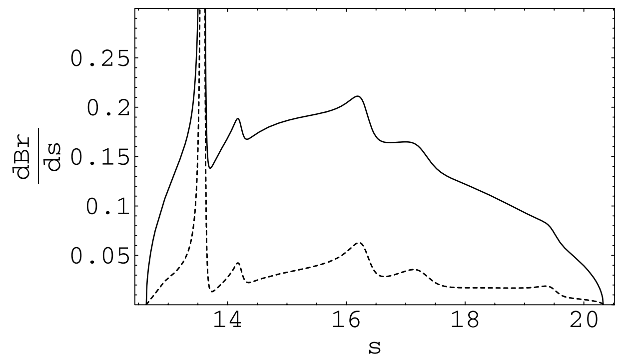

Fig. 2. Same as Fig. 1 for the $\Lambda_b \to \Lambda\tau^+\tau^-$ process

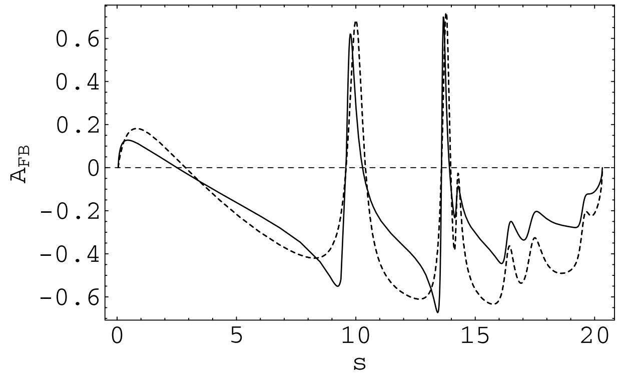

Fig. 3. The forward–backward asymmetry versus s (in GeV^2) for the process $\Lambda_b \to \Lambda\mu^+\mu^-$. The legends are the same as in Fig. 1

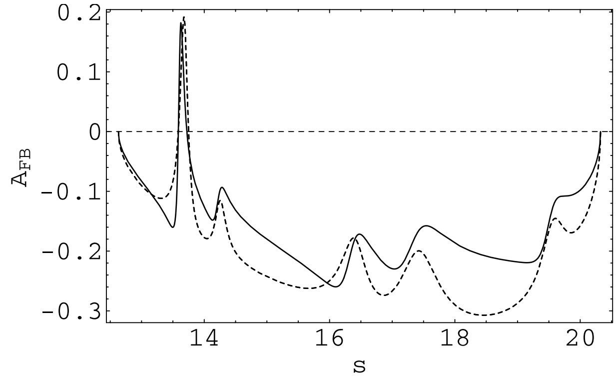

Fig. 4. Same as Fig. 3 for the $\Lambda_b \to \Lambda\tau^+\tau^-$ process

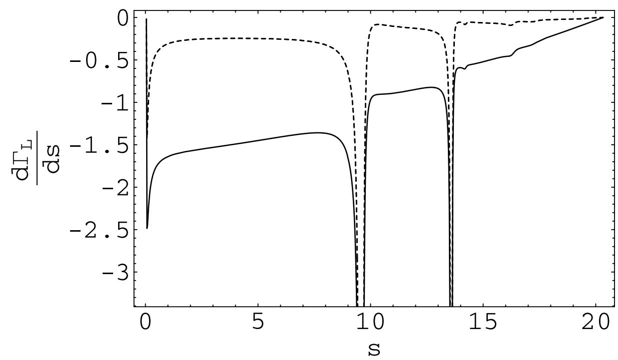

Fig. 5. The differential decay width distribution dΓ_L/ds (in units of $\tau_{\Lambda_b} \times 10^6$ GeV^{-2}) versus s (in GeV^2) of $\Lambda_b \to \Lambda\mu^+\mu^-$ for the longitudinal polarized Λ. The legends are the same as in Fig. 1

approach, where the q^2 dependence of the various form factors are given as

$$F(q^2) = \frac{F(0)}{1 - a_F(q^2/m_{\Lambda_b}^2) + b_F\left(q^2/m_{\Lambda_b}^2\right)^2}\,. \quad (35)$$

The values of the parameters $F_i(0)$, a and b are summarized in Table 1.

With these values we plot the differential decay rate (17) for $\Lambda_b \to \Lambda l^+l^-$ for $l = \mu, \tau$, against s, which are depicted in Figs. 1 and 2. It can be noted from the figures that there is considerable enhancement in the decay rates due to the non-universal Zbs coupling. The forward–backward asymmetries (20) are plotted in Figs. 3 and 4. It can be seen from Fig. 3 that the zero position of A_{FB} shifts towards the left for $\Lambda_b \to \Lambda\mu^+\mu^-$ process due to the NP effect, however there is no such deviation in the $\Lambda_b \to \Lambda\tau^+\tau^-$ process. The differential decay rate for the longitudinal polarized Λ are shown in Figs. 5 and 6. These distributions are very similar to the differential decay rates (Figs. 1 and 2) but with opposite sign. We also find that there is no significant difference in the longitudinal polarization of Λ due to the NP effect. The distribution of the normal polarization components are shown in Figs. 7 and 8. From these figures one can observe that the normal polarization is very small in the region with low momentum transfer and can have significant values in the large momentum transfer region. The transverse polarization P_T is found to be identically zero in this model as the structure of the Hamiltonian is same as that of the SM, i.e., of the $(V - A)(V - A)$ form.

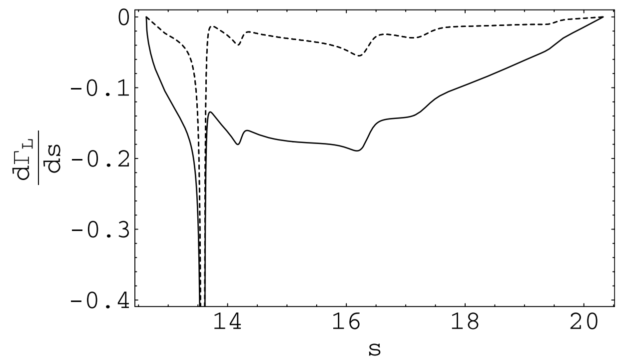

Fig. 6. Same as Fig. 5 for the $\Lambda_b \to \Lambda\tau^+\tau^-$ process

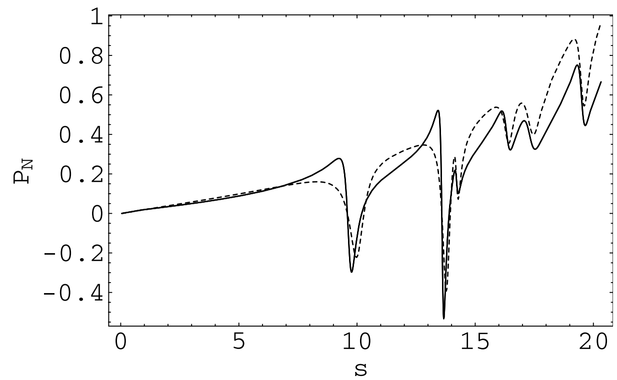

Fig. 7. Normal polarization P_N versus s (in GeV^2) for the $\Lambda_b \to \Lambda\mu^+\mu^-$ process

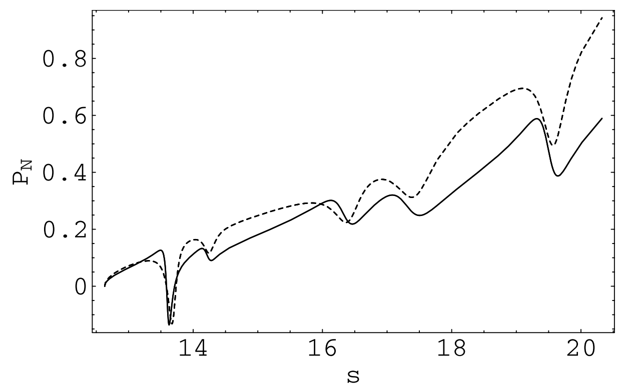

Fig. 8. Same as Fig. 7 for the $\Lambda_b \to \Lambda\tau^+\tau^-$ process

We now proceed to calculate the total decay rates for $\Lambda_b \to \Lambda l^+l^-$ for which it is necessary to eliminate the backgrounds coming from the resonance regions. This can be done by using the following veto windows so that the backgrounds coming from the dominant resonances $\Lambda_b \to \Lambda J/\psi(\psi')$ with $J/\psi(\psi') \to l^+l^-$ can be elimenated:

$$\Lambda_b \to \Lambda\mu^+\mu^- : \; m_{J/\psi} - 0.02 < m_{\mu^+\mu^-} < m_{J/\psi} + 0.02;$$
$$m_{\psi'} - 0.02 < m_{\mu^+\mu^-} < m_{\psi'} + 0.02$$
$$\Lambda_b \to \Lambda\tau^+\tau^- : \; m_{\psi'} - 0.02 < m_{\tau^+\tau^-} < m_{\psi'} + 0.02.$$

Using these veto windows we obtain the branching ratios for semileptonic rare Λ_b decays which are presented in Table 2. It is seen from the table that the branching ratios

Table 2. The branching ratios (in units of 10^{-6}) for various decay processes

Decay modes	$\mathcal{B}^{\mathrm{SM}}$	$\mathcal{B}^{\mathrm{NP}}$
$\Lambda_b \to \Lambda\mu^+\mu^-$	4.55	20.9
$\Lambda_b \to \Lambda\tau^+\tau^-$	0.17	0.92

obtained in the model with the non-universal Z bosons are reasonably enhanced from the corresponding SM values.

5 Conclusion

In this paper we have studied the rare baryonic semileptonic decays $\Lambda_b \to \Lambda l^+l^-$ in the model in which the fermion sector of the SM is extended by an extra vector-like down quark. The importance of this model is that it allows FCNC transitions to occur at the tree level. For the process under consideration, it exhibits the $b \to s$ FCNC transition at the tree level by emitting one Z boson. Here we have studied the effect of this non-universal Zbs couplings on the decay rates and forward–backward asymmetries of $\Lambda_b \to \Lambda\, l^+l^-$ process. Furthermore, we have also studied the differential decay rates when the final Λ baryon is longitudinally polarized and for the normal polarization of the final Λ baryon. We found that in this model the branching ratios of $\Lambda_b \to \Lambda l^+l^-$ can differ significantly from their corresponding SM values. The forward–backward asymmetries are also found to differ from that of the SM expectation due to this non-universal Z mediated FCNC. Moreover, in this model the zero point of the F_{AB} is found to be shifted towards the left. For the polarized Λ, we found that the decay distribution is similar to that of the unpolarized one but with opposite orientation. Furthermore, no significant change in P_N is observed, and P_T is found to be identically zero.

To conclude, we note that the effect of the Z mediated FCNC in the vector-like down quark model can enhance the branching ratio reasonably for the $\Lambda_b \to \Lambda l^+l^-$ and also change the forward–backward asymmetries in these modes. The polarized variables can be studied experimentally to distinguish various new physics models. The transverse polarization (P_T) in $\Lambda_b \to \Lambda l^+l^-$, if found zero (or nonzero) can confirm or (rule out) this new physics scenario.

Acknowledgements. AKG was supported by the National Science Council of the Republic of China under Contract Number NSC-93-2811-M-007-060. The work of RM was partly supported by the Department of Science and Technology, Government of India, through Grant No. SR/FTP/PS-50/2001.

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