Chiral Structure of the $b \to c$ Charged Current

and

the Semileptonic $\Lambda_b$ Decay

Minoru Tanaka

Theory Group, National Laboratory for High Energy Physics (KEK),

Oho 1-1, Tsukuba, Ibaraki 305, Japan

Using the heavy quark effective theory, chiral structure of the $b \to c$ charged current can be determined by the semileptonic $\Lambda_b$ decay with satisfactory theoretical accuracy. We define an asymmetry which is sensitive to the chirality of the $b \to c$ charged current. We show that this asymmetry has no theoretical uncertainty in the heavy quark limit. It is also shown that $1/m_c$ correction to this asymmetry is small and disappears at the kinematical point of zero momentum transfer.
I. INTRODUCTION

The most conspicuous feature of the weak charged current in the standard model is its pure left-handedness. It is implemented by the transformation properties of the chiral fermions: The left-handed fermions are assigned to doublets of the weak $SU(2)$ and the right-handed ones are to singlets.

Another important aspect of the charged current is the mixing. In the standard model, there are three generations of quarks and leptons, and these quarks mix with each other, while the mixing of the leptons has no physical meanings. This quark mixing leads to a $3 \times 3$ matrix (Kobayashi-Maskawa matrix [1]) in the quark charged current, resulting in nine components of the charged current in the quark sector. The quark mixing does not alter the above-mentioned chiral structure of the charged current in the standard model. In other words, the standard model predicts that all nine quark flavor (and three leptonic) components of the weak charged current have the identical chiral structure, pure left-handedness.

This prediction, however, may be altered if extended models are considered. For example, in the Left-Right gauge models based on the $SU(2)_L \times SU(2)_R \times U(1)$ weak gauge group [2], the exchange of the right-handed $W$ boson and the mixing between the left-handed $W$ boson and the right-handed one cause the right-handed component in the charged current. If the model is manifestly Left-Right symmetric, the above flavor universality of the chiral structure of the charged current is maintained although the pure left-handedness is lost.

While, even the flavor universality is not preserved in the non-manifest models because of the difference between the left-handed quark mixing matrix and its right-handed counterpart. In this kind of models, the mass of the right-handed $W$ boson can be relatively light ($\sim$ a few hundred GeV) and non-negligible part of some flavor components of the charged current may be caused by the right-handed interaction [3].

Recently, an extreme model along this line was proposed by Gronau and Wakaizumi [4]. In their model, $b$ quark decay is caused purely by a right-handed current.

One can suppose another example in which both the pure left-handedness and the flavor
universality are broken. It is the standard model with the vector-like fourth generation. This model also contains the right-handed charged current which originates from the SU(2) gauge interaction of the right-handed doublet fermions in the fourth generation. An interesting speculation on this kind of model was made by Fritzsch [5]. Following his speculation, the phenomenon of the parity violation is interpreted as a low energy phenomenon and related to the smallness of the usual quark and lepton masses. The parity violation becomes maximal in the limit of vanishing fermion masses. Therefore, the magnitude of the right-handed charged current is expected to be larger in heavier quark sectors. This means the loss of the flavor universality of the chiral structure of the charged current. Actually, it is shown in ref. [5] that the strength of the right-handed charged current which consists of \( q \) and \( q' \) quarks is proportional to the factor \( \sqrt{m_q m_{q'}} \).

Following the above arguments, the chiral structure of all the nine quark (and three leptonic) charged currents should be examined separately. It, however, is not straightforward to determine the chirality of the quark charged currents experimentally, because we observe hadrons not quarks. We have to extract the quark couplings from the observed hadronic interaction. In spite of this difficulty, the chiral structures of the charged currents which contain only the light quarks are well constrained to be left-handed. Actually, the result of inelastic neutrino interaction [6] suggests \( |g_R/g_L| \lesssim 0.1 \) for the \( d \to u \) charged current, and the PCAC consideration on the \( K \to 2\pi \) and \( K \to 3\pi \) decays gives limits as \( |g_R/g_L| \lesssim 0.004 \) for the \( d \to u \) and \( s \to u \) charged current [7].

On the other hand, since we cannot use the chiral symmetry for heavy quarks (\( c, b \) and \( t \)) in addition to the difficulty of experiments like deep inelastic scattering due to the CKM suppression, there are poorer evidences of the chiralities of the charged currents involving the heavy quarks. In fact, the experimental bound on the \( c \to d, s \) currents is \( |g_R/g_L| \lesssim 0.3 \) [8], and \( |g_R/g_L| \lesssim 0.5 \) for \( b \to c \) current has been given recently by CLEO using \( B \to D^*\ell\nu \) decay [9]. We have no direct limits on the other charged currents involving \( b \) and/or \( t \). It should be noticed that the test made by CLEO assumed the left-handedness of the leptonic charged current, and it cannot discriminate the model proposed in ref. [4] from the standard
model. In this sense, this test is incomplete \cite{1,10}.

From the theoretical point of view, recently a remarkable progress has been made in treating hadrons which contain a quark of much heavier mass than the QCD scale. It was pointed out that an effective theory of QCD with $N_h$ heavy quarks has new symmetry $SU(2N_h)$, associated with the flavor and the spin rotations of the heavy quarks \cite{11}. Such a theory is called as the heavy quark effective theory. This symmetry, especially its spin rotation part, is expected to play a crucial role to determine the chiral structure of the charged current involving the heavy quarks.

In this paper, we discuss a method to detect the effect of the right-handed $b \rightarrow c$ current in the semileptonic decay of $\Lambda_b$, which is supposed to be the lightest bottom baryon, into $\Lambda_c$. In the following analysis, we utilize the heavy quark effective theory where we regard both $b$ and $c$ quarks as heavy. The bottom decay is caused by this current for the most part and its detailed study may lead us to something beyond the standard model.

Moreover, we stress that the $\Lambda_{b,c}$ baryons are simpler systems than the $B^{(*)}$ and $D^{(*)}$ mesons in a standpoint of the heavy quark effective theory, because the light degrees of freedom form a zero-spin system. The simplicity of these baryons leads to the most important result of our analysis that the theoretical uncertainty is small enough to see the chirality of the $b \rightarrow c$ current. The $1/m_c$ correction is known to be controlled by one dimensionful parameter \cite{12}, and we found that, as for the asymmetry which we will define in the following section, even the effect of $1/m_c$ correction through this parameter vanishes at the kinematical point of zero momentum transfer.

In section 2, we present our formalism of the differential decay rate, and define an asymmetry which is sensitive to the chirality of the $b \rightarrow c$ current. Section 3 includes the implication of the heavy quark limit and the numerical result in this limit. In section 4, the $1/m_c$ correction is discussed. We state two remarks and our conclusion in section 5.
II. FORMALISM

In this section, we present an expression of a double differential decay rate of $\Lambda_b \to \Lambda_c \ell \bar{\nu}$ decay and define an asymmetry which is a good probe of the chirality of the $b \to c$ charged current. Here we assume a polarization of the initial $\Lambda_b$ and sum up with the final state spins. The initial polarization is realized in $e^+e^- \to Z \to b\bar{b}$ followed by the hadronization $b \to \Lambda_b$ [13]. Note that the semileptonic $\Lambda_b$ decay in this process has already been observed at LEP [14]. We work with this production process of $\Lambda_b$ in mind, and therefore we choose the frame in which the initial $\Lambda_b$ is running [15].

We start with the definition of form factors which can appear in $\Lambda_b \to \Lambda_c$ transition by weak vector and axial-vector currents:

$$\langle \Lambda_c(v', s')| \bar{c}\gamma_\mu b | \Lambda_b(v, s) \rangle \equiv \bar{u}_{\Lambda_c}(v', s')(F_1\gamma_\mu + F_2v_\mu + F_3v'_\mu)u_{\Lambda_b}(v, s),$$  \hspace{1cm} (1)

$$\langle \Lambda_c(v', s')| \bar{c}\gamma_\mu \gamma_5 b | \Lambda_b(v, s) \rangle \equiv \bar{u}_{\Lambda_c}(v', s')(G_1\gamma_\mu + G_2v_\mu + G_3v'_\mu)\gamma_5u_{\Lambda_b}(v, s),$$  \hspace{1cm} (2)

where $v = p_b/m_{\Lambda_b}$ and $v' = p_c/m_{\Lambda_c}$ are the four-velocity of the $\Lambda_b$ and $\Lambda_c$ respectively, and $F_i$’s and $G_i$’s are functions of $v \cdot v'$.

Here, we consider the differential decay rate of $\Lambda_b \to \Lambda_c \ell \bar{\nu}$ with respect to the energy of $\Lambda_c$ ($E_c$) and the momentum transfer squared ($q^2 = (p_b - p_c)^2$). One can write this double differential rate as

$$\frac{d\Gamma}{dx_c dq^2} = J(q^2) + PK(q^2)(x_c - \bar{x}_c),$$  \hspace{1cm} (3)

where $x_c = E_c/E_b$, $\bar{x}_c = p_b \cdot p_c/m_{\Lambda_b}^2$, $E_b$ is the energy of the initial $\Lambda_b$, and $P$ is the initial $\Lambda_b$ polarization which is, as will be explained, equal to that of the initial $b$ quark in the limit of infinite $m_b$. The polarization of $b$ quark which comes from the above mentioned process $Z \to b\bar{b}$ is given by

$$P = \frac{2G_V G_A}{G_V^2 + G_A^2} \simeq -0.93,$$  \hspace{1cm} (4)
where $G_V$ and $G_A$ are the vector and the axial-vector coupling constants in $Zb\bar{b}$ vertex respectively, and we used $\sin^2 \theta_W = 0.233$. Note that the functions $J(q^2)$ and $K(q^2)$ are independent of $x_c$. The range of $x_c$ is given by $x_c^{\text{min}} \leq x_c \leq x_c^{\text{max}}$ with

$$x_c^{\text{max, min}} = \bar{x}_c \pm \frac{\lambda}{2}, \quad \lambda = \beta \frac{\sqrt{w_+ w_-}}{m_{\Lambda_b}^2},$$

(5)

where $\beta = |p_b|/E_b$ is the rapidity of $\Lambda_b$ and $w_\pm = (m_{\Lambda_b} \pm m_{\Lambda_c})^2 - q^2$.

$J(q^2)$ and $K(q^2)$ are written in terms of the couplings of the $b \to c$ vector and axial-vector current (denoted by $g_V$ and $g_A$ respectively [16]), and the form factors $F_i$’s and $G_i$’s. Their explicit forms are given in the appendix. Note that the leptonic couplings $f_V$ and $f_A$ appear in an overall factor, since we neglected the lepton mass and integrated the lepton system completely.

A genuine parity violating effect in the $b \to c$ charged current is expected to be observed through the correlation between the spin of $\Lambda_b$ and the momentum of $\Lambda_c$. Since the second term in eq. (3) expresses this correlation, an energy asymmetry which picks it up is defined as

$$A(q^2) \equiv \frac{\int_{x_c^{\text{max}}}^{x_c^{\text{max}}} \frac{d\Gamma}{dx_c dq^2} dx_c - \int_{x_c^{\text{min}}}^{x_c^{\text{min}}} \frac{d\Gamma}{dx_c dq^2} dx_c}{\int_{x_c^{\text{max}}}^{x_c^{\text{max}}} \frac{d\Gamma}{dx_c dq^2} dx_c + \int_{x_c^{\text{min}}}^{x_c^{\text{min}}} \frac{d\Gamma}{dx_c dq^2} dx_c}$$

(6)

$$= \frac{\lambda}{4} \bar{p} \frac{K(q^2)}{J(q^2)}. \quad (7)$$

Note that we do not integrate with $q^2$ in the r. h. s. of eq. (6) because $J(q^2)$ and $K(q^2)$ involve the unknown form factors in eqs. (1) and (2). As explained below, however, $A(q^2)$ itself involves no unknown functions in the heavy quark limit, because all form factors are proportional to an unknown function in this limit [17]. Moreover, it is worth while pointing out that $A(q^2)$ is independent of the leptonic couplings. Therefore we can evaluate the above asymmetry without any theoretical uncertainty in the heavy quark limit.
III. IMPLICATION OF THE HEAVY QUARK LIMIT

Here, we discuss the implication of the heavy quark limit on the asymmetry $A(q^2)$. In general, we cannot evaluate $A(q^2)$ because it depends on the unknown form factors defined in eqs. (1) and (2). This difficulty, however, greatly reduces in the heavy quark limit.

In the limit that $m_b, m_c \to \infty$, the six form factors in eqs. (1) and (2) can be written by only one unknown function of $q^2$ [17]:

$$F_1 = G_1 \equiv \zeta(q^2), \quad F_2 = F_3 = G_2 = G_3 = 0.$$  \hspace{1cm} (8)

As seen in the above equation (8), the spin of heavy $\Lambda_Q$ hyperon is carried by the heavy quark. This is also true for $Z \to b\bar{b}$ followed by $b \to \Lambda_b$. Therefore, we can use eq. (4) as the polarization of the $\Lambda_b$ which comes from $Z$ decay in the heavy $b$ quark limit.

Using eq. (8), (23) and (24), we get

$$J(q^2) = \frac{|f_V|^2 + |f_A|^2}{192\pi^3 \beta E_b} \zeta(q^2)^2 \left[|g_V|^2(w_+w_- + 3q^2w_-) + |g_A|^2(w_+w_- + 3q^2w_+)\right],$$  \hspace{1cm} (9)

$$K(q^2) = \frac{|f_V|^2 + |f_A|^2}{96\pi^3 \beta^2 E_b} \zeta(q^2)^2 (g_V g_A^* + g_A^* g_V) m_{\Lambda_b}^2 (m_{\Lambda_b}^2 - m_{\Lambda_c}^2 - 2q^2).$$  \hspace{1cm} (10)

Then, according to eq. (7), we get an expression for $A(q^2)$ in the heavy quark limit:

$$A(q^2) = -\mathcal{P} \frac{(|g_L|^2 - |g_R|^2)(m_{\Lambda_b}^2 - m_{\Lambda_c}^2 - 2q^2)\sqrt{w_+w_-}}{(|g_L|^2 + |g_R|^2) \{2w_+w_- + 3q^2(w_+ + w_-)\} - 3(g_L g_R^* + g_R g_L^*)q^2(w_+ - w_-)},$$  \hspace{1cm} (11)

where $g_L = g_V - g_A$, $g_R = g_V + g_A$. As is mentioned above, this expression has no theoretical uncertainty if $g_R/g_L$ is given. The numerical result of this expression is shown in fig. 1. Throughout this paper, we use $m_{\Lambda_c} = 2.285\text{GeV}$ [18] and $m_{\Lambda_b} = 5.640\text{GeV}$ [19] in numerical calculations.

In fig. 1, $A(q^2)$ is plotted for pure left-handed case ($g_R/g_L = 0$), 30% right-handed contamination cases ($g_R/g_L = \pm 0.3$) and pure right-handed case ($g_L/g_R = 0$). In the cases of 30% right-handed contamination, it is assumed that the left-handed coupling and the right-handed one are relatively real and the two relative sign possibilities are taken into...
account. Fig. 1 shows that we can easily discriminate between the pure left-handed case and the pure right-handed case by measuring $A(q^2)$ for smaller $q^2$. An experiment with enough accuracy will give an upper bound for the strength of right-handed current, or will find a right-handed current.

From eq. (11), we can immediately find that $A(q^2)$ takes the following values at the kinematical points of zero momentum transfer and zero recoil:

$$A(0) = -\frac{1}{2} \mathcal{P} \frac{|g_L|^2 - |g_R|^2}{|g_L|^2 + |g_R|^2},$$

$$A((m_{\Lambda_b} - m_{\Lambda_c})^2) = 0.$$  

(12)

At this stage, we conclude that we can clearly distinguish the left-handed case from the right-handed case by measuring the asymmetry $A(q^2)$ near the kinematical point of zero momentum transfer, at least in the heavy quark limit. It should be noticed that the phase space volume is finite at the zero momentum transfer, while it vanishes at the point of zero recoil.

**IV. 1/m$_c$ CORRECTION**

In this section, we discuss the $1/m_c$ correction to the result of the previous section. The charm quark mass is not so heavy compared with a typical QCD scale that the $1/m_c$ correction is expected to be the dominant one, and one should take it into account.

Including the leading $1/m_c$ correction, eq. (8) is modified as follows [12]:

$$F_1 = (1 + \Delta) \zeta(q^2), \quad G_1 = \zeta(q^2), \quad F_2 = G_2 = -\Delta \zeta(q^2), \quad F_3 = G_3 = 0,$$

(14)

where

$$\Delta = \frac{\bar{\Lambda}}{m_c} \left(\frac{1}{1 + v \cdot v'}\right),$$

(15)

and $\bar{\Lambda} = m_{\Lambda_b} - m_b = m_{\Lambda_c} - m_c$. $\bar{\Lambda}$ is an unknown parameter and is estimated as $\bar{\Lambda} \sim 0.7$GeV. Note that the form factors are still proportional to one unknown function $\zeta(q^2)$. This means
that the asymmetry \( A(q^2) \) does not involve any unknown functions even with the \( 1/m_c \) correction although it depends on the unknown parameter \( \bar{\Lambda} \).

Using eq. (14), one can write down an analytic formula for \( A(q^2) \) in terms of \( \bar{\Lambda} \). It, however, is too lengthy to be presented here. We show only the numerical result in fig. 2.

In fig. 2, \( A(q^2) \) for the pure left-handed case and the pure right-handed case are plotted with and without \( 1/m_c \) correction. (\( \bar{\Lambda} = 0.7 \)GeV and \( \bar{\Lambda} = 0 \) respectively.) This figure shows that the \( 1/m_c \) correction decreases quickly with approaching to zero momentum transfer and vanishes at the point of zero momentum transfer.

Actually, it is easy to demonstrate the latter analytically. At the point \( q^2 = 0 \), we get the same result as eq. (12) even in the presence of \( 1/m_c \) correction. Using eqs. (23), (24) and (14), one gets

\[
J(0) = \frac{|f_V|^2 + |f_A|^2}{384\pi^3\beta E_b} \zeta(0)^2 (|g_L|^2 + |g_R|^2)(m_{\Lambda_b}^2 - m_{\Lambda_c}^2)^2 \left\{ 1 + \Delta_0 \left( 1 - m_{\Lambda_c}/m_{\Lambda_b} \right) \right\} , \quad (16)
\]

\[
K(0) = \frac{|f_V|^2 + |f_A|^2}{192\pi^3\beta^2 E_b} \zeta(0)^2 (|g_L|^2 - |g_R|^2)m_{\Lambda_b}^2(m_{\Lambda_b}^2 - m_{\Lambda_c}^2) \left\{ 1 + \Delta_0 \left( 1 - m_{\Lambda_c}/m_{\Lambda_b} \right) \right\} , \quad (17)
\]

where \( \Delta_0 = \Delta|_{q^2=0} \) and we neglected the terms of higher order in \( \Delta_0 \). Using eqs. (9), (16) and (17), one finds the same result of eq. (12) \[20\].

From the above argument, we conclude that the \( 1/m_c \) correction can be negligible if one does not go near the zero recoil point \( (q^2 = (m_{\Lambda_b} - m_{\Lambda_c})^2) \). Especially, it vanishes at the point of zero momentum transfer \( (q^2 = 0) \).

V. CONCLUSION

Before summarizing our result, two remarks are in order:

(i) **Effect of fragmentation**: One needs to know the momentum of the initial \( \Lambda_b \) to measure the asymmetry \( A(q^2) \). If the initial \( b \) quark hadronizes into a \( \Lambda_b \) without other hadrons, the absolute value of the \( \Lambda_b \) momentum is simply given by \( \sqrt{(m_Z/2)^2 - m_{\Lambda_b}^2} \). In general, however, \( b \) quarks in \( Z \) decay hadronize with several hadrons, *i.e.* they form
jets. Then the magnitude of the $\Lambda_b$ momentum is no longer a constant. It varies according to some fragmentation function $D(z)$ where $z$ is the energy or momentum fraction carried by $\Lambda_b$ \cite{21}. This affects a measurement of the double differential rate of eq. (3).

To discuss the fragmentation effect, we consider the model of Peterson et al. \cite{22}. In this model, the fragmentation function is given by

\[ D(z) = Nz^{-1} \left( 1 - \frac{1}{z} - \frac{\epsilon}{1-z} \right)^{-2}, \quad (18) \]

where $N$ is a normalization factor and $\epsilon$ is a parameter. We can write $\epsilon$ as $\epsilon \simeq m_q^2/m_Q^2$, where $m_q$ is the mass of light quark and $m_Q$ is that of the heavy quark. Eq. (18) describes the experimental data of $\Lambda_c$ production well \cite{23}, and the scaling of $\epsilon$ as $m_Q^{-2}$ has been observed by comparing charm and bottom fragmentation \cite{24}. Therefore, the use of this model for a qualitative discussion seems to be legitimate.

The maximum value of $D(z)$ in eq. (18) is given at

\[ z_{max} = 1 + \frac{\epsilon}{2} - \frac{1}{2} \sqrt{\epsilon(\epsilon + 4)} \]
\[ = 1 - \sqrt{\epsilon} + \cdots. \quad (19) \]

\[ = 1 - \sqrt{\epsilon} + \cdots. \quad (20) \]

Therefore the fragmentation effect seems to be $1/m_b$ effect in $\Lambda_b$ production. Since we ignored $1/m_b$ effect in the arguments of the previous sections, the fragmentation effect can be neglected, at least in the formal point of view.

While the average of $z$ which is distributed according to the fragmentation function in eq. (18) cannot be expanded in $\sqrt{\epsilon}$:

\[ \langle z \rangle = 1 + \frac{2}{\pi} \sqrt{\epsilon \log \epsilon} + \cdots. \quad (21) \]

This equation suggests that the effect of fragmentation is not so small. To be more precise, Monte Carlo study seems to be needed to extract the asymmetry in eq. (3) from experimental data. This is beyond the scope of the present paper.
(ii) $\Lambda_b$ from polarized $e^+e^-$ collision near threshold: To get rid of the fragmentation effect, one may produce $\Lambda_b$ in $e^+e^-$ collision near $\Lambda_b\bar{\Lambda}_b$ threshold. In this case, however, polarized beams are needed to get the necessary polarization of $\Lambda_b$. In the collision of right-handed electron and left-handed positron, the polarization of $b$ quark is given by

$$P = \frac{2s \cos \theta}{s(1 + \cos^2 \theta) + 4m_b^2(1 - \cos^2 \theta)},$$

(22)

where $s$ is the center-of-mass energy squared, and $\theta$ denotes the angle between the direction of the incoming electron momentum and the that of the outgoing $b$ quark. One can use eq. (22) for $\Lambda_b$ polarization if $1/m_b$ correction is neglected.

If one approaches close to the threshold, the produced $\Lambda_b$ is almost stopped. In the rest frame of $\Lambda_b$, the asymmetry defined in eq. (3) has no meanings. However, we can observe the same physical effect, i.e. the spin-momentum correlation, by measuring the angular distribution of $\Lambda_c$.

To summarize, we discussed $\Lambda_b \to \Lambda_c \ell \bar{\nu}$ decay with an initial polarization using the heavy quark effective theory. We defined an asymmetry which was sensitive to the chirality of the $b \to c$ charged current and independent of that of leptonic charged currents. In the heavy quark limit, this asymmetry has no theoretical uncertainty. The $1/m_c$ correction to it is negligibly small if one does not go near the kinematical point of zero recoil. Moreover, the $1/m_c$ correction vanishes at the point of zero momentum transfer.

According to the above result, we conclude that the investigation of the semileptonic $\Lambda_b$ decay with an initial polarization will provide a good test of the left-handedness of the $b \to c$ current and lead to a limit or an evidence on existence of a right-handed $b \to c$ current with satisfactory theoretical accuracy.

ACKNOWLEDGMENTS

The author would like to thank Prof. Hikasa, Prof. Körner and Prof. Wakaizumi for useful discussions.
**EXPLICIT FORMS OF $J(q^2)$ AND $K(q^2)$**

Here, we give the explicit forms of $J(q^2)$ and $K(q^2)$ which appear in eq. (3).

\[
J(q^2) = \frac{|f_V|^2 + |f_A|^2}{192\pi^3\beta E_b}
\times \left[ |g_V|^2 \left( w_+ w_- \left\{ F_1^2 + 2(m_{\Lambda_b} + m_{\Lambda_c})F_1 F_+ + w_+ F_+^2 \right\} + 3q^2 w_- F_1^2 \right) + |g_A|^2 \left( w_+ w_- \left\{ G_1^2 - 2(m_{\Lambda_b} - m_{\Lambda_c})G_1 G_+ + w_+ G_+^2 \right\} + 3q^2 w_- G_1^2 \right) \right],
\]

(23)

\[
K(q^2) = \frac{|f_V|^2 + |f_A|^2}{96\pi^3\beta^2 E_b} (g_V g_A^* + g_V^* g_A) m_{\Lambda_b}^2
\times \left[ (m_{\Lambda_b}^2 - m_{\Lambda_c}^2 - 2q^2)F_1 G_1 - (m_{\Lambda_b} + m_{\Lambda_c})w_- F_1 G_+ \\
+ (m_{\Lambda_b} - m_{\Lambda_c})w_+ F_+ G_1 - w_+ w_- F_+ G_+ \right],
\]

(24)

where $f_V$ and $f_A$ denote the leptonic vector and axial-vector couplings, $F_+ = (F_2/m_{\Lambda_b} + F_3/m_{\Lambda_c})/2$ and $G_+ = (G_2/m_{\Lambda_b} + G_3/m_{\Lambda_c})/2$. Note that only $F_1$, $G_1$ and the combinations $F_+$ and $G_+$ survive in the above expressions because of vanishing lepton masses.
REFERENCES

* A JSPS fellow.

† E-mail address: minoru@theory.kek.jp

[1] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. 49, 652 (1973).

[2] R. N. Mohapatra and J. C. Pati, Phys. Rev. D11, 566, 2558 (1975);
     R. N. Mohapatra and G. Senjanovic, Phys. Rev. D12, 1502 (1975).

[3] F. I. Olness and M. E. Ebel, Phys. Rev. D30, 1034 (1984);
     P. Langacker and S. Uma Sankar, Phys. Rev. D40, 1569 (1989);
     D. London and D. Wyler, Phys. Lett. B232, 503 (1989);
     H. Nishiura, E. Takasugi and M. Tanaka, Prog. Theor. Phys. 84, 116 (1990); 85, 343 (1991).

[4] M. Gronau and S. Wakaizumi, Phys. Rev. Lett. 68, 1814 (1992);
     M. Gronau and S. Wakaizumi, in B decays, ed. S. L. Stone
     (World Scientific,Singapore,1992) p. 479.

[5] H. Fritzsch, preprint MPI-PAE/PTh 19/91 (1991).

[6] H. Abramowicz et al., Z. Phys. C12, 225 (1982).

[7] J. F. Donoghue and B. R. Holstein, Phys. Lett. 113B, 382 (1982).

[8] H. Abramowicz et al., Z. Phys. C15, 19 (1982).

[9] CLEO Collaboration, S. Sanghera et al., preprint CLNS 92/1156 (1992).

[10] M. Gronau and S. Wakaizumi, Phys. Lett. B280, 79 (1992).

[11] N. Isgur and M. B. Wise, Phys. Lett. B232, 113 (1989); B237 527 (1990).

[12] H. Georgi, B. Grinstein and M. B. Wise, Phys. Lett. B252, 456 (1990).

[13] T. Mannel and G. A. Schuler, Phys. Lett. B279, 184 (1992).
[14] ALEPH Collaboration, D. Decamp et al., Phys. Lett. B278, 209 (1992); preprint CERN-PPE/92-73;
OPAL Collaboration, P. D. Acton et al., Phys. Lett. B281, 394 (1992).

[15] Recently there appeared papers which discussed the lepton energy spectrum in this
process to test the chirality of $b \to c$ coupling. They regarded $\Lambda_b$ decay as free quark one.
J. F. Amundson et al., preprint CALT-68-1804, EFI 92-38-REV (1992);
M. Gronau and S. Wakaizumi, preprint SLAC-PUB-5892, CERN-TH.6620/92.

[16] Some trivial factors like $G_F$ may be absorbed into these couplings.

[17] N. Isgur and M. B. Wise, Nucl. Phys. B348, 276 (1991);
T. Mannel, W. Roberts and Z. Ryzak, Nucl. Phys. B355, 38 (1991).

[18] Particle Data Group, Phys. Rev. D45, S1 (1992).

[19] UA1 Collaboration, C. Albajar et al., Phys. Lett. B273, 540 (1991).

[20] A similar stability against the $1/m$ correction has been pointed out by Körner and
Krämer in a different context.
J. G. Körner and M. Krämer, Phys. Lett. B275, 495 (1992).

[21] Precise definitions of $z$ depend on authors.

[22] C. Peterson et al., Phys. Rev. D27, 105 (1983).

[23] ARGUS Collaboration, H. Albrecht et al., Phys. Lett. B207, 109 (1988);
T. Bowcock et al., Phys. Rev. Lett. 55, 923 (1985);
CLEO Collaboration, P. Avery et al., Phys. Rev. D43, 3599 (1991).

[24] J. Chrin, Z. Phys. C36, 163 (1987).
FIGURES

FIG. 1. $A(q^2)$ in the heavy quark limit for pure left-handed case (solid line), 30% right-handed contamination cases (dashed and dash-dotted lines), and pure right-handed case (dotted line).

FIG. 2. $A(q^2)$ for pure left-handed case with (dashed line) and without (solid line) $1/m_c$ correction, and that for pure right-handed case with (dash-dotted line) and without (dotted line) $1/m_c$ correction.
$A(q^2)$

Fig. 1
Fig. 2