Chiral structure of the $b \to c$ charged current and semileptonic Λ_b decay

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Using the heavy quark effective theory, the chiral structure of the $b \to c$ charged current can be determined by the semileptonic Λ_b decay with small theoretical uncertainties. 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 the $1/m_c$ correction to this asymmetry is small and disappears at the kinematical point of zero momentum transfer.

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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 be doublets of the weak SU(2) and the right-handed ones to be 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 leptons has no physical meaning. This quark mixing leads to a 3×3 matrix (Kobayashi-Maskawa matrix [1]) in the quark sector, 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, i.e., pure left handedness.

This prediction, however, may be altered if some 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 induce the right-handed components in the charged current. If the model is manifestly left-right symmetric, the above-mentioned flavor universality of the chiral structure of the charged current is maintained although the pure left handedness is lost.

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

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

One can conceive another example in which both the pure left handedness and the flavor universality are broken. It is the standard model with a vector like fourth generation. This model also contains a right-handed charged current which originates from the right-handed doublet fermions in the fourth generation through the SU(2) gauge interaction. An interesting speculation on this kind of model was made by Fritzsch [5]. According to his speculation, parity violation is interpreted as a low-energy phenomenon and related to the smallness of the usual quark and lepton masses. 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 a heavier quark sector. This means a loss of the flavor universality in 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 $\sqrt{m_q m_{q'}}$.

The above arguments suggest that the chiral structures of all the nine quark (and three leptonic) charged currents should be examined separately. However, it is not straightforward to determine the chirality of a *quark* charged current experimentally, because we observe *hadrons* and not quarks. We have to extract the information on the quark couplings from the observed hadronic interactions. 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 experimental data of inelastic neutrino interaction [6] suggests $|g_R/g_L| \lesssim 0.1$ for the $d \to u$ charged current, and the partially conserved axial-vector current (PCAC) consideration on the $K \to 2\pi$ and $K \to 3\pi$ decays gives $|g_R/g_L| \lesssim 0.004$ for the $d \to u$ and $s \to u$ charged currents [7].

On the other hand, because of the lack of chiral symmetry for heavy quarks (c, b, and t) and the difficulty of experiments such as deep inelastic scattering due to the

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Cabibbo-Kobayashi-Maskawa (CKM) suppression, there is poorer evidence of the chiralities of the charged currents involving 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 the $b \to c$ 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 noted 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 [4, 10].

From the theoretical point of view, recently remarkable progress has been made in treating a hadron which contains a quark with much heavier mass than the QCD scale. It was pointed out that the effective theory of QCD with N_h heavy quarks has a new symmetry $\mathrm{SU}(2N_h)$, associated with the flavor and the spin rotations of heavy quarks [11]. This effective theory is called the heavy quark effective theory. This symmetry, especially its spin rotation part, is expected to play a crucial role in determining 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 $b \to c$ current in the semileptonic decay of Λ_b into Λ_c. Λ_b is supposed to be the lightest bottom baryon. 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 from the 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 $b \to c$ current. The $1/m_c$ correction is known to be controlled by one dimensionful parameter [12], and we found that as for the asymmetry, which we will define in the following section, even the effect of the $1/m_c$ correction through this parameter vanishes at the kinematical point of *zero* momentum transfer.

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

II. FORMALISM

In this section, we present an expression of the 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 that the initial Λ_b is polarized and sum over 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 Λ_b decay in this process has already been observed at the CERN e^+e^- collider LEP [14]. We proceed with this production process of Λ_b in mind, and therefore we choose the frame in which the initial Λ_b is running.

Recently there appeared papers which discussed the lepton energy spectrum in this process to test the chirality of $b \to c$ coupling. They regarded Λ_b decay as a free quark decay [15]. We expect, however, dependence on the momentum transfer squared (denoted by q^2) in $\Lambda_b \to \Lambda_c$ transition form factors. If one wishes to handle a quantity which involves q^2 integration, he or she has to assume certain form factors [16]. This leads to undesirable theoretical uncertainties. Hence, we fix q^2.

Moreover, as is stated in the second paper of Ref. [15], the lepton energy distribution is affected by the chiral structure of the leptonic current. This makes things complicated. Since we are interested in the chirality of the quark current in this paper, we integrate over the lepton phase space.

It will be shown in the following arguments how these procedures will make things independent of the hadronic form factors and the lepton couplings.

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

$$\langle \Lambda_c(v', s')|\, \bar{c}\gamma_\mu b\, |\Lambda_b(v, s)\rangle$$

$$\equiv \bar{u}_{\Lambda_c}(v', s')(F_1\gamma_\mu + F_2 v_\mu + F_3 v'_\mu) u_{\Lambda_b}(v, s)\,, \quad (1)$$

$$\langle \Lambda_c(v', s')|\, \bar{c}\gamma_\mu \gamma_5 b\, |\Lambda_b(v, s)\rangle$$

$$\equiv \bar{u}_{\Lambda_c}(v', s')(G_1\gamma_\mu + G_2 v_\mu + G_3 v'_\mu)\gamma_5 u_{\Lambda_b}(v, s)\,, \quad (2)$$

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

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

$$\frac{d\Gamma}{dx_c dq^2} = J(q^2) + \mathcal{P} K(q^2)(x_c - \bar{x}_c)\,, \quad (3)$$

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 Λ_b, and $\mathcal{P}$ is the initial Λ_b polarization, which is, as will be explained, equal to that of the initial b quark in the infinite m_b limit. The polarization of b quarks which come from the above mentioned process $Z \to b\bar{b}$ is given by

$$\mathcal{P} = \frac{2G_V G_A}{G_V^2 + G_A^2} \simeq -0.93\,, \quad (4)$$

where G_V and G_A are the vector and the axial-vector coupling constants in $Zb\bar{b}$ vertex, respectively. We used $\sin^2\theta_W = 0.233$, assuming the standard structure of the neutral current, for a numerical illustration.

In the models such as the left-right model, the structure of the neutral current is modified in general. Of course, one can determine the b quark polarization without relying on the theory of the electroweak interaction by measuring G_A/G_V through the forward-backward asymmetries. The magnitude of the modification, however, must be small and the value in Eq. (4) is plausible because of the recent precise measurement on the Z^0 pole. In fact, the weak mixing angle determined by the $b\bar{b}$ forward-backward asymmetry is consistent with the values determined by other measurements. Rigorously speaking, the value in Eq. (4) has some experimental error [17].

Note that the functions $J(q^2)$ and $K(q^2)$ are independent of x_c. The range of x_c is given by $x_c^{\min} \le x_c \le x_c^{\max}$ with

$$x_c^{\max,\min} = \bar{x}_c \pm \frac{\lambda}{2}\,, \quad \lambda = \beta\frac{\sqrt{w_+w_-}}{m_{\Lambda_b}^2}\,, \tag{5}$$

where $\beta = |\boldsymbol{p}_b|/E_b$ is the rapidity of Λ_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 currents (denoted by g_V and g_A, respectively [18]) 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 the overall factor, since we neglected the lepton mass and integrated out 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 Λ_b and the momentum of Λ_c. Since the second term in Eq. (3) expresses this correlation, the energy asymmetry which picks it up is defined as

$$A(q^2) \equiv \frac{\displaystyle\int_{\bar{x}_c}^{x_c^{\max}} \frac{d\Gamma}{dx_c dq^2}dx_c - \int_{x_c^{\min}}^{\bar{x}_c} \frac{d\Gamma}{dx_c dq^2}dx_c}{\displaystyle\int_{\bar{x}_c}^{x_c^{\max}} \frac{d\Gamma}{dx_c dq^2}dx_c + \int_{x_c^{\min}}^{\bar{x}_c} \frac{d\Gamma}{dx_c dq^2}dx_c} \tag{6}$$

$$= \frac{\lambda}{4}\mathcal{P}\frac{K(q^2)}{J(q^2)}\,. \tag{7}$$

Note that we do not integrate over q^2 on the rhs of Eq. (6) because $J(q^2)$ and $K(q^2)$ involve the unknown form factors in Eqs. (1) and (2), as was stated at the beginning of this section. 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 [19]. Moreover, as was mentioned above, $A(q^2)$ is independent of the leptonic couplings. Therefore we can evaluate the above asymmetry without any theoretical uncertainties 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, is greatly reduced 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 using only one unknown function $\zeta(q^2)$ [19]:

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

As can be seen in the above Eq. (8), the spin of heavy Λ_Q hyperon is carried by the heavy quark. This can be applied for the Λ_b's which come from the $Z \to b\bar{b}$ decays. Therefore, we can use Eq. (4) as the polarization of the Λ_b's which come from the Z decays in the heavy b quark limit.

Using Eqs. (8), (A1), and (A2), we get

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

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

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

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

where $g_L = g_V - g_A$, $g_R = g_V + g_A$. As was mentioned above, this expression has no theoretical uncertainty if g_R/g_L is given. The numerical values of this expression are shown in Fig. 1. Throughout this paper, we use $m_{\Lambda_c} = 2.285$ GeV [20] and $m_{\Lambda_b} = 5.640$ GeV [21] in the numerical calculations.

In Fig. 1, $A(q^2)$ is plotted for the pure left-handed case ($g_R/g_L = 0$), the 30% right-handed contamination case ($g_R/g_L = \pm 0.3$), and the pure right-handed case ($g_L/g_R = 0$). In the case of 30% right-handed contamination, it is assumed that the left-handed coupling and the right-handed one are relatively real and the two possibilities of the relative sign are taken into account. Figure 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

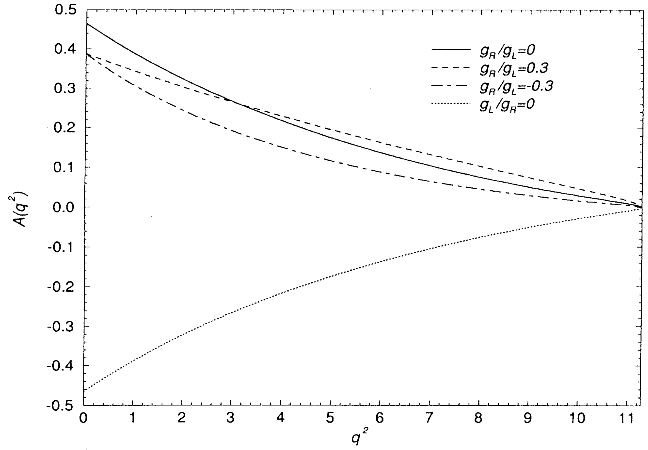

FIG. 1. $A(q^2)$ in the heavy quark limit for the pure left-handed case (solid line), the 30% right-handed contamination case (dashed and dash-dotted lines), and the pure right-handed case (dotted line).

accuracy will give an upper bound for the strength of the right-handed current or will find the 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:

$$A(0) = -\frac{1}{2}\mathcal{P}\frac{|g_L|^2 - |g_R|^2}{|g_L|^2 + |g_R|^2}, \tag{12}$$

$$A([m_{\Lambda_b} - m_{\Lambda_c}]^2) = 0. \tag{13}$$

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 noted that the phase space volume is finite at zero momentum transfer, while it vanishes at the point of zero recoil [cf. Eq. (5)].

IV. $1/m_c$ CORRECTION

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

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

$$\begin{aligned} &F_1 = (1+\Delta)\zeta(q^2)\,,\; G_1 = \zeta(q^2)\,, \\ &F_2 = G_2 = -\Delta\,\zeta(q^2)\,,\; F_3 = G_3 = 0\,, \end{aligned} \tag{14}$$

where

$$\Delta = \frac{\bar{\Lambda}}{m_c}\left(\frac{1}{1+v\cdot v'}\right), \tag{15}$$

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 an analytic formula for $A(q^2)$ in terms of $\bar{\Lambda}$. It is, however, too lengthy to be presented here. We show only the numerical values in Fig. 2.

In Fig. 2, $A(q^2)$ for the pure left-handed case and the pure right-handed case is plotted with and without the $1/m_c$ correction. ($\bar{\Lambda} = 0.7$ GeV and $\bar{\Lambda} = 0$, respectively.) This figure shows that the $1/m_c$ correction decreases quickly as the momentum transfer decreases, and it 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 the $1/m_c$ correction. Using Eqs. (A1), (A2), and (14), one gets

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

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

where $\Delta_0 = \Delta|_{q^2=0}$ and we neglected the terms of higher order in Δ_0. Using Eqs. (7), (16), and (17), one finds the same result as Eq. (12) [22].

From the above argument, we conclude that the $1/m_c$ correction is negligible if one does not go near the zero recoil point ($q^2 = [m_{\Lambda_b} - m_{\Lambda_c}]^2$). In particular, 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 Λ_b to measure the asymmetry $A(q^2)$. If the initial b quark hadronizes into a Λ_b without other hadrons, the absolute value of the Λ_b momentum is simply given by $\sqrt{(m_Z/2)^2 - m_{\Lambda_b}^2}$. In general, however, b quarks in Z decays hadronize with several hadrons; i.e., they form jets. Then the magnitude of the Λ_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 Λ_b [23]. This affects the measurement of the double differential rate in Eq. (3).

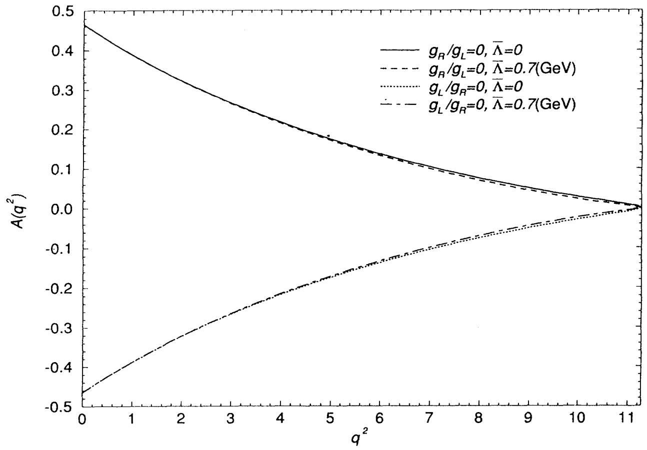

FIG. 2. $A(q^2)$ for the pure left-handed case with (dashed line) and without (solid line) the $1/m_c$ correction, and that for the pure right-handed case with (dash-dotted line) and without (dotted line) the $1/m_c$ correction.

To discuss the fragmentation effect, we consider the model of Peterson *et al.* [24]. In this model, the fragmentation function is given by

$$D(z) = Nz^{-1}\left(1 - \frac{1}{z} - \frac{\epsilon}{1-z}\right)^{-2}, \tag{18}$$

where N is the normalization factor and $\epsilon \simeq m_q^2/m_Q^2$, where m_q is the mass of light quark and m_Q is that of the heavy quark. Equation (18) describes the experimental data of Λ_c production well [25], and the scaling of ϵ as m_Q^{-2} has been observed by comparing charm and bottom fragmentation [26]. 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

$$z_{\max} = 1 + \frac{\epsilon}{2} - \frac{1}{2}\sqrt{\epsilon(\epsilon+4)} \tag{19}$$

$$= 1 - \sqrt{\epsilon} + \cdots . \tag{20}$$

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

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

$$\langle z\rangle = 1 + \frac{2}{\pi}\sqrt{\epsilon}\ln\epsilon + \cdots . \tag{21}$$

This equation suggests that the effect of fragmentation is not so small. To be more precise, a Monte Carlo study seems to be needed to extract the asymmetry in Eq. (6) from experimental data. This is beyond the scope of the present paper.

(ii) *Λ_b from polarized e^+e^- collision near threshold*: To get rid of the fragmentation effect, one may produce Λ_b in e^+e^- collision near the $\Lambda_b\bar\Lambda_b$ threshold. In this case, however, polarized beams are needed to get the necessary polarization of Λ_b. In collisions of right-handed electrons and left-handed positrons, the polarization of the b quark is given by

$$\mathcal{P} = \frac{2s\cos\theta}{s(1+\cos^2\theta) + 4m_b^2(1-\cos^2\theta)}, \tag{22}$$

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

If one approaches close to the threshold, the produced Λ_b is almost at rest. In the rest frame of Λ_b, the asymmetry defined in Eq. (6) has no meaning. However, we can observe the same physical effect, i.e., the spin-momentum correlation, by measuring the angular distribution of Λ_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 Λ_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 evidence of the existence of the right-handed $b \to c$ current with very small theoretical uncertainties.

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APPENDIX: 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):

$$J(q^2) = \frac{|f_V|^2+|f_A|^2}{192\pi^3\beta E_b}\Bigg[\,|g_V|^2\left[w_+w_-\left\{F_1^2 + 2(m_{\Lambda_b}+m_{\Lambda_c})F_1F_+ + w_+F_+^2\right\} + 3q^2w_-F_1^2\right]$$
$$+ |g_A|^2\left[w_+w_-\left\{G_1^2 - 2(m_{\Lambda_b}-m_{\Lambda_c})G_1G_+ + w_-G_+^2\right\} + 3q^2w_+G_1^2\right]\Bigg], \tag{A1}$$

$$K(q^2) = \frac{|f_V|^2 + |f_A|^2}{96\pi^3\beta^2 E_b}(g_V g_A^* + g_V^* g_A) m_{\Lambda_b}^2 \Big[(m_{\Lambda_b}^2 - m_{\Lambda_c}^2 - 2q^2) F_1 G_1 - (m_{\Lambda_b} + m_{\Lambda_c}) w_- F_1 G_+ + (m_{\Lambda_b} - m_{\Lambda_c}) w_+ F_+ G_1 - w_+ w_- F_+ G_+ \Big] , \qquad \text{(A2)}$$

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 and G_1 and the combinations F_+ and G_+ survive in the above expressions because of vanishing lepton masses.

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