Test of the chiral structure of the top-bottom charged current by the process $b \rightarrow s\gamma$

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The chiral structure of the top-bottom charged current is examined by using the process $b \rightarrow s\gamma$. The new CLEO results on the radiative B decay constrain the possible deviation from the $V-A$ coupling to be less than a few percent in the sector of the W boson and top and bottom quarks.

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One of the remarkable predictions of the $SU(2)_L \times U(1)_Y$ electroweak interactions is the universal $V-A$ structure of the interactions between the $SU(2)_L$ gauge bosons and quarks and leptons. The pure $V-A$ structure of the $SU(2)_L$ gauge interactions implies that the left- and right-handed components of the quarks are particles with different quantum numbers and it might suggest that each chiral component of the quarks is a pointlike rather than some composite object. The possible violation of the $V-A$ structure would provide useful information on the new properties of the quarks at higher energy scales and also on the origin of parity violation. In particular, the interactions of the top quark are of special importance, since the top quark is far more massive than the other fermions, and its interactions may be quite sensitive to new physics originating at higher energy scales [1]. In this paper, the chiral structure of the top-bottom charged current is examined by using the measured process $b \rightarrow s\gamma$. A right-handed coupling of the top and bottom quarks to the W boson, when treated on the basis of a reasonable ansatz, gives rise to a large contribution to the decay $b \rightarrow s\gamma$. This enhanced radiative decay arises from a chirality flip induced by the large top quark mass in the intermediate states [2]. The violation of the $V-A$ structure is thus constrained to be less than a few percent by the current measurement of $b \rightarrow s\gamma$ by the CLEO Collaboration [3,4].

The coupling of the W boson to the quarks in the standard model is given by

$$\mathcal{L}_{SM}=\frac{g_2}{2\sqrt{2}}(\bar{u},\bar{c},\bar{t})\gamma^{\mu}(1-\gamma_5)W_{\mu}^{+}V\begin{pmatrix} d\\ s\\ b\end{pmatrix}+\text{H.c.}\ , \qquad (1)$$

where V is the Cabibbo-Kobayashi-Maskawa matrix. Here we introduce a phenomenological coupling of the right-handed top and bottom quarks to the W boson,

$$\mathcal{L}_{RH}=\frac{g_2}{2\sqrt{2}}V_{tb}f_R^{tb}\bar{t}\gamma^{\mu}(1+\gamma_5)bW_{\mu}^{+}+\text{H.c.}\ , \qquad (2)$$

in addition to the interaction (1). Using the interaction Lagrangians (1) and (2), the effective Hamiltonian for weak radiative B-meson decay is calculated as

$$H_{\text{eff}}=\frac{e}{16\pi^2}\frac{2}{\sqrt{2}}G_F V_{tb}V_{ts}^{*}m_b\bar{s}_L\sigma^{\mu\nu}b_R F_{\mu\nu}\times\left\{A_{SM}(x)+f_R^{tb}\frac{m_t}{m_b}A_{RH}(x)\right\}\ , \qquad (3)$$

where the masses of the light quarks u, d, s, and c are neglected. In Eq. (3), m_b is the mass of the bottom quark, $F_{\mu\nu}$ is the electromagnetic field strength tensor, and A_{SM} and A_{RH} are the functions of the top quark mass squared divided by the W-boson mass squared. The contribution A_{SM} from the standard model is given by [5]

$$A_{SM}(x)=\frac{1}{(1-x)^4}Q_t\left\{\frac{x^4}{4}-\frac{3}{2}x^3+\frac{3}{4}x^2+\frac{x}{2}+\frac{3}{2}x^2\ln(x)\right\}+\frac{1}{(1-x)^4}\left\{\frac{x^4}{2}+\frac{3}{4}x^3-\frac{3}{2}x^2+\frac{x}{4}-\frac{3}{2}x^3\ln(x)\right\}, \qquad (4)$$

with $x=m_t^2/m_W^2$, and the right-handed interaction (2) yields the contribution

$$A_{RH}(x)=\frac{1}{(1-x)^3}Q_t\left\{-\frac{x^3}{2}-\frac{3}{2}x+2+3x\ln(x)\right\}+\frac{1}{(1-x)^3}\left\{-\frac{x^3}{2}+6x^2-\frac{15}{2}x+2-3x^2\ln(x)\right\}, \qquad (5)$$

where Q_t is 2/3. (The effects of the strong interaction are taken into account later.)

In computing the Hamiltonian (3), we first write the Feynman amplitudes by using the unitary gauge propagator for the W boson to avoid the contributions of negative metric components. We then put all the internal lines on the mass shell in the lowest order triangle diagrams, or equivalently we write unsubtracted dispersion relations for all the internal lines. For example, we begin with the expression

$$\bar{s}\gamma_{\alpha}(c_V-c_A\gamma_5)(\not{l}+m_t)\gamma_{\beta}(c_V'-c_A'\gamma_5)b\left[g^{\alpha\nu}-\frac{k'^{\alpha}k'^{\nu}}{m_W^2}\right]\times e\epsilon^{\mu}\Gamma_{\mu\nu\lambda}\left[g^{\lambda\beta}-\frac{k^{\lambda}k^{\beta}}{m_W^2}\right], \qquad (6)$$

$$\Gamma_{\mu\nu\lambda}=\{-(k+k')_{\mu}g_{\nu\lambda}+(k'+q)_{\lambda}g_{\mu\nu}+(k-q)_{\nu}g_{\mu\lambda}\}\ ,$$

for the amplitude, where only the numerator is written. In Eqs. (6), l is the four-momenta of the top quark, q_μ and ϵ^μ are the four-momenta and the polarization vector of the outgoing photon ($\epsilon^\mu q_\mu=0$), respectively, and k and $k'=k-q$ are the four-momenta of the two W-boson lines. The numerator of the top quark propagator $(\not l+m_t)$ in Eq. (6) is rewritten as a spin sum of the top quark wave functions as

$$\bar{s}\gamma_\alpha(c_V-c_A\gamma_5)\left\{\sum_{\text{spin}} t\bar{t}\right\}\gamma_\beta(c_V'-c_A'\gamma_5)b \ . \tag{7}$$

Using the on-shell conditions for three momenta q, k, and $k'=k-q$ and $\epsilon^\mu q_\mu=0$, the product of the electromagnetic vertex factor $\epsilon^\mu\Gamma_{\mu\nu\lambda}$ and the numerator factors of the gauge boson propagators can be decomposed as

$$eg^{\alpha\nu}\epsilon^\mu\Gamma_{\mu\nu\lambda}g^{\lambda\beta}-em_W g^\alpha_\mu\epsilon^\mu\frac{k^\beta}{m_W}-\frac{k'^\alpha}{m_W}\epsilon^\mu e m_W g^\beta_\mu + \frac{k'^\alpha}{m_W}\epsilon^\mu e(k+k')_\mu\frac{k^\beta}{m_W} \ . \tag{8}$$

The momentum factors k'^α/m_W and k^β/m_W in Eq. (8) yield the couplings of the unphysical scalar to the quarks, when combined with the expression (7):

$$\bar{s}\gamma_\alpha(c_V-c_A\gamma_5)t\frac{k'^\alpha}{m_W} =\frac{1}{m_W}\bar{s}\{m_s(c_V-c_A\gamma_5)-m_t(c_V+c_A\gamma_5)\}t \ ,$$
$$\frac{k^\beta}{m_W}\bar{t}\gamma_\beta(c_V'-c_A'\gamma_5)b =\frac{1}{m_W}\bar{t}\{-m_t(c_V'-c_A'\gamma_5)+m_b(c_V'+c_A'\gamma_5)\}b \ , \tag{9}$$

where the equations of motion are used for strange, bottom and top quarks. By this way, the second and third terms in Eq. (8) give the couplings of the unphysical scalar to the photon and W boson in the 't Hooft–Feynman gauge, while the fourth term gives the interaction of the unphysical scalar to the photon. We thus recover the Feynman rules written in 't Hooft–Feynman gauge for the standard model couplings. Corresponding to the right-handed coupling (2), we find an extra effective coupling of the unphysical scalar described by

$$\mathcal{L}_\phi=\frac{g_2}{2\sqrt{2}m_W}V_{tb}f_R^{tb}\bar{t}\{m_t(1+\gamma_5) -m_b(1-\gamma_5)\}b\phi^+ +\text{H.c.} \tag{10}$$

When all the quarks are supposed to be elementary, the above prescription is justified if one constructs a vector-like model, as in Ref. [1], which incorporates the interaction (2). If one supposes that the right-handed coupling (2) arises from the possible composite nature of top and bottom quarks, for example, the coupling (10) may be understood as an effective coupling induced by the gauge invariance and unitarity in the more fundamental level.

This calculational prescription gives a well-defined finite contribution of the interaction (2) to the Hamiltonian (3), without anomalously enhancing the contributions of the longitudinal component of the W boson. The contribution of the right-handed interaction to the Hamiltonian (3) is enhanced by the factor m_t/m_b because of the chirality flip proportional to the large top quark mass in the intermediate states [2]: To be more precise, the chirality flip amplitude is relatively *suppressed* by the pure $V-A$ currents in the standard model.

Using the Hamiltonian (3), the branching fraction $B(b\to s\gamma)$ is computed following the procedure of Ref. [6] by normalizing the decay width $\Gamma(b\to s\gamma)$ to the semileptonic decay width $\Gamma(b\to ce\bar{\nu})$,

$$B(b\to s\gamma)=\frac{\Gamma(b\to s\gamma)}{\Gamma(b\to ce\bar{\nu})}B(b\to ce\bar{\nu}) \ , \tag{11}$$

and using $B(b\to ce\bar{\nu})\simeq 11\%$ [7]. The leading QCD corrections to the Hamiltonian (3) are significant [6,8]. These QCD corrections in the standard model were computed in Ref. [9] by analyzing the operator mixing between the magnetic moment operator in Eq. (3) and the four Fermi operators involving the quarks lighter than the W bosons, e.g., $\bar{s}\gamma^\mu c\bar{c}\gamma_\mu b$. Since the running of the coupling constant is small from the scale m_t to m_W, and the same set of composite operators appear below the scale of m_W in our analysis, we incorporate these QCD corrections following the prescription of Ref. [9]. Including the QCD corrections, the rate for $\Gamma(b\to s\gamma)$ normalized to the semileptonic rate is given by

$$\frac{\Gamma(b\to s\gamma)}{\Gamma(b\to ce\bar{\nu})}=\frac{3\alpha}{2\pi\rho(m_c/m_b)(1+\delta_{\text{QCD}})}\eta^{-32/33}\left\{\frac{116}{135}(\eta^{10/23}-1)+\frac{116}{378}(\eta^{28/23}-1)+A_{\text{SM}}+f_R^{tb}\frac{m_t}{m_b}A_{\text{RH}}\right\}^2 , \tag{12}$$

where $\eta=\alpha_s(m_b^2)/\alpha_s(m_W^2)\simeq 1.8$. In Eq. (12), $\rho(m_c/m_b)$ and δ_{QCD} are the phase suppression factor and the QCD corrections to the semileptonic decay, respectively. These factors are evaluated as $\rho(m_c/m_b)=0.447$ and $\delta_{\text{QCD}}=(2\alpha_s(m_b^2)3\pi)f(m_c/m_b)$, where $f(m_c/m_b)=2.41$ [10]. [We take $\alpha_s(m_b^2)=0.23$.] In Fig. 1, the branching fraction $B(b\to s\gamma)$ is given as a function of f_R^{tb} for the three values of the top quark mass, 110, 140, and 170 GeV. The upper and lower bounds on $B(b\to s\gamma)$, $0.6\times10^{-4}\le B(b\to s\gamma)\le 5.4\times10^{-4}$, obtained by the CLEO Collaboration, constrain the parameter f_R^{tb} to be $-0.05\lesssim f_R^{tb}\lesssim 0.01$, with the vicinity of $f_R^{tb}=-0.02$ being excluded.

For comparison, we carry out a preliminary investigation of the effects of the right-handed couplings of the top and bottom quarks to the $SU(2)_L$ gauge bosons on the ρ

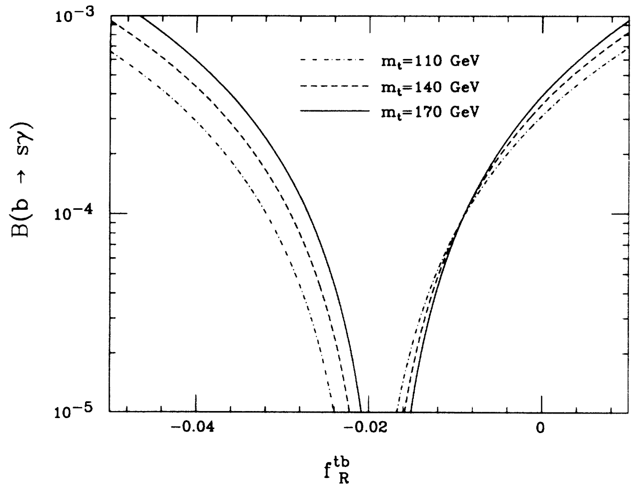

FIG. 1. The branching fraction $B(b \rightarrow s\gamma)$ as a function of f_R^{tb} for the three values of the mass of the top quark, 110, 140, and 170 GeV.

parameter [11]. For definiteness, the phenomenological interaction of the form

$$\mathcal{L} = \frac{g_2}{2\sqrt{2}} f_R^{tb} \bar{t} \gamma^\mu (1+\gamma_5) b W_\mu^+ + \text{H.c.} + \frac{g_2}{4} \{ f_R^{tt} \bar{t} \gamma^\mu (1+\gamma_5) t - f_R^{bb} \bar{b} \gamma^\mu (1+\gamma_5) b \} W_\mu^3 , \quad (13)$$

is assumed, where W_μ^3 is the neutral $SU(2)_L$ gauge boson. The interaction (13) yields the following corrections to the ρ parameter:

$$\Delta\rho = \frac{G_F N_C m_t^2}{2\sqrt{2}\pi^2} \left[f_R^{tt} + \frac{m_b^2}{m_t^2} f_R^{bb} - 2 f_R^{tb} \frac{m_b}{m_t} \right] \ln \left[\frac{\Lambda^2}{m_t^2} \right] , \quad (14)$$

where the dimensional regularization is used in the calculation, and the pole $1/\epsilon$ is replaced by the logarithm of a cutoff squared Λ^2. In Eq. (14), we retain the leading corrections involving $\ln(\Lambda^2)$ and terms linear in f_R. The experimental limit on the ρ parameter is $\Delta\rho = (-0.1 \pm 2.7) \times 10^{-3}$ [12]. To compare (14) with the experimental bound, we tentatively choose $\Lambda = 10$ TeV for the threshold energy of some new physics, and $m_t = 140$ GeV. The parameter f_R^{bb} affects the $Z \rightarrow b\bar{b}$ cross section at the tree level, and thus is severely constrained [13]. So we set $f_R^{bb} = 0$ in Eq. (14). The ρ parameter leads to a rather tight bound $|f_R^{tt}| \lesssim 0.01 \sim 0.02$ for f_R^{tt} [14]. However, the bound on f_R^{tb} from the ρ parameter is not so stringent; $|f_R^{tb}| \lesssim 0.2$. (The other electroweak parameters yield the bounds on f_R weaker than those obtained from the above analysis.) Here we note that if the right-handed couplings are induced in some perturbatively controllable renormalizable theories, the bounds on f_R^{tb} and f_R^{tt} obtained in the analysis of the ρ parameter would be weaker than those obtained above, because the large logarithm will be replaced by some moderate function due to the decoupling phenomena. On the other hand, our analysis on $b \rightarrow s\gamma$ is expected to remain valid in those models.

Finally we briefly review the experimental bounds on the right-handed charged currents involving the other quarks. For the up-down and up-strange charged currents, the bound $g_R/g_L \lesssim 0.1$ is obtained [15]. (Here, $g_{L,R}$, respectively denote the coupling constants of left- and right-handed quarks to the W boson.) For the charged currents involving heavy (c, b, and t) quarks, the bound $g_R/g_L \lesssim 0.3$ is obtained for the down-charm and strange-charm charged currents [16], and $g_R/g_L \lesssim 0.5$ is obtained for the bottom-charm charged current [17].

In summary, the possible right-handed top-bottom charged current is investigated by using the process $b \rightarrow s\gamma$. The new CLEO results combined with our analysis imply that the top-bottom charged current is left-handed within the accuracy of a few percent. It is interesting that the chiral structure of the top-bottom charged current is, although indirectly, constrained most stringently among the charged currents involving the heavy (c, b, and t) quarks.

Note added. After submitting our manuscript, the paper by D. Cocolicchio, G. Costa, G. L. Fogli, J. H. Kim, and A. Masiero, Phys. Rev. D **40**, 1477 (1989) came to our attention; they studied the effect of the right-handed charged currents on the $b \rightarrow s\gamma$ in the left-right symmetric models. Though the physics contents discussed in their paper are different, our function A_{RH} in Eq. (5) can be compared to their result; their analytic expression [$f_{2\gamma}^{LR}$ in Eq. (19)] disagrees with ours and diverges at $m_t = m_W$. After the submission of our paper, there appeared a paper by P. Cho and M. Misiak [Phys. Rev. D (to be published)] who investigate the same subject as D. Cocolicchio *et al.* Our expression in Eq. (5) agrees with that of Cho and Misiak. We are grateful to K. S. Babu for informing us about the paper by Cho and Misiak.

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