Enhanced Electroweak Penguin Amplitude in $B \to VV$ Decays

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(Received 21 December 2005; published 10 April 2006)

We discuss a novel electromagnetic penguin contribution to the transverse helicity amplitudes in B decays to two vector mesons, which is enhanced by two powers of m_B/Λ relative to the standard penguin amplitudes. This leads to unique polarization signatures in penguin-dominated decay modes such as $B \to \rho K^*$ similar to polarization effects in the radiative decay $B \to K^*\gamma$ and offers new opportunities to probe the magnitude and chirality of flavor-changing neutral current couplings to photons.

DOI: 10.1103/PhysRevLett.96.141801 PACS numbers: 12.60.−i, 13.20.He

Introduction.—Decays of B mesons into two charmless mesons provide an abundant source of information on flavor- and CP-violating phenomena in the weak interactions of quarks. In particular, decays to two vector mesons ($B \to VV$) can shed light on the helicity structure of these interactions through polarization studies. While predicted to be fundamentally $V - A$ in the standard model (SM), a deviation from this expectation cannot currently be excluded. The first observations of $B \to VV$ decays show no anomalies in the helicity structure but point to a reduced amount of longitudinal polarization in penguin-dominated decays [1]. This has led to theoretical studies that reconsider strong interactions effects in $B \to VV$ decays [2–4] or invoke new fundamental interactions [5].

Any particular $B \to VV$ decay is characterized by the three helicity amplitudes A_0 (longitudinal), A_-, and A_+. A quark model or naive factorization analysis [6] leads to the expectation that, for $\bar{B}$, i.e., b-quark, decay, the helicity amplitudes are in proportions

$$A_0{:}A_-{:}A_+ = 1{:}\frac{\Lambda}{m_b}{:}\left(\frac{\Lambda}{m_b}\right)^2, \tag{1}$$

with $\Lambda \approx 0.5$ GeV the strong-interaction scale and $m_b \approx 5$ GeV the bottom quark mass. This expectation has been parametrically (not necessarily numerically) confirmed [2] in the framework of QCD factorization, which provides a theoretical basis for the heavy-quark expansion of B decays to charmless mesons [7]. The hierarchy (1) of helicity amplitudes follows from the $V - A$ structure of the standard weak interactions.

In this Letter, we point out and discuss an effect which has been neglected in all previous studies of $B \to VV$ but which substantially alters the prediction for polarization observables. The effect is connected with electromagnetic penguin transitions and appears only for neutral vector mesons. It leads to the unique feature that the transverse electroweak penguin amplitude is dominated by the electromagnetic dipole operator providing a signature similar to polarization in radiative decays $B \to K^*\gamma$ [8], but which is easier to access experimentally.

The effect in question is related to the two diagrams shown in Fig. 1. When the vector meson V_2 is transversely polarized, there exists a large contribution to the decay amplitude due to the small virtuality $m_{V_2}^2$ of the intermediate photon propagator. This is in contrast to the case of longitudinal polarization, where the photon propagator is canceled, and the amplitude is local on the scale m_b [9]. The large transverse amplitude is best described by a short-distance transition $b \to D\gamma$ ($D = d, s$), followed by the transition of the low-virtuality photon ($q^2 \ll m_b^2$) to the neutral vector meson. We shall perform a factorization analysis of the amplitude below.

The calculation of the diagrams in Fig. 1 is straightforward. The weak interactions are given in terms of the standard effective Hamiltonian [10]. We use the conventions of Ref. [11] but generalize the electromagnetic dipole operators to include both chiralities

$$\mathcal{H}_{\rm eff} = \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^{(D)} \sum_{a=-,+} C_{7\gamma}^a Q_{7\gamma}^a + \ldots, \tag{2}$$

$$Q_{7\gamma}^{\mp} = -\frac{e\bar{m}_b}{8\pi^2} \bar{D}\sigma_{\mu\nu}(1 \pm \gamma_5) F^{\mu\nu} b, \tag{3}$$

where $\lambda_p^{(D)} = V_{pb}V_{pD}^*$. The ellipses denote other operators (see [11]). In the SM, $C_{7\gamma}^+$ is suppressed by a factor m_D/m_b; hence, $Q_{7\gamma}^+$ is usually neglected. The remaining term is then simply denoted by $C_{7\gamma}Q_{7\gamma}$. However, in

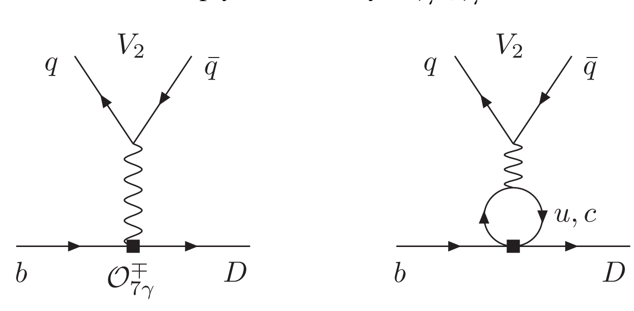

FIG. 1. Leading contributions to $\Delta\alpha_{3,\rm EW}^{p\mp}(V_1V_2)$ defined in the text.

generic extensions of the SM, there is no reason to expect a suppression of additional contributions to $C_{7\gamma}^{+}$ relative to $C_{7\gamma}^{-}$. The coupling of the photon to the quark electric charge in V_2 implies that the diagrams in Fig. 1 contribute to the electroweak penguin amplitude in the general flavor decomposition of hadronic two-body decay amplitudes. Adopting the α_i notation of Ref. [9] extended to allow for the three helicity amplitudes of $B \to VV$, the new contribution to the transverse electroweak penguin amplitudes is

$$\Delta\alpha_{3,\mathrm{EW}}^{p\mp}(V_1V_2) = \mp\frac{2\alpha_{\mathrm{em}}}{3\pi}C_{7\gamma,\mathrm{eff}}^{\mp}R_{\mp}\frac{m_B\bar{m}_b}{m_{V_2}^2}, \qquad (4)$$

with $C_{7\gamma,\mathrm{eff}}^{\mp}$ taking into account the effect of quark loop diagrams (see Fig. 1). $R_{\mp}$ is a ratio of tensor to (axial) vector $B \to V_1$ form factors such that R_- equals 1 in the heavy-quark limit [12], while R_+ is of order m_b/Λ. We note the large enhancement factor $m_B\bar{m}_b/m_{V_2}^2 \sim (m_b/\Lambda)^2$, which implies that the first hierarchy in (1) is inverted, rendering the negative-helicity amplitude A_- leading over the longitudinal amplitude A_0 in the heavy-quark limit. Of course, for real values of m_b/m_{V_2}, this enhancement is compensated by the small electromagnetic coupling $\alpha_{\mathrm{em}} = e^2/(4\pi)$. For instance, for neutral ρ mesons, we obtain $\Delta\alpha_{3,\mathrm{EW}}^{p-}(K^*\rho) \approx 0.02$. This should be compared to the uncorrected negative-helicity electroweak penguin amplitude

$$\alpha_{3,\mathrm{EW}}^{p-}(K^*\rho) = C_7 + C_9 + \frac{C_8 + C_{10}}{N_c} + \ldots \approx -0.01 \qquad (5)$$

and the leading QCD penguin amplitude

$$\hat{\alpha}_4^{c-}(\rho K^*) = C_4 + \frac{C_3}{N_c} + \ldots \approx -0.055. \qquad (6)$$

The C_i are Wilson coefficients for the various penguin operators in the effective Hamiltonian [10], and the ellipses denote the 1-loop corrections in QCD factorization [4], which we have taken into account in the numerical estimates. In the SM, the corresponding positive-helicity amplitudes are suppressed by about an order of magnitude relative to the negative-helicity ones as explained above.

There are strong-interaction corrections to the leading-order expression (4) from gluon exchange between the quark lines in the second diagram in Fig. 1 and also through hard interactions with the spectator quark (not shown in the figure) in the B meson. Because of factorization as discussed below, these corrections modify only the effective $b \to D\gamma$ transition at leading order in the expansion in Λ/m_b. They have been computed in next-to-leading order in the context of factorization of exclusive radiative B decays [13] and can be incorporated by substituting $C_{7\gamma}^{-} \to \mathcal{C}_7^{\prime}$ [first paper of Ref. [13], Eq. (62)]. Turning this argument around, the absolute value of $\Delta\alpha_{3,\mathrm{EW}}^{c-}(K^*V_2)$ can be obtained from the branching fraction of $B \to K^*\gamma$ via

$$|\Delta\alpha_{3,\mathrm{EW}}^{c-}(K^*V_2)| = \frac{2\alpha_{\mathrm{em}}}{3\pi}R_-\frac{m_B^2}{m_{V_2}^2} \times\left(\frac{\Gamma(B\to K^*\gamma)}{\frac{G_F^2|V_{ts}^*V_{tb}|^2}{8\pi^3}\frac{\alpha_{\mathrm{em}}}{4\pi}m_B^5T_1^{K^*}(0)^2}\right)^{1/2}, \qquad (7)$$

with $T_1^{K^*}(0) \approx 0.28$ a tensor form factor. This results in $|\Delta\alpha_{3,\mathrm{EW}}^{c-}(K^*\rho)| = 0.023$, close to the leading-order estimate from (4).

We therefore conclude that the new radiative contribution to the negative-helicity electroweak penguin amplitude is at least twice as large (and opposite in sign) as was previously assumed. For penguin-dominated $b \to s$ transitions, it is almost half the size of the leading QCD penguin amplitude and should, therefore, have visible impact on polarization measurements. In the case of new interactions generating $C_{7\gamma}^{+}$, the corresponding contribution to the positive-helicity amplitude (4) should be observed against a very small standard model background.

Factorization analysis.—Since the existence of an amplitude violating the power counting (1) may appear surprising, we sketch how this amplitude emerges and factorizes in soft-collinear effective theory (SCET) [14]. The notation and method of the following discussion is similar to the one in Ref. [15]. After integrating out the scale m_b, SCET formalizes the interaction of the static b-quark field h_v with collinear fields for the lightlike direction n_-, in which meson V_1 moves, and collinear fields for the lightlike direction n_+ of meson V_2. Let χ denote the collinear quark field corresponding to V_2, and let V_2 be the meson that does not pick up the spectator quark from the B meson. The leading quark bilinears that have nonvanishing overlap with $\langle V_2|$ are

$$\bar{\chi}\not{n}_-(1\mp\gamma_5)\chi, \qquad \bar{\chi}\not{n}_-\gamma_\perp^\mu(1\pm\gamma_5)\chi. \qquad (8)$$

The subscript $\perp$ denotes projection of a Lorentz vector on the plane transverse to the two light-cone vectors $n_\mp$. Both operators scale as λ^4 according to the SCET scaling rules; the first overlaps only with the longitudinal polarization state of V_2, the second only with a transverse vector meson. However, the second operator is not generated by the $V-A$ interactions of the SM (at least at the tree and 1-loop level). This implies the power suppression of $A_\mp$ relative to A_0 in (1), since the leading contribution to transverse polarization now involves an operator with an additional derivative $D_\perp \sim \lambda^2 \sim \Lambda/m_b$.

This reasoning ignores electromagnetic effects. Including QED in SCET, there is a collinear photon field with unsuppressed interactions with collinear quarks (of the same direction). Only the transverse photon field is truly a degree of freedom of the theory, since the other two components are either gauge artifacts or can be eliminated by the field equations. Hence, there is an additional operator $eA_{\gamma\perp}^\mu = W_\gamma^\dagger iD_{\gamma\perp}^\mu W_\gamma$ (where W_γ is an electromagnetic Wilson line formally required to make the operator gaugeinvariant), which overlaps only with a transversely polarized vector meson. To first order in the electromagnetic coupling, the matrix element can be computed exactly, yielding

$$\langle V_2|[W_\gamma^\dagger iD_{\gamma\perp}^\mu W_\gamma](0)|0\rangle = -\frac{2i}{3}a_{V_2}\frac{e^2 f_{V_2}}{m_{V_2}}\epsilon_\perp^{*\mu}, \quad (9)$$

with $\epsilon_\perp^\mu$ a transverse polarization vector, f_{V_2} the decay constant, and a_{V_2} a constant that depends on the quark-flavor composition of V_2, $a_\rho = 3/2$, $a_\omega = 1/2$, $a_\phi = -1/2$. [The convention for the covariant derivative corresponding to (3) is $iD_\gamma^\mu = i\partial^\mu + e_q A_\gamma^\mu$, with e_q the quark electric charge.] The crucial point is that the operator $W_\gamma^\dagger iD_{\gamma\perp}^\mu W_\gamma$ scales with λ^2; hence, this contribution to $A_\mp$ is a factor m_b/Λ larger than even the longitudinal amplitude A_0. Thus, we find the tree-level matching equation (see also [16])

$$Q_{7\gamma}^\mp \to -\frac{\bar{m}_b m_B}{4\pi^2}[\bar{\xi}W\gamma_{\perp\mu}(1\mp\gamma_5)h_v](0)[W_\gamma^\dagger iD_{\gamma\perp}^\mu W_\gamma](0), \quad (10)$$

valid as an equation for the $\langle V_1 V_2|\ldots|\bar{B}\rangle$ matrix element. In SCET, only soft fields can couple to the two brackets representing collinear field products in the two different directions. But since the photon operator in the second bracket is a color singlet, the soft fields decouple, and the matrix element of the right-hand side of (10) falls apart into (9) and $\langle V_1|\bar{\xi}W\gamma_\perp^\mu(1\mp\gamma_5)h_v|\bar{B}\rangle$, which is proportional to the SCET form factor $\xi_\perp$ [12] at maximal recoil. Equation (10) has to be amended by radiative corrections as well as a second operator structure with an additional transverse derivative in the first bracket. This is very similar to heavy-to-light form factors [15]; in fact, these corrections simply restore the QCD tensor form factor. Combining (9) and (10), we therefore find

$$\langle V_1 V_2|C_{7\gamma}^\mp Q_{7\gamma}^\mp|\bar{B}\rangle = im_{V_2}m_B 2T_1^{V_1}(0)f_{V_2}a_{V_2} \times\left(\mp\frac{2\alpha_{\rm em}}{3\pi}\right)C_{7\gamma}^\mp\frac{m_B\bar{m}_b}{m_{V_2}^2}, \quad (11)$$

which on accounting for the normalization of $\alpha_{3,\rm EW}^{p,h}$ [4,9] reproduces (4). The previous equation should be understood such that the matrix element of $Q_{7\gamma}^-$ ($Q_{7\gamma}^+$) takes the value given only when both V_1 and V_2 have negative (positive) helicity but is zero otherwise. In general, the four-quark operators from the effective weak Hamiltonian also contribute to the matching coefficient of the SCET operator on the right-hand side of (10), and including further spectator-scattering effects replaces $C_{7\gamma}^-$ by $\mathcal{C}_7'$ as discussed above.

The $B\to\rho K^*$ system.—We now focus on the eight $B\to\rho K^*$ decay modes, where the electroweak penguin amplitude is largest relative to the leading QCD penguin amplitude ($a_\rho = 3/2$). Assuming isospin symmetry, the ρK^* system is described by six complex strong-interaction parameters for each helicity $h = 0, -, +$. Neglecting the color-suppressed electroweak penguin amplitude and the doubly Cabibbo-Kobayashi-Maskawa (CKM) suppressed QCD penguin amplitude is a good approximation for elucidating the effect of the new (color-allowed) electroweak penguin contribution; hence, we write

$$\begin{aligned} A_h(\rho^-\bar{K}^{*0}) &= P_h, \\ \sqrt{2}A_h(\rho^0 K^{*-}) &= [P_h + P_h^{\rm EW}] + e^{-i\gamma}[T_h + C_h], \\ A_h(\rho^+ K^{*-}) &= P_h + e^{-i\gamma}T_h, \\ -\sqrt{2}A_h(\rho^0\bar{K}^{*0}) &= [P_h - P_h^{\rm EW}] + e^{-i\gamma}[-C_h] \end{aligned} \quad (12)$$

and define $x_h = X_h/P_h$, where P_h is the QCD penguin amplitude. The tree amplitudes T_h, C_h are suppressed by the CKM factor $\epsilon_{\rm KM} = |V_{ub}V_{us}^*|/|V_{cb}V_{cs}^*| \sim 0.025$. Assuming $\gamma = 70°$ is known, one can obtain P_h from an angular analysis of the $\rho^-\bar{K}^{*0}$ final state, t_h from $\rho^\pm K^{*\mp}$, and $p_h^{\rm EW}$ and c_h from the remaining four decay modes. In principle, this allows for a determination of $P_h^{\rm EW}$, which can be compared to the theoretical result. In practice, a complete amplitude analysis will be experimentally difficult.

The sensitivity to the electroweak penguin amplitude is made apparent in CP-averaged helicity-decay rate ratios such as

$$S_h \equiv \frac{2\bar{\Gamma}_h(\rho^0\bar{K}^{*0})}{\bar{\Gamma}_h(\rho^-\bar{K}^{*0})} = |1 - p_h^{\rm EW}|^2 + \Delta_h, \quad (13)$$

where Δ_h depends on c_h (and mildly on $p_h^{\rm EW}$) and vanishes for $c_h \to 0$. To estimate S_-, we assume that the positive-helicity amplitudes are negligible as predicted in the SM and use the observed $\rho^-\bar{K}^{*0}$ branching fraction and longitudinal polarization fraction f_L to determine the magnitude of P_0 and P_-. We shall also assume that the phase of $p_h^{\rm EW}$ is not more than 30° away from 0 or π. Writing $p_h^{\rm EW} = [P_h^{\rm EW}/T_h]\times t_h$, this amounts to the assumption that no large CP asymmetries will be found in $B\to\rho^\pm K^{*\mp}$. For all other quantities, we perform a calculation in the QCD factorization framework. In this procedure, there is a considerable uncertainty in P_- due to the discrepant experimental results on $f_L(\rho^+K^{*0})$ [1], which may result in an overestimate of P_- and, hence, an underestimate of $p_-^{\rm EW}$. It is therefore not excluded that the electromagnetic penguin effect is more pronounced than in the following theoretical estimates. Keeping this in mind, we find $\mathrm{Re}(p_-^{\rm EW}) = -0.23 \pm 0.08[+0.14^{+0.04}_{-0.05}]$ and $\Delta_- = -0.0 \pm 0.2$, yielding

$$S_- = 1.5 \pm 0.2\,[0.7 \pm 0.1]. \quad (14)$$

Here (and below) the numbers in brackets refer to the calculation without the new electromagnetic penguin contribution. Despite the current large theoretical uncertainties, which could be removed with more experimental data, Eq. (14) clearly shows the impact of this contribution on polarization observables. The effect is even more significant for the ratio of the two final states with neutral ρ

mesons, as S_-/S'_- [(15) below] changes by a factor of about 4 whether or not the electromagnetic penguin contribution is included, but for this ratio the tree contamination is also larger. Data are not currently available to test (14), but we may instead consider

$$S'_h \equiv \frac{2\bar{\Gamma}_h(\rho^0\bar{K}^{*-})}{\bar{\Gamma}_h(\rho^-\bar{K}^{*0})} = |1 + p_h^{\rm EW}|^2 + \Delta'_h. \qquad (15)$$

Following the same strategy as above, we obtain $\Delta'_- = -0.1 \pm 0.0$, and $S'_- = 0.5 \pm 0.1$ [1.2 ± 0.1]. In the absence of direct CP asymmetries, S'_h is directly related to the corresponding ratio of polarization fractions $f'_h \equiv f_h(\rho^0\bar{K}^{*-})/f_h(\rho^-\bar{K}^{*0})$. Including a theoretical estimate of the CP asymmetries, we obtain

$$f'_0 = 1.3 \pm 0.1\,[1.1 \pm 0.1], \qquad (16)$$

$$f'_- = \frac{1 - f_L(\rho^0\bar{K}^{*-})}{1 - f_L(\rho^-\bar{K}^{*0})} = 0.4 \pm 0.1\,[0.8 \pm 0.1]. \qquad (17)$$

This can be compared to the experimental values $f'_0|_{\rm exp} = 1.45^{+0.64}_{-0.58}$, $f'_-|_{\rm exp} = 0.12^{+0.44}_{-0.11}$ [1].

The electromagnetic penguin contribution also applies to the $B \to \phi K^*$ modes, though the effect is smaller by $a_\phi/a_\rho(m_\rho/m_\phi)^2 = 0.19$. This helps to understand the small observed $f_L(\phi\bar{K}^*) \approx 0.5$ by lowering the theoretical result by about 0.05 but cannot account for the difference in f_L between the $\phi\bar{K}^*$ and $\rho^-\bar{K}^{*0}$ final states, which therefore must be due to different transverse QCD penguin amplitudes.

Finally, we comment on the possibility of detecting the presence of new flavor-changing neutral currents in the form of an electromagnetic penguin operator with opposite chirality $Q^+_{7\gamma}$. For this analysis, one must isolate experimentally the positive-helicity amplitudes. Theoretically, all positive-helicity amplitudes are suppressed, except for the electromagnetic penguin contribution $\Delta P_+^{\rm EW}$ to the electroweak penguin amplitude. In the naive factorization approximation $X_+ = rX_-$, where r is a Λ/m_b-suppressed form factor ratio, while $\Delta P_+^{\rm EW} \approx C^+_{7\gamma}/C^-_{7\gamma}\Delta P_-^{\rm EW}$ is suppressed only by the ratio of Wilson coefficients [see (11)]. A conservative analysis of the $b \to s\gamma$ branching fraction constrains $C^+_{7\gamma}/C^-_{7\gamma} < 0.5$; hence, it is possible that the suppression is weak. This would lead to $P_+^{\rm EW} \gg P_+$, in which case the positive-helicity-decay rates of the $\rho^0 K^*$ final states are much larger than the $\rho^\pm K^*$ ones. A complete angular analysis of the ρK^* system should allow a determination of $p_+^{\rm EW}$ even when it is not dominant, possibly allowing a limit on $C^+_{7\gamma}/C^-_{7\gamma}$ of order $r \approx 0.1$.

In conclusion, we discussed an electromagnetic penguin contribution to nonleptonic B decays that has previously been overlooked. It is the largest contribution to the negative-helicity electroweak penguin amplitude and substantially modifies the theoretical expectations for polarization observables in $b \to s$ penguin-dominated decays, in particular, to the $\rho^0 K^*$ final states. These observables may therefore be of considerable interest to the search for electromagnetic flavor-changing neutral currents with chirality equal or opposite to the SM.

We thank G. Buchalla, A. Kagan, and M. Neubert for comments. The work of M. B. is supported by the DFG SFB/TR 9 "Computergestützte Theoretische Teilchenphysik," the work of J. R. by GIF under Grant No. I-781-55.14/2003. D. Y. acknowledges support from the Japan Society for the Promotion of Science.

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