Polarization in charmless $B \rightarrow VV$ decays

Fulvia De Fazio

Istituto Nazionale di Fisica Nucleare, Sezione di Bari, Italy

Abstract. Recent data for $B$ decays to two light vector mesons show that the longitudinal amplitude dominates in $B^0 \rightarrow \rho^+ \rho^-$, $B^+ \rightarrow \rho^0 \rho^-$, and $B \rightarrow \phi K^0$ decays and not in $B \rightarrow \phi K^0$, $B^+ \rightarrow \phi K^+$ transitions. Using factorization, we consider rescattering mediated by charmed resonances, finding that in $B \rightarrow \phi K^0$, it can be responsible for the suppression of the longitudinal amplitude. A similar result is found for $B \rightarrow \rho K$.

Recently, decay widths and polarization fractions of several $B$ decays to two light vector mesons were measured\textsuperscript{[1],[2]}. The branching ratios are of order $10^{-5}$. The measured polarization fractions, collected in Table\textsuperscript{[1]}, show that in penguin induced $B \rightarrow \phi K$ transitions the longitudinal amplitude does not dominate. However, using factorization and the heavy quark limit one can show that the light $V V$ final state should be mainly longitudinally polarized. Actually, the decay $B^0 \rightarrow \phi K^0$ is described by the amplitude $A \left( B^0 (p) \rightarrow \phi (q) ; \epsilon \right) K^0 \left( \rho^0 (\eta) \right) = A_{00} \epsilon \eta + A_2 \left( \epsilon , p \right) (\eta , q) + i A \epsilon \alpha \beta \gamma \delta \epsilon \alpha \eta p \gamma q$, with $\epsilon (q, \lambda)$, $\eta (p, 0, \lambda)$ the $\phi$, $K$ polarization vectors and $\lambda = 0, 1$ the helicities. Since the $B$ meson is spinless, the final mesons have the same helicity. In terms of $A_{0,1,2}$ (describing $S$, $P$, $D$ wave decays, respectively) the helicity amplitudes $A_L$, $A$ read: $A_L = \{ 0 \ \ 0 \ ; \ R \ ; \ L \}$, $A = A_{0,1,2} = (M_0 M_{K_B}, M_0 M_{B_B}, M_B M_B^* \rho)$. The transverse amplitudes are defined as $A_{k\ell} = (\rho_A, A) = 1$, while the polarization fractions are $f_\ell = \hat{\alpha}, \hat{\beta}, \hat{\gamma}, i = L; k; ?$.

Considering the effective weak Hamiltonian inducing the $\bar{b} \rightarrow \bar{s} s \bar{s}$ transitions\textsuperscript{[3]} the amplitude $A \left( B^0 \rightarrow \phi K^0 \right)$ admits a factorized form $A_{fact} (B^0 \rightarrow \phi K^0) = (G_F = \frac{2}{3} V_{ub} V_{ts} a_W) K^0 (p, \eta) \bar{s} (\eta) \bar{s} (\eta) \bar{s} (\eta) \bar{s} (\eta) \bar{s} (\eta) \bar{s} (\eta)$, with $a_W$ a combination of Wilson coefficients. Using $i \phi (q, \epsilon) \bar{s} \gamma^\mu s \bar{s} \gamma^\mu s i = f_\phi M_\phi \epsilon^\mu$, and

\[
\begin{align*}
&\frac{A_2 (q^2)}{(M_B + M_K)} (\eta, p) (p + \bar{p}) (2 M_K) (M_B + M_K) (A_0 (q^2)) (A_0 (q^2)) (\eta, p) (q^2) ; \\
&\frac{M_B + M_K}{(M_B + M_K)} (\eta, p) (p + \bar{p}) (2 M_K) (M_B + M_K) (A_0 (q^2)) (A_0 (q^2)) (\eta, p) (q^2) ; (1)
\end{align*}
\]

$\left(q = p \quad p^0\right)$, one can write the polarization fractions and check that, for large $M_B$: $A_L \propto \frac{M_B}{M_B} \left( A_1 (M_B^2) + A_2 (M_B^2) + A_2 (M_B^2) \right), A_2 \propto M_B M_B \left( M_B^2 \right), A_2 \propto M_B \left( M_B^2 \right)$. For $M_B \rightarrow \infty$, $q^2 = 0$ it was found that: $A_2 = A_1 = V = A_1 = 1$\textsuperscript{[4]}, giving $f_L = f_? = 1$. Using generalized factorization, considering the $a_i$ as effective parameters, one may reproduce the experimental branching ratio, but not the polarization fractions, since the dependence on the $a_i$ cancels in the ratios. Hence, to explain the small $f_L$ one has to look either at finite mass corrections and effects beyond factorization, or at new
TABLE 1. Polarization fractions in charmless $B^0 \to VV$ transitions.

| Mode          | Pol. fraction | Belle [1] | BaBar [2] | Average |
|---------------|---------------|-----------|-----------|---------|
| $B^+ \to \phi K^+$ | $f_L$         | $0.49 \pm 0.13$ | $0.05$ | $0.05$ | $0.46 \pm 0.12$ | $0.03$ | $0.47 \pm 0.09$ |
| $B^0 \to \phi K^0$ | $f_L$         | $0.82 \pm 0.11$ | $0.03$ | $0.05$ | $0.52 \pm 0.02$ | $0.02$ | $0.52 \pm 0.04$ |
| $B^0 \to \rho^0 K^+$ | $f_L$         | $0.52 \pm 0.07$ | $0.05$ | $0.05$ | $0.96 \pm 0.04$ | $0.04$ |
| $B^+ \to \rho^0 \rho^+$ | $f_L$         | $0.95 \pm 0.11$ | $0.02$ | $0.97 \pm 0.03$ | $0.04$ | $0.96 \pm 0.07$ |
| $B^0 \to \rho^- \rho$ | $f_L$         | $0.98 \pm 0.08$ | $0.03$ | $0.02$ | $0.02$ |

The decay $B^0 \to \phi K^0$ can also be induced by rescattering through $B^+ \to D_s^+ \bar{D}_s^- \phi$. Such effects could be sizeable since they involve Wilson coefficients of current-current operators \( \mathcal{O} (1) \), while the coefficients of penguin operators in $B^+ \to \phi K^+$ are $\mathcal{O} (10^{-2})$. Besides, there is no CKM suppression \( \langle \bar{V}_t V_s \rangle \). Vectors $D_s^+ \bar{D}_s^- \phi$ can be estimated using an effective Lagrangian describing the interactions of heavy hadrons with light vector mesons \cite{11}. We write $\bar{D}_s^- \bar{D}_s^- \phi = (G_\phi = \Lambda^2 \mathcal{M}_{D_s}^2) \bar{D}_s^- \bar{D}_s^- \phi$, with $\Lambda \sim 1$, as there is empirical evidence that factorization works in these modes. In the heavy quark limit the above matrix elements involve the Isgur-Wise function $\xi$ and a single quantity $f_{D_s} = f_{D_s}$; we use $\xi(y) = 2(l + y)/2$ and $f_{D_s} = 240$ MeV.

Since the exchanged mesons are off-shell, we write the couplings $g_{\phi} \equiv g_{10} F \psi$, $g_{d0}$ being on-shell couplings and $F \psi = (\Lambda^2 \mathcal{M}_{D_s}^2) = (\Lambda^2 \mathcal{M}_{D_s}^2) \mu$ to satisfy QCD counting rules. The relative sign of rescattering and factorized amplitude is unknown, as well as the role of diagrams involving excitations, hence we fix $\Lambda = 2 \ 3$ GeV analyzing the sum $\mathcal{A} = \mathcal{A}_{\text{fact}} + r \mathcal{A}_{\text{resc}}$ in terms of the parameter $r$. In $\mathcal{A}_{\text{fact}}$ we use the $B^+ \to \phi K^+$ form factors in \cite{12,13}, with Wilson coefficients from \cite{3}. The result is shown in fig.\cite[10]. For the model \cite{12}, $r' = 0.08$ gives the experimental branching ratio and $f_L' = 0.55$. Using the form factors in \cite{13}, we reproduce the branching ratio for $r' = 0.05$, though increasing $f_L$. However, this conclusion depends on the value of the Wilson coefficients. For smaller $a_W$, in both cases a similar long-distance contribution is required, reducing $f_L$.

Our conclusion is that rescattering can modify the helicity amplitudes in penguin dominated modes. On the other hand, such effects are too small to affect $B^+ \to \rho \rho$ decays. Actually, while in the tree diagram in $B^0 \to \rho^+ \rho^-$ the CKM factor $(V_{ub} V_{ud})$ has similar size to that in the rescattering diagrams $(V_{cb} V_{cd})$, the Wilson coefficient in current-current transition is $\mathcal{O} (1)$. We expect to observe FSI effects in colour-suppressed and other penguin induced $B^+ \to VV$ decays. For $B^+ \to \rho^0 K^+$, including the rescattering term, we get $f_L' = 0.7$, i.e. smaller (though compatible within 2-σ) than the datum in Table \cite{1}. Hence, our approach can give a small $f_L B^+ \phi K^+$ at the price of a
FIGURE 1. Dependence of branching ratio and polarization fractions of $B^0 \to \phi K^0$ on the long distance term. $B \to K$ form factors in [12] (left) and [13] (right) are used in the factorized amplitude. $r = 0$ corresponds to absence of rescattering. The three curves in (b) refer to $f_L$ (continuous), $f_\tau$ (dashed) and $f_k$ (dot-dashed). The horizontal lines represent the data for the branching ratio (a) and for $f_L$ (b).

smaller $f_L (B^+ \to \rho K^0)$. For $B^+ \to K^0 \rho^+$ very recent measurements reported: $f_L = 0.79 \pm 0.08 \pm 0.04 \pm 0.02$ [14], $f_L = 0.50 \pm 0.09 \pm 0.05 \pm 0.07$ [15].

Other analyses of non factorizable effects have been proposed [10, 16]; they will not be discussed here. It is only worth mentioning the common conclusion that there seems to be no need of non standard mechanisms to understand the polarization fractions in $B \to VV$ transitions, even though more refined studies are required.

I thank P. Colangelo and T.N. Pham for collaboration. Partial support from the EC Contract No. HPRN-CT-2002-00311 (EURIDICE) is acknowledged.

REFERENCES

1. K.-F. Chen et al. [BELLE Collab.], Phys. Rev. Lett. 91 (2003) 201801; J. Zhang et al. [BELLE Collab.], Phys. Rev. Lett. 91 (2003) 221801; K.Abe et al., [BELLE Collab.], hep-ex/0408141.
2. B. Aubert et al. [BABAR Collab.], Phys. Rev. Lett. 91 (2003) 171802; Phys. Rev. D 69 (2004) 031102; [hep-ex/0408017].
3. A. Ali, G. Kramer and C. D. Lu, Phys. Rev. D 58 (1998) 094009.
4. J. Charles et al., Phys. Rev. D 60 (1999) 014001.
5. See Y. Grossman, Int. J. Mod. Phys. A 19 (2004) 907 and references therein.
6. P. Colangelo, F. De Fazio and T. N. Pham, Phys. Lett. B 597, 291 (2004).
7. P. Colangelo et al., Z. Phys. C 45 (1990) 575.
8. C. Isola et al., Phys. Rev. D 64 (2001) 014029; Phys. Rev. D 68 (2003) 114001.
9. P. Colangelo et al., Phys. Lett. B 542 (2002) 71; Phys. Rev. D 69 (2004) 054023.
10. M. Ladisa et al., hep-ph/0409268; H. Y. Cheng et al., hep-ph/0409317.
11. R. Casalbuoni et al., Phys. Lett. B 292 (1992) 371.
12. P. Colangelo et al., Phys. Rev. D 53 (1996) 3672 [Erratum-ibid. D 57 (1998) 3186].
13. P. Ball, eConf C0304052 (2003) WG101 [arXiv:hep-ph/0306251].
14. B. Aubert [BABAR Collab.], hep-ex/0408093.
15. K. Abe et al. [BELLE Collab.], hep-ex/0408102.
16. A. L. Kagan, Phys. Lett. B 601 (2004) 151; W. S. Hou and M. Nagashima, hep-ph/0408007; H. n. Li and S. Mishima, hep-ph/0411146; H. n. Li, hep-ph/0411305.