$B \rightarrow K_1 \gamma$ and tests of factorization for two-body non leptonic *B* decays with axial-vector mesons

G. Nardulli

Dipartimento di Fisica dell'Università di Bari, Italy Istituto Nazionale di Fisica Nucleare, Sezione di Bari, Italy

T. N. Pham

Centre de Physique Théorique, Centre National de la Recherche Scientifique, UMR 7644, École Polytechnique, 91128 Palaiseau Cedex, France

The large branching ratio for $B \to K_1 \gamma$ recently measured at Belle implies a large $B \to K_1$ transition form factor and large branching ratios for non leptonic B decays involving an axial-vector meson. In this paper we present an analysis of two-body B decays with an axial-vector meson in the final state using naive factorization and the $B \to K_1$ form factors obtained from the measured radiative decays. We find that the predicted $B \to J/\psi K_1$ branching ratio is in agreement with experiment. We also suggest that the decay rates of $B \to K_1 \pi$, $B \to a_1 K$ and $B \to b_1 K$ could be used to test the factorization ansatz.

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I. INTRODUCTION

Our analysis is based on the recent announcement from the Belle collaboration concerning the first measurement of the branching ratio \mathcal{B} for B decay into $K_1(1270) \gamma$ [1]:

$$\mathcal{B}(B^+ \to K_1^+(1270)\gamma) = (4.28 \pm 0.94 \pm 0.43) \times 10^{-5} , \qquad (1)$$

together with an upper bound on $K_1(1400)$:

$$\mathcal{B}(B^+ \to K_1^+(1400)\gamma) < 1.44 \times 10^{-5} \text{ (at 90\% C.L.)}.$$
 (2)

These results should be compared to $B \to K^* \gamma$. The decay fractions measured by CLEO [2], BaBar [3] and Belle [4] Collaborations result in average branching ratios $\mathcal{B}(B^0 \to K^{*0}\gamma) = (4.17 \pm 0.23) \times 10^{-5}$ and $\mathcal{B}(B^+ \to K^{*+}\gamma) = (4.18 \pm 0.32) \times 10^{-5}$.

The large measured $\mathcal{B}(B^+ \to K_1^+(1270)\gamma)$ is a surprise since recent calculations [5, 6, 7, 8] predict a branching ratio smaller than the measured value by a factor ≈ 4 , though a previous calculation [9] gives a larger branching ratio, in the range $(1-4) \times 10^{-5}$, not too far from the measured value. However the small tensor $B \to K_1$ form factor for the radiative decays $B \to K_1 \gamma$ obtained by these recent calculations implies also a tiny branching ratio for non leptonic two-body B decays with axial-vector meson in the final state. Therefore one would expect a small branching ratio for $B \to J/\psi K_1$. This is in contrast with the large measured value [10] for this decay. This value is comparable with the $B \to J/\psi K^*$ branching ratio, which implies that the form factors for the transitions $B \to K_1(1270)$ and $B \to K^*$ should be similar in size in order to explain the large branching ratios for both the radiative and the non leptonic $B \to K_1(1270)$ decays. The aim of the present letter is to present arguments to show that this is indeed the case. We employ naive factorization and the heavy quark symmetry to relate the tensor form factor of the radiative transition to the form factors that describe non-leptonic decays. From the measured radiative decay rates as well as recent data on branching ratios and polarizations for $B \to J/\psi K^*$ decays, we find that the predicted $\mathcal{B}(B \to J/\psi K_1(1270))$ agrees with the experimental results.

From the data (1), (2) we also derive some straightforward predictions for a few non leptonic decay channels involving light strange or non-strange axial-vectors in the final state. This can be achieved by making use only of naive factorization and relations obtained from the Heavy Quark Effective Theory (HQET) [11, 12, 13]. Our approach is therefore in the spirit of the chiral effective theory for heavy mesons, see e.g. [14, 15, 16, 17, 18], and for a review [19], where a similar approach was used to relate a number of decay channels of heavy mesons using the approximate symmetries of HQET.

II. $B \rightarrow K_1$ RADIATIVE DECAYS AND THE MIXING ANGLE

The $K_1(1270)$ and $K_1(1400)$ are strange axial-vector resulting from a mixing of ${}^{3}P_1$ and ${}^{1}P_1$ states. Following PDG [20], we denote by K_{1A} and K_{1B} the ${}^{3}P_1$ and ${}^{1}P_1$ states of K_1 . Thus we have

$$K_1(1270) = K_{1A}\sin\theta + K_{1B}\cos\theta,$$

$$K_1(1400) = K_{1A}\cos\theta - K_{1B}\sin\theta.$$
(3)

The mixing angle θ has been determined up to a fourfold ambiguity, see [8] and, previously, [21]. The masses of K_{1A} and K_{1B} , can be determined by the relations [21]

$$m_{K_{1A}}^2 = m_{K_1(1270)}^2 + m_{K_1(1400)}^2 - m_{K_{1B}}^2$$

$$2m_{K_{1B}}^2 = m_{b_1(1235)}^2 + m_{h_1(1380)}^2 . aga{4}$$

 K_{1B} belongs to the same nonet as the states $b_1(1235)$, $h_1(1170)$ and $h_1(1380)$; K_{1A} , $a_1(1260)$, $f_1(1285)$ and $f_1(1400)$ are also in the same nonet, different from the previous one. Besides (3) we also have

$$\cos 2\theta = \frac{m_{K_{1B}}^2 - m_{K_{1A}}^2}{m_{K_1(1270)}^2 - m_{K_1(1400)}^2} \,. \tag{5}$$

Using (3) and (5) and restricting to $0 < \theta < 90^{\circ}$ we get only two solutions [21]:

Sol.[a]:
$$\theta = 32^{\circ}$$
, $(m_{K_{1B}}, m_{K_{1A}}) = (1310, 1367) \text{ MeV}$,
Sol.[b]: $\theta = 58^{\circ}$, $(m_{K_{1B}}, m_{K_{1A}}) = (1367, 1310) \text{ MeV}$. (6)

These results give a clue for understanding the Belle results. In fact, for any reasonable computational scheme the form factors $T_1(0)$ that determine the radiative decays $B \to K_{1A}\gamma$ and $B \to K_{1B}\gamma$ should be almost identical. This is confirmed by the dynamical calculation of Ref. [8] that gives for this ratio

$$\frac{T_1^{B \to K_{1B}}(0)}{T_1^{B \to K_{1A}}(0)} = 1.2 , \qquad (7)$$

where the form factor is defined by

$$\langle K_1(p',\,\epsilon) | \bar{s}\sigma_{\mu\nu}(1+\gamma_5)q^{\nu}b | B(p) \rangle = i\epsilon_{\mu\nu\rho\sigma}\epsilon^{*\nu}p^{\rho}p'^{\sigma}2 T_1(q^2) + [\epsilon^*_{\mu}(m_B^2 - m_{K_1}^2) - (\epsilon^* \cdot q)(p+p')_{\mu}]T_2(q^2) + [(\epsilon^* \cdot q)q_{\mu} - \frac{q^2}{m_B^2 - m_{K_1}^2}(p+p')_{\mu}]T_3(q^2) .$$

$$(8)$$

Here $T_1(0) = T_2(0)$, while T_3 does not contribute to the radiative decay. A similar definition holds also for $B \to K^*$. We note that, from experiment,

$$y \equiv \frac{T_1^{B \to K_1(1270)}(0)}{T_1^{B \to K^*}(0)} = \sqrt{\left(\frac{m_B^2 - m_{K^*}^2}{m_B^2 - m_{K_1}^2}\right)^3 \frac{\mathcal{B}(B \to K_1(1270)\gamma)}{\mathcal{B}(B \to K^*\gamma)}} \approx 1.06 .$$
(9)

As to $K_1(1400)$ we get

$$\frac{\mathcal{B}(B \to K_1(1400)\gamma)}{\mathcal{B}(B \to K_1(1270)\gamma)} = \left(\frac{m_B^2 - m_{K_1}^2(1400)}{m_B^2 - m_{K_1}^2(1270)}\right)^3 \left|\frac{T_1^{B \to K_{1A}}(0) - \tan\theta \ T_1^{B \to K_{1B}}(0)}{T_1^{B \to K_{1B}}(0) + \tan\theta \ T_1^{B \to K_{1A}}(0)}\right|^2 .$$
(10)

Assuming the value (7) we can predict, from the Belle result (1) the value for $\mathcal{B}(B \to K_1(1400)\gamma)$. The result is in Table I. Both the solutions obtained are in agreement with the upper limit (2).

III. K_1 LEPTONIC DECAY CONSTANT

 K_1 leptonic decay constant can be derived from τ decays. Let us denote by \mathcal{A} a generic axial-vector meson, i.e. one of the following states: K_{1A} , K_{1B} , a_1 , b_1 . We also denote by P, $P^{(\prime)}$ the pseudoscalar mesons, and we use the following definition for the matrix elements of weak currents:

$$\langle 0 | A_{\mu} | P(p) \rangle = i f_P p_{\mu} , \qquad \langle \mathcal{A}(\varepsilon, p) | A_{\mu} | 0 \rangle = f_{\mathcal{A}} m_{\mathcal{A}} \varepsilon_{\mu}^* . \qquad (11)$$

From the $\tau \to K_1$ data [20] we get

$$f_{K_1(1270)} = 171 \text{ MeV}, \qquad f_{K_1(1400)} = 126 \text{ MeV}$$
(12)

Using the mixing angle we derived and SU(3) symmetry we get

Sol.[a]
$$(\theta = 32^0)$$
: $(f_{b_1}, f_{a_1}) = (74, 215) \text{ MeV}$,
Sol.[b] $(\theta = 58^0)$: $(f_{b_1}, f_{a_1}) = (-28, 223) \text{ MeV}$. (13)

We note that these values might be useful to compute weak decays with non strange axial vector in the final state.

IV.
$$B \to K_1 J/\psi$$

For the decay $B \to K^* J/\psi$ and $B \to K_1 J/\psi$ we have the experimental result reported in Table I and we may ask if they are compatible with the Belle result (1).

We use a simple scaling relations, based on HQET, which allows to relate the form factors for the transition $B \to K^*$ via V-A current to those describing transitions by a tensor current. At large q^2 it relates the $A(q^2)$ and $V_1(q^2)$ form factors defined by

$$<\mathcal{A}(\epsilon,p')|V^{\mu} - A^{\mu}|P(p) > = +i(m_{P} + m_{\mathcal{A}})\epsilon^{*\mu}V_{1}(q^{2}) - i\frac{(\epsilon^{*} \cdot q)}{m_{P} + m_{\mathcal{A}}}(p+p')^{\mu}V_{2}(q^{2}) - i(\epsilon^{*} \cdot q)\frac{2m_{\mathcal{A}}}{q^{2}}q^{\mu}\left[V_{3}(q^{2}) - V_{0}(q^{2})\right] - \frac{2A(q^{2})}{m_{P} + m_{\mathcal{A}}}\epsilon^{\mu\nu\alpha\beta}\epsilon^{*}_{\nu}p_{\alpha}p'_{\beta}$$
(14)

where

$$V_3(q^2) = \frac{m_A - m_P}{2m_A} V_2(q^2) + \frac{m_A + m_P}{2m_A} V_1(q^2)$$
(15)

and $V_3(0) = V_0(0)$, with $T_1(q^2)$ in (8) as follows

$$T_1(q^2) = \frac{q^2 + m_B^2 - m_{K_1}^2}{2m_B} \cdot \frac{A(q^2)}{m_B + m_{K_1}} - \frac{m_B + m_{K_1}}{2m_B} V_1(q^2) .$$
(16)

Moreover we assume that the effect of substituting K^* with K_1 is identical in the radiative and in the non leptonic decay, in other words that each form factor for the $B \to K_1$ transition is given by the corresponding form factor for $B \to K^*$ multiplied by the same factor y, once the change of parity between the two strange mesons is taken into account. On this basis we predict

$$\frac{\mathcal{B}(B \to K_1(1270)J/\psi)}{\mathcal{B}(B \to K^*J/\psi)} \frac{\mathcal{B}(B \to K^*\gamma)}{\mathcal{B}(B \to K_1(1270)\gamma)} = \frac{p_{K_1}}{p_{K^*}} \left(\frac{m_B^2 - m_{K^*}^2}{m_B^2 - m_{K_1}^2}\right)^3 \left(x_{\parallel} + x_{\perp} \frac{p_{K_1}^2}{p_{K^*}^2} + x_L \cdot \frac{m_{K^*}^2}{m_{K_1}^2}\right)$$
(17)

Here (we use the BaBar data [22])

$$x_{\parallel} = \frac{\Gamma_{\parallel}(B \to K^* J/\psi)}{\Gamma(B \to K^* J/\psi)} = 0.24 \pm 0.04$$

$$x_{\perp} = \frac{\Gamma_{\perp}(B \to K^* J/\psi)}{\Gamma(B \to K^* J/\psi)} = 0.16 \pm 0.03$$

$$x_L = \frac{\Gamma_L(B \to K^* J/\psi)}{\Gamma(B \to K^* J/\psi)} = 0.60 \pm 0.04$$
(18)

while p_{K^*} (resp. p_{K_1}) is the c.m momentum of K^* (resp. K_1) for the nonleptonic decay $B \to K^* J/\psi$ (resp. $B \to K_1 J/\psi$).

The r.h.s of eq. (18) has the numerical value r.h.s. = 0.64, while

$$l.h.s. = \begin{cases} 0.94 & (neutral mode) \\ 1.30 & (charged mode) \end{cases}$$
(19)

with experimental uncertainties of around 50%. Thus we see that the experimental results for $B \to K_1(1270)\gamma$ and $B \to K_1(1270) J/\psi$ are compatible within the errors. We report in Table I our prediction. Similar arguments apply to the decay $B \to K_1(1400) J/\psi$. Also these results can be found in Table I.

V.
$$B \to K_1 \pi$$

For $B \to K_1 \pi$ decays, if q_{K_1} and q_{K^*} are respectively the c.m. momenta of K_1 and K^* in the reactions $B \to K_1 \pi$ and $B \to K^* \pi$, one gets, using factorization:

$$\frac{\mathcal{B}(B^+ \to K_1^0 \pi^+)}{\mathcal{B}(B^+ \to K^{*\,0} \pi^+)} = \frac{\mathcal{B}(B^0 \to K_1^+ \pi^-)}{\mathcal{B}(B^0 \to K^{*\,+} \pi^-)} = \left(\frac{q_{K_1}}{q_{K^*}}\right)^3 \frac{m_{K^*}^2}{m_{K_1}^2} \left(\frac{F_1^{B \to \pi}(m_{K_1}^2) f_{K_1} m_{K_1}}{F_1^{B \to \pi}(m_{K^*}^2) f_{K^*} m_{K^*}}\right)^2 .$$
(20)

Here we use the form factor F_1 defined by

$$\langle P'(p')|V_{\mu}|P(p)\rangle = F_1(q^2) \left[(p_{\mu} + p'_{\mu}) - \frac{m_P^2 - m_{P'}^2}{q^2} q_{\mu} \right] + F_0(q^2) \frac{m_P^2 - m_{P'}^2}{q^2} q_{\mu}$$
(21)

and a simple pole formula, with a pole mass equal to m_{B^*} , for the q^2 behavior of the $F_1^{B\to\pi}$ form factor. The results obtained by (20) are reported in Table I and represent an interesting test of factorization. It is indeed quite possible that both $B \to K^*\pi$ and $B \to K_1\pi$ decays take non-factorizable contributions from long distance operators formally suppressed in the m_b limit, see e.g. [23], or power corrections in QCD Factorization [24]. In this case the predictions of the last four rows in Table I would get significant violations, pointing to non-factorizable contributions to the decay amplitude.

The reactions with a π^0 in the final state: $B^+ \to K_1^+ \pi^0$ and $B^+ \to K_1^+ \pi^0$ involve two form factors F_1 and V_0 and different combinations of Wilson coefficients and CKM matrix elements. As explained in the introduction the main purpose of this letter is to pick up a few decay channels involving light and strange axial-vector mesons in the final state whose rates can be predicted using only the Belle results [25], τ decay rates and the factorization hypothesis. On this basis we skip these channels leaving a complete analysis to a future paper.

VI. $B \to \mathcal{A}_1 K$

Also in this case we have some clear predictions based on factorization:

$$\frac{\mathcal{B}(B^+ \to a_1^+ K^0)_{\text{fact.}}}{\mathcal{B}(B^+ \to \rho^+ K^0)_{\text{fact.}}} \approx \left(\frac{q_{a_1}}{q_{\rho}}\right)^3 \left(\sin\theta \frac{V_0^{B \to K_1(1270)}(m_K^2)}{A_0^{B \to \rho}(m_K^2)} + \cos\theta \frac{V_0^{B \to K_1(1400)}(m_K^2)}{A_0^{B \to \rho}(m_K^2)}\right)^2 R_+ \quad (22)$$

$$\frac{\mathcal{B}(B^+ \to b_1^+ K^0)_{\text{fact.}}}{\mathcal{B}(B^+ \to \rho^+ K_{\text{fact.}}^0)} \approx \left(\frac{q_{b_1}}{q_{\rho}}\right)^3 \left(\cos\theta \frac{V_0^{B \to K_1(1270)}(m_K^2)}{A_0^{B \to \rho}(m_K^2)} - \sin\theta \frac{V_0^{B \to K_1(1400)}(m_K^2)}{A_0^{B \to \rho}(m_K^2)}\right)^2 R_+$$
(23)

$$\frac{\mathcal{B}(B^0 \to a_1^- K^+)_{\text{fact.}}}{\mathcal{B}(B^0 \to \rho^- K^+)_{\text{fact.}}} \approx \left(\frac{q_{a_1}}{q_{\rho}}\right)^3 \left(\sin\theta \frac{V_0^{B \to K_1(1270)}(m_K^2)}{A_0^{B \to \rho}(m_K^2)} + \cos\theta \frac{V_0^{B \to K_1(1400)}(m_K^2)}{A_0^{B \to \rho}(m_K^2)}\right)^2 R_-$$
(24)

$$\frac{\mathcal{B}(B^0 \to b_1^- K^+)_{\text{fact.}}}{\mathcal{B}(B^0 \to \rho^- K^+)_{\text{fact.}}} \approx \left(\frac{q_{b_1}}{q_{\rho}}\right)^3 \left(\cos\theta \frac{V_0^{B \to K_1(1270)}(m_K^2)}{A_0^{B \to \rho}(m_K^2)} - \sin\theta \frac{V_0^{B \to K_1(1400)}(m_K^2)}{A_0^{B \to \rho}(m_K^2)}\right)^2 R_-$$
(25)

where the subscript means that we consider only factorizable contributions. V_0 has been defined in (14) and, if $|V\rangle$ is a vector meson state,

$$< V(\epsilon, p')|V^{\mu} - A^{\mu}|P(p) > = -i(m_{P} + m_{V})\epsilon^{*\mu}A_{1}(q^{2}) + i\frac{(\epsilon^{*} \cdot q)}{m_{P} + m_{V}}(p + p')^{\mu}A_{2}(q^{2}) + i(\epsilon^{*} \cdot q)\frac{2m_{V}}{q^{2}}q^{\mu}\left[A_{3}(q^{2}) - A_{0}(q^{2})\right] + \frac{2V(q^{2})}{m_{P} + m_{\mathcal{A}}}\epsilon^{\mu\nu\alpha\beta}\epsilon^{*}_{\nu}p_{\alpha}p'_{\beta}$$
(26)

and

$$A_3(q^2) = \frac{m_V - m_P}{2m_V} A_2(q^2) + \frac{m_V + m_P}{2m_V} A_1(q^2)$$
(27)

with $A_3(0) = A_0(0)$; q_{a_1} and q_{b_1} are the c.m. momenta of a_1 and b_1 respectively; the factors R_{\pm} are defined below.

It is a well known fact that factorization terms give small contribution to the decay rates $B^+ \to \rho^+ K^0$, $B^0 \to \rho^- K^+$; for example, for the $B^0 \to \rho^- K^+$ channel, the experimental result $\mathcal{B}(B^0 \to \rho^- K^+) =$ $(7.3 \pm 1.8) \times 10^{-6}$ [20] is larger by one order of magnitude than theoretical predictions based on factorization [23], [26]. This is mainly due to the large cancellation between the penguin contributions appearing in the denominator of the two factors R_{\pm} . These two factors differ by 1 because of the different parity of the vector and axial-vector mesons. The penguin operators O_6 and O_8 distinguish the two parities and therefore

$$R_{+} = \left(\frac{a_{4} - \frac{a_{10}}{2} + \frac{(2 a_{6} - a_{8}) m_{K}^{2}}{(m_{b} - m_{d})(m_{s} + m_{d})}}{a_{4} - \frac{a_{10}}{2} - \frac{(2 a_{6} - a_{8}) m_{K}^{2}}{(m_{b} + m_{d})(m_{s} + m_{d})}}\right)^{2}, \qquad (28)$$

$$R_{-} = \left(\frac{a_4 + a_{10} + \frac{2(a_6 + a_8)m_K^2}{(m_b - m_u)(m_s + m_u)}}{a_4 + a_{10} - \frac{2(a_6 + a_8)m_K^2}{(m_b + m_u)(m_s + m_u)}}\right)$$
(29)

For numerical evaluation of these coefficients we take [27]: $c_3 = 0.013$, $c_4 = -0.029$, $c_5 = 0.009$, $c_6 = -0.033$, $c_7/\alpha = 0.005$, $c_8/\alpha = 0.060$, $c_9/\alpha = -1.283$, $c_{10}/\alpha = 0.266$, with $a_i = c_i + \frac{c_{i-1}}{3}$ (i=even). The other two Wilson coefficients $c_2 = 1.105$, $c_1 = -0.228$ are of no interest here. Moreover, for the current quark masses we use the values $m_b = 4.6$ GeV, $m_u = 4$ MeV $m_d = 8$ MeV, $m_s = 0.150$ GeV. We get therefore

$$R_+ \approx 160, \qquad R_- \approx 80.$$
 (30)

Following the same procedure of Section IV we evaluate the ratio of form factors as follows.

$$\frac{V_0^{B \to K_1}(m_K^2)}{A_0^{B \to \rho}(m_K^2)} \approx \frac{V_0^{B \to K_1}(0)}{A_0^{B \to \rho}(0)} = y \frac{m_{K^*}}{m_{K_1}} \frac{m_B + m_{K_1} - (m_B - m_{K_1})z}{m_B + m_{K^*} - (m_B - m_{K^*})z}$$
(31)

Here y is defined, for $K_1(1270)$ by (9); a similar expression holds for $K_1(1400)$ and $y_{K_1(1400)} = 0.14$ for $\theta = 32^o$ and $y_{K_1(1400)} = 0.35$ for $\theta = 58^o$. The factor z is defined as

$$z = \frac{A_2^{B \to \rho}(0)}{A_1^{B \to \rho}(0)} \approx \frac{A_2^{B \to K^*}(0)}{A_1^{B \to K^*}(0)}$$
(32)

We take the value z = 0.93 intermediate between the value z = 0.9 predicted by light cone sum rules [28] and z = 0.95 given by the BWS model [29]. Although the phase space and the ratio of form factors act as suppressing factors, the big enhancement given by R_{\pm} can produce very large predictions for the decays $B \rightarrow a_1, b_1 K$. As a matter of fact we get for the four ratios in eqns. (22)-(25) results of the order $\approx (59, 76, 29, 38)$ for the solution $\theta = 32^{\circ}$ and $\approx (147, 9, 73, 4)$ for the solution $\theta = 58^{\circ}$. This means that factorization terms give sizeable contributions to these decays, and especially to $B \rightarrow a_1^+ K^0$ for both values of the mixing angle. Our conclusion is that, in view of these results, two-body nonleptonic B decays with a kaon and a light non-strange axial vector meson in the final state represent interesting decay channels with expected large branching ratios. Significant experimental deviations from the the abovementioned ratios would point to specific violations of the factorization model.

It would be tempting to extend the present analysis to the case of B transitions to other orbitally excited K mesons. For example two decay modes with $K_2^*(1430)$ in the final state have been measured: $B \to K_2^*(1430)\gamma$ and $B \to K_2^*(1430) J/\psi$. The radiative transitions $B \to K_2^*(1430)$ have been investigated by some authors. In Ref. [30] HQET is used and the strange quark is treated as heavy, which is however a rather crude approximation. As a result, these authors predict $\mathcal{B}(B \to K_2^*(1430)\gamma)/\mathcal{B}(B \to K_1(1270)\gamma) = 3$, which is at odds with the data, though the predicted $B \to K_2^*(1430)\gamma$ branching ratio is in agreement with experiments. On the contrary in [31] the s quark is considered light. Also this relativistic quark model reproduces correctly the $B \to K_2^*(1430)\gamma$ decay mode, but predicts a too small branching ratio for $B \to K_1(1270)\gamma$. This again brings up the problem with the small predicted radiative branching ratio involving $K_1(1270)$ state, the motivation for the present work. There are also more recent calculations [8] with results in agreement with experiment for the $B \to K_2^*(1430)\gamma$ branching ratios, obtained by various techniques such as light cone sum rules or covariant relativistic quark models as given in the Table V of [8]. An analysis of the $B \to K_2^*(1430) J/\psi$ decay mode is performed in [32]. These authors use the ISGW2 quark model [33], [34] and find results that are however sensitive to the model-dependent form factors. Tests of factorization would be therefore desirable also for these channels. However it must be said that the extension of the present study of $B \to K_1$ transitions to the $B \to K_2^*(1430)$ decay modes cannot be immediate. To study $B \to K_1$ transitions we have used data from $B \to K^*(892)$ transitions and made some further hypotheses, based on the chiral similarity between 1⁻ and 1⁺ states (see the discussion on the $B \to K_1 J/\psi$ channel presented above). For $B \to K_2^*(1430)$ we are in a less favorable situation and some further assumption has to be made. We plan to come back to this issue in the future.

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TABLE I: Theoretical branching ratios and comparison with experiment; [a] and [b] refer to the two possible values of the mixing angle between the K_{1A} and K_{1B} states, $\theta = 32^{\circ}$ and $\theta = 58^{\circ}$ respectively. The experimental values for the processes in lines 1, 2, 3, 5, 9 and 10 are used as inputs.

Process	\mathcal{B} (theory)	exp.
$B^+ \to K^{*+} \gamma$	input	$(4.18 \pm 0.31) \times 10^{-5}$ (av. of [2, 3?])
$B^0 \to K^{*0} \gamma$	input	$(4.17 \pm 0.23) \times 10^{-5}$ (av. of [2, 3?])
$B^+ \to K_1^+(1270)\gamma$	input	$(4.28 \pm 0.94 \pm 0.43) \times 10^{-5}$ [25]
$B^+ \to K_1^+(1400)\gamma$	7.7 × 10 ⁻⁷ [a] 4.4 × 10 ⁻⁶ [b]	$< 1.44 \times 10^{-5}$ [25]
$B^+ \to K^{*+} J/\psi$	input	$(1.35 \pm 0.10) \times 10^{-3}$ [20]
$B^+ \to K_1^+(1270) J/\psi$	0.89×10^{-3}	$(1.8 \pm 0.5) \times 10^{-3}$ [?]
$B^0 \to K_1^0(1270) J/\psi$	0.89×10^{-3}	$(1.3 \pm 0.5) \times 10^{-3}$ [?]
$B^+ \to K_1^+(1400)J/\psi$	1.4×10^{-5} [a] 8.1×10^{-5} [b]	$< 5 \times 10^{-4} \ [20]$
$B^+ \rightarrow K^{*0}\pi^+$	input	$\left(1.9^{+0.6}_{-0.8}\right) \times 10^{-5}[20]$
$B^0 \rightarrow K^{*+}\pi^-$	input	$\left(1.6^{+0.6}_{-0.5}\right) \times 10^{-5} \left[20\right]$
$B^+ \to K_1^0(1270)\pi^+$	1.0×10^{-5}	==
$B^0 \to K_1^+(1270)\pi^-$	0.85×10^{-5}	==
$B^+ \to K_1^0(1400)\pi^+$	0.54×10^{-5}	$< 2.6 \times 10^{-4}$ [20]
$B^0 \to K_1^+(1400)\pi^-$	0.46×10^{-5}	$< 1.1 \times 10^{-3}$ [20]

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