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$\bar{B}^0$ decays to $D^{(*)0}\eta$ and $D^{(*)0}\eta'$

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Abstract. We consider the hadronic B decays $\bar{B}^0 \to D^0\eta$, $\bar{B}^0 \to D^{*0}\eta$, $\bar{B}^0 \to D^0\eta'$, $\bar{B}^0 \to D^{*0}\eta'$ in the framework of a quark-flavour basis and factorization. The formalism allows one to compute the decays to the η meson and to relate them to those of the η'. Measuring the branching ratios of these processes may shed light on the nature of the η–η' mixing. On the experimental side, only upper limits on the branching ratios are known at present.

1 Introduction

The hadronic B decays $\bar{B}^0 \to D^0\eta$, $\bar{B}^0 \to D^{*0}\eta$, $\bar{B}^0 \to D^0\eta'$, $\bar{B}^0 \to D^{*0}\eta'$ are of considerable interest both theoretically and experimentally. Measuring the branching ratios of these processes may shed light on the nature of the η–η' mixing. On the experimental side, only limits on these decays are known at present from CLEO [1]. A theoretical estimate of $\bar{B}^0 \to D^{(*)0}\eta$, $\bar{B}^0 \to D^{(*)0}\eta'$ is important also because these processes are a background for the semi–inclusive processes $\bar{B}^0 \to X_s\eta$, $\bar{B}^0 \to X_s\eta'$. For a recent experimental analysis, see [2].

The processes $\bar{B}^0 \to D^{(*)0}\eta$, $\bar{B}^0 \to D^{(*)0}\eta'$ proceed via the internal spectator diagram shown in Fig. 1. They are colour suppressed since the colour of the quarks from the virtual W has to match the colour of the produced c quark from $b \to c$ and the spectator antiquark coming from the B meson. Here we consider the quark component η_q of the η particle in the so-called quark-flavour basis [3]:

$$\begin{pmatrix} \eta \\ \eta' \end{pmatrix} = \begin{pmatrix} \cos\phi & -\sin\phi \\ \sin\phi & \cos\phi \end{pmatrix} \begin{pmatrix} \eta_q \\ \eta_s \end{pmatrix}, \quad (1)$$

where the Fock state description of $\eta_{q,s}$ is

$$|\eta_q\rangle = \Psi_q \frac{|u\bar{u} + d\bar{d}\rangle}{\sqrt{2}} + \cdots \quad (2)$$

$$|\eta_s\rangle = \Psi_s |s\bar{s}\rangle + \cdots \quad (3)$$

The ellipses indicate higher Fock states, like glue states, while the wave functions $\Psi_{q,s}$ are the light-cone wave functions of the parton states. Our assumption is that the η_q component is the one to be considered in the factorization diagram resulting from Fig. 1 (shrinking the W propagator to a point). We will exploit the relation between η, η' and η_q to compute the amplitude for the process $\bar{B}^0 \to D^{(*)0}\eta'$.

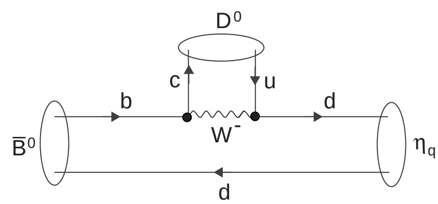

Fig. 1. Flavour diagram for $\bar{B}^0 \to \eta(\eta')$. The η_q component allows one to bridge between the two processes

We recall the form factor parametrisation:

$$\langle\eta(q_\eta)|V^\mu(q)|\bar{B}^0(p)\rangle = \left[(p+q_\eta)^\mu + \frac{m_\eta^2 - m_B^2}{q^2}q^\mu\right]F_1(q^2) - \left[\frac{m_\eta^2 - m_B^2}{q^2}q^\mu\right]F_0(q^2), \quad (4)$$

with $F_1(0) = F_0(0)$. Moreover we need the matrix elements

$$\langle\mathrm{VAC}|V^\mu|D(p)\rangle = \mathrm{i} f_D p^\mu, \quad (5)$$

$$\langle\mathrm{VAC}|V^\mu|D^*(\epsilon,p)\rangle = \epsilon^\mu m_{D^*} f_{D^*}, \quad (6)$$

and we will use the values [4] $f_D = 200\,\mathrm{MeV}$, $f_{D^*} = 230\,\mathrm{MeV}$.

Adopting the Bauer–Stech–Wirbel approach to factorization [5] we can consider the diagram in Fig. 1 describing the η coupling through its $d\bar{d}$ component. The amplitude we need to compute is

$$G_{\bar{B}^0 D^0 \eta} = \langle\eta D^0|H^{\mathrm{BSW}}|\bar{B}^0\rangle = \cos\phi\langle\eta_q D^0|H^{\mathrm{BSW}}|\bar{B}^0\rangle = \frac{G_\mathrm{F}}{\sqrt{2}} V_{cb}V_{ud}^* a_2 f_D F_0^{(B\to\eta_q)} \times (m_D^2)(m_B^2 - m_{\eta_q}^2)\frac{\cos\phi}{\sqrt{2}}, \quad (7)$$

where ϕ is determined in [3] to be $\phi = 39.3°$ and a_2 is the phenomenological coefficient of the Bauer–Stech–Wirbel effective Hamiltonian $a_2 = 0.29$. It is significantly larger than the leading-order result $a_2^{\rm LO} \sim 0.12$ corresponding to naive factorization. Such a large correction is not surprising. The QCD factorization formula of [6] cannot be used to compute rigorously a_2 in these decays; however, a rough estimate in the limit $m_c \ll m_b$ gives $a_2 \sim 0.25$ [7]. The factor $1/2^{1/2}$ in (7) accounts for the fact that only the $d\bar{d}$ component enters in our diagram. The strange quark component is absent.

The amplitude $G_{\bar{B}^0 D^0 \eta'}$ is then connected to (7) through

$$G_{\bar{B}^0 D^0 \eta'} = \tan\phi G_{\bar{B}^0 D^0 \eta}. \tag{8}$$

In an analogous way we find

$$\begin{aligned} G_{\bar{B}^0 D^{*0}\eta} &= \langle \eta D^{*0} | H^{\rm BSW} | \bar{B}^0 \rangle = \cos\phi \langle \eta_q D^{*0} | H^{\rm BSW} | \bar{B}^0 \rangle \\ &= \frac{G_F}{\sqrt{2}} V_{cb} V_{ud}^* a_2 f_{D^*} m_{D^*} F_1^{(B\to\eta_q)}(m_{D^*}^2)\epsilon \\ &\quad \cdot (p + q_{\eta_q}) \frac{\cos\phi}{\sqrt{2}}. \end{aligned} \tag{9}$$

The term $\epsilon \cdot (p + q_\eta)$ is easily evaluated summing over the polarisations of D^* once the square modulus of this amplitude is considered.

2 NS model

In order to get an idea about the amount of model dependence, we will consider two different models for the form factors. We start by considering a model by Neubert and Stech (NS) [4] which is a parametrisation based on simple assumptions, in reasonable agreement with the available data and with most theoretical predictions. In the case of heavy-to-light transitions, the form factors no longer obey the symmetry relations valid in the heavy-to-heavy decays. One can account for this by introducing, for each form factor, a function $\xi_i(w)$ replacing the Isgur–Wise function:

$$F_1(q^2) = \frac{m_B + m_M}{2\sqrt{m_B m_M}} \xi_1(w), \tag{10}$$

$$F_0(q^2) = \frac{2\sqrt{m_B m_M}}{m_B + m_M} \frac{w+1}{2} \xi_0(w), \tag{11}$$

where

$$w = v_B \cdot v_M = \frac{m_B^2 + m_M^2 - q^2}{2 m_B m_M}. \tag{12}$$

The maximum value of $w = w_{\rm max}$ is obtained at $q^2 = 0$. For an estimate of the functions $\xi_i(w)$ a simple pole model is used:

$$\xi_1(w) = \sqrt{\frac{2}{w+1}} \frac{1}{1+r} \frac{w_{\rm max} - w_1}{w - w_1}, \tag{13}$$

$$\xi_0(w) = \sqrt{\frac{2}{w+1}} \frac{1}{1 + r \dfrac{w-1}{w_{\rm max}-1}} \tag{14}$$

Table 1. Values of the form factors used in the calculation. These are used in slightly different formulas for computing the amplitudes. The CQM form factors are indeed extracted using a Goldstone-like coupling for the η particle and refer only to the $d\bar{d}$ component, while those of the NS model refer to the $u\bar{u} + d\bar{d}$ (see Sect. 3). In order to compare the two models on the same footing the form factors of the CQM model have to be multiplied by $2^{1/2}$

	$F_0^{B\to\eta}(0)$	$F_0^{B\to\eta}(m_D^2)$	$F_1^{B\to\eta}(m_{D^*}^2)$
CQM model	0.15	0.17	0.23
NS model	0.26	0.28	0.33

Table 2. Theoretical predictions and 90% C.L. upper limits on branching ratios. The errors quoted for the CQM model only refer to the variation of the parameters in the model and not to the model dependence of the result

Decay mode	@90% C.L. [1]	NS model	CQM model
$\bar{B}^0 \to D^0\pi$	$< 1.2\times10^{-4}$	0.77×10^{-4}	$1.3^{+0.4}_{-0.3} \times 10^{-4}$
$\bar{B}^0 \to D^{*0}\pi$	$< 4.4\times10^{-4}$	1.05×10^{-4}	$1.1 \pm 0.3 \times 10^{-4}$
$\bar{B}^0 \to D^0\eta$	$< 1.3\times10^{-4}$	0.50×10^{-4}	$0.44 \pm 0.02 \times 10^{-4}$
$\bar{B}^0 \to D^{*0}\eta$	$< 2.6\times10^{-4}$	0.60×10^{-4}	$0.70 \pm 0.04 \times 10^{-4}$
$\bar{B}^0 \to D^0\eta'$	$< 9.4\times10^{-4}$	0.32×10^{-4}	$0.30 \pm 0.02 \times 10^{-4}$
$\bar{B}^0 \to D^{*0}\eta'$	$< 14\times10^{-4}$	0.41×10^{-4}	$0.47 \pm 0.04 \times 10^{-4}$

where

$$r = \frac{(m_B - m_V)^2}{4 m_B m_V} \left(1 + \frac{4 m_B m_V}{M_0^2 - (m_B - m_V)^2} \right), \tag{15}$$

where V is the vector meson in the same doublet as the scalar meson involved in the decay described by the F_0 form factor (for the η this corresponds to ω). Moreover

$$w_1 = \frac{m_B^2 + m_M^2 - M_1^2}{2 m_B m_M}, \tag{16}$$

where M_i is the mass of the nearest resonance with the appropriate spin-parity quantum numbers (for $M_0 = M_{B^{**}} = 5.754\,{\rm GeV}$ it is the 0^+ pole and for $M_1 = M_{B^*} = 5.325\,{\rm GeV}$ it is the 1^- pole).

These form factors satisfy the relation

$$F_1(0) = F_0(0), \tag{17}$$

and scale as $\xi_i(w) \sim w^{-3/2}$ for large w, in accordance with the scaling rules obtained in [8].

The $B \to \eta$ form factors we obtain in this model are given in Table 1. The results for the branching ratios are given in Table 2. We also compute the corresponding $B \to \pi$ decays in order to allow a comparison to be made with experimental data and theoretical predictions that is independent of the assumptions concerning the η mixing.

3 CQM model

The computation of the form factor F_0 and F_1 can be carried out with the aid of a constituent quark model

(CQM) [9]. This model has shown to be particularly suitable for the study of heavy meson decays. Since its Lagrangian describes Feynman rules for the vertices (heavy meson)–(heavy quark)–(light quark) [10], transition amplitudes are computable via simple constituent quark loop diagrams in which mesons appear as external legs. We determine $F_0^{B\to\eta_8}$ considering the η_8 as a Goldstone boson. Therefore one has to take into account the relations between η_8 and η_q for the process under study. To begin we observe that

$$q^\mu\langle\eta_8|V^\mu|\bar{B}^0\rangle = (m_B^2 - m_\eta^2)F_0^{(B\to\eta_8)}, \tag{18}$$

where q^μ is the momentum carried by the current V^μ. For our purposes η is essentially the Goldstone octet component η_8:

$$\langle\eta|V^\mu|\bar{B}^0\rangle \simeq \langle\eta_8|V^\mu|\bar{B}^0\rangle. \tag{19}$$

On the other hand the amplitude related to the process in Fig. 1 is

$$\langle\eta|V^\mu|\bar{B}^0\rangle = \cos\phi\langle\eta_q|V^\mu|\bar{B}^0\rangle, \tag{20}$$

if we also take into account a small mixing angle, θ_8 in the notation of [3], we write

$$\langle\eta_8|V^\mu|\bar{B}^0\rangle = \frac{\cos\phi}{\cos\theta_8}\langle\eta_q|V^\mu|\bar{B}^0\rangle. \tag{21}$$

Again we stress that here we are neglecting the singlet component and we assume that the $B\to\eta$ matrix element which enters in the factorization diagram is the one containing only the $d\bar{d}$ components. In other words we can write

$$q^\mu\langle\eta_8|V^\mu|\bar{B}^0\rangle = \frac{\cos\phi}{\cos\theta_8}(m_B^2 - m_{\eta_q}^2)F_0^{(B\to\eta_q)}, \tag{22}$$

in such a way that

$$F_0^{(B\to\eta_q)} = \frac{\cos\theta_8}{\cos\phi}\frac{(m_B^2 - m_\eta^2)}{(m_B^2 - m_{\eta_q}^2)}F_0^{(B\to\eta_8)}. \tag{23}$$

Notice that θ_8 is determined in [3] to be $\theta_8 = -21.0°$. For the F_1 form factor we obtain

$$F_1^{(B\to\eta_q)}(m_{D^*}^2) = \frac{\epsilon\cdot(p+q_{\eta_8})}{\epsilon\cdot(p+q_{\eta_q})}\frac{\cos\theta_8}{\cos\phi}F_1^{(B\to\eta_8)}(m_{D^*}^2). \tag{24}$$

We compute the η particle via its η_8 component with a Goldstone-like coupling in the meson-quark loop computation. The interaction Lagrangian generating this coupling has the form

$$\bar{\psi}_a\gamma\cdot\mathbf{A}_{ab}\gamma_5\psi_b, \tag{25}$$

where the indices a, b are the u, d, s quark indices, ψ being a triplet of flavour-SU_3. The structure of the $\mathbf{A}$ matrix is

$$\mathbf{A}_\mu = -\frac{\mathrm{i}}{2}(\xi\partial_\mu\xi^\dagger - \xi^\dagger\partial_\mu\xi), \tag{26}$$

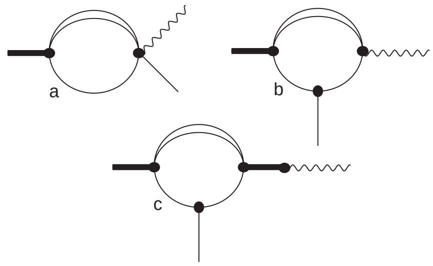

Fig. 2a–c. CQM diagrams. The heavy line represents the incoming heavy meson. The double line is the heavy quark, the curly line represents the current insertion. **a** is the non-derivative diagram in which the η is coupled at the same vertex as the current, **b** is the direct diagram and **c** is the polar diagram: an intermediate heavy meson state is taken into account

and $\xi = \mathrm{e}^{\mathrm{i}\pi/f}$. The π matrix has the well-known form

$$\pi = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}} & \pi^+ & K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}} & K^0 \\ K^- & \overline{K^0} & -\frac{2}{\sqrt{6}}\eta_8 \end{pmatrix}. \tag{27}$$

When the $d\bar{d}$ component is taken into account the Feynman rule for the vertex is $q^\mu\gamma_\mu\gamma_5/(6^{1/2}f_\pi)$. We isolate three contributions to the $F_{0,1}$ form factors: a *non-derivative* contribution, a *direct* and a *polar* contribution [11]; see Fig. 2. The direct and the non-derivative contributions are represented in the CQM model by loop diagrams in which the internal lines are heavy quark and light quark propagators, while the external legs contain the incoming meson and the vector current. In the non-derivative diagram the η external leg is attached at the same vertex (heavy quark)–(light quark)–(current) due to the structure of the CQM interaction Lagrangian. In the direct diagram the η is instead attached to the light quark internal line (see [11]). The polar diagram allows for an intermediate polar state separating the current insertion from the loop diagram. Each of these diagrams contributes to the form factors $F_{0,1}$ in a calculable way. Summing up these contributions and imposing the condition $F_1(0) = F_0(0)$, which eliminates the spurious singularity in the form factor decomposition, one obtains

$$F_i(q^2) = \alpha_j(q^2)(F_i^{\mathrm{nd}} + F_i^{\mathrm{dir}}(q^2) + F_i^{\mathrm{pol}}(q^2)), \tag{28}$$

where $i = 0, 1$. The functions $\alpha_j(q^2)$, $j = 0, 1$ are needed to impose the condition $F_1(0) = F_0(0)$. This is because the sum in parentheses reproduces the form factor on the l.h.s. of the previous equation only at q^2_{max}, i.e., $\alpha(q^2_{\mathrm{max}}) \simeq 1$. Following [11] we choose

$$\alpha_j(q^2) = 1 + \alpha_j\frac{q^2 - m_B^2 - m_\eta^2}{2m_B\Lambda_\chi} = 1 - \alpha_j E_\eta/\Lambda_\chi, \tag{29}$$

where $\Lambda_\chi = 1\,\mathrm{GeV}$. In such a way we enforce that $\alpha(q^2_{\mathrm{max}}) \simeq 1$. E_η is the η energy in the B rest frame. We borrow α_0

from [11] where we had $\alpha_0 = 0.27$. The form factor used in the computation of the branching ratios are of course those given in (28). The non-derivative (nd) contributions to the form factors are given by

$$F_0^{\rm nd} = \frac{f_B}{\sqrt{6}f_\pi}, \quad F_1^{\rm nd} = \frac{f_B}{2\sqrt{6}f_\pi}. \tag{30}$$

The direct contributions are

$$F_1^{\rm Dir}(q^2) = \frac{S_1+S_2}{2}, \tag{31}$$

$$F_0^{\rm Dir}(q^2) = \frac{S_1 - S_2 + A(q^2)(S_1+S_2)}{2A(q^2)}, \tag{32}$$

where

$$A(q^2) = \frac{(m_B^2 - m_\eta^2)}{q^2},$$

and

$$S_1 = \frac{2}{\sqrt{6}f_\pi}\sqrt{Z_H m_H}\left(\frac{m m_\eta^2}{m_H}Z + \left(\frac{2m}{m_H}v\cdot q_\eta - \frac{m_\eta^2}{m_H}\right)\Omega_1 - \frac{2v\cdot q_\eta}{m_H}\Omega_4 - \frac{2m_\eta}{m_H}\Omega_6\right),$$

$$S_2 = \frac{2}{\sqrt{6}f_\pi}\sqrt{Z_H m_H}\Big(m^2 Z - 2m\Omega_1 - m_\eta \Omega_2 + 2\Omega_3 + \Omega_4 - \Omega_5\Big).$$

In these expressions we have

$$v\cdot q_\eta = \frac{m_B^2 + m_\eta^2 - q^2}{2m_H}. \tag{33}$$

The symbols Z, Z_H, Ω_i indicate integral expression listed in [10]. These are functions of Δ_H, the difference between the heavy meson and the heavy quark constituent mass $m_H - m_Q$, $\Delta = \Delta_H - v\cdot q_\eta$ and $v\cdot q_\eta$, v being the four velocity of the B meson. q^μ is the momentum carried by the current insertion in the CQM loop diagram, m is the constituent light quark mass, $m = 300\,\text{MeV}$. The suffix H is referred to the $H = (0^-, 1^-)$ multiplet predicted by heavy quark effective theory.

The polar contributions are

$$F_1^{\rm pol}(q^2) = \frac{\hat{F}g}{\sqrt{6}f_\pi\sqrt{m_B}}\frac{1}{1-q^2/m_{B^*}^2}, \tag{34}$$

$$F_0^{\rm pol}(q^2) = \frac{1}{m_B^2 - m_\eta^2}\left(\frac{h m_\eta \sqrt{m_B}\hat{F}^+}{\sqrt{6}f_\pi}\right)\frac{1}{1-q^2/m_{B^{**}}^2}, \tag{35}$$

where B^{**} is the 0^+ state of the multiplet $S = (0^+, 1^+)$, while B^* is the 1^- state of the H multiplet. h is the coupling constant $HS\pi$ and in CQM one finds $h = -0.76 \pm 0.13$ while $g = 0.46 \pm 0.04$, from $HH\pi$. We use the central values for all these constants and the value $\Delta_H = 0.4\,\text{GeV}$. Variations of Δ_H in the range 0.3–0.5 GeV induce few % variations in the form factors (see Table 2). $\hat{F}$ and $\hat{F}^+$ are defined as the decay constants of the H and S heavy mesons respectively.

Note that the form factors calculated in this section already include the factor taking into account that the process we are considering only involves the $d\bar{d}$ component of η through the numerical coefficient of the Feynman rule. Therefore when combining (7) and (23) one has to be careful to remove a factor $1/(2^{1/2})$ from (7):

$$G_{\bar{B}^0 D^0 \eta} = \frac{G_{\rm F}}{\sqrt{2}} V_{cb}V_{ud}^* a_2 f_D F_0^{(B\to\eta_8)}(m_D^2)(m_B^2 - m_\eta^2)\cos\theta_8 \tag{36}$$

in terms of the η_8 state. From (24) and (9) we get in a similar way

$$G_{\bar{B}^0 D^{*0}\eta} = \frac{G_{\rm F}}{\sqrt{2}} V_{cb}V_{ud}^* a_2 f_{D^*} m_{D^*} \times F_1^{(B\to\eta_8)}(m_{D^*}^2)\epsilon\cdot(p+q_\eta)\cos\theta_8. \tag{37}$$

Numerical results are given in Table 1 and 2. The two models give similar results for the form factors and branching ratios. The amount of model dependence due to the form factors is difficult to quantify. By comparing the two models one can conclude that they affect the branching ratios at the level of $\sim \pm 0.1 \times 10^{-4}$. The a_2 coefficient is considered fixed at the value 0.29 in the calculation. As it enters as a_2^2 in the branching ratio, changing its value to 0.25 may add to the error on the branching ratio up to $\sim \pm 0.2 \times 10^{-4}$. However if the quark picture used in the calculation is correct, such a large variation is not expected as $\bar{B}^0 \to D^0\eta$ is in this case to be treated on the same footing as $\bar{B}^0 \to D^0\pi$.

4 Conclusion

In order to test the idea of relating the η and η' using the quark-flavour basis in the decays considered in this work, experimental data will have to measure the ratio of Br's η'/η. If different from what estimated, this probably means that the quark mechanism (see Fig. 1) is not sufficient to explain the decays into η'. This would suggest we should look at some other mechanism, like the gluon anomaly explored in [12] where the large observed production of η' in $B \to X_s\eta'$ decays has been studied. The mechanism suggested there is based on the subprocess $b \to sg^* \to s\eta' g$ where the virtual gluon emerging from the standard model penguin couples to η' via a gluon anomaly vertex $g^* g\eta'$. The possibility that the other gluon g is emitted by the light quark inside a B meson has been examined in [13]. The $B \to K\eta'$ decay has been recently examined in the context of perturbative QCD in [14]. In the case of $B \to D\eta'$ studied here it is more difficult to imagine some gluonic mechanism of the kind described before. Anyway we cannot exclude this possibility since the η' can have a large glue component [13]; see also [15]. Under the assumptions considered in this work $\bar{B}^0 \to D^0\eta$

is similar to $\bar{B}^0 \to D^0\pi$. The calculation is related to the one for $\bar{B}^0 \to D^0\eta'$ using the quark-flavour basis. We have checked that two different ways of computing the form factors give similar results.

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Note added in proof: Very recently there was a phenomenological determination of the a_2 coefficient [16] based on new data from Belle [17]. This seems to be in contrast with all previous determinations based on the factorization approach [6] addressing the problem of final-state interaction in the $B \to D\pi$ channel. Such a problem could arise also in this analysis where the $B \to D\eta$ is considered as the starting point, and the factorization value for a_2 is assumed. We will consider this aspect in future work.