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Branching ratios for $B_{d,s} \to J/\psi\,\eta$ and $B_{d,s} \to \eta\,\ell^+\ell^-$, extracting γ from $B_{d,s} \to J/\psi\,\eta$, and possibilities for constraining C_{10A} in semileptonic B decays

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Abstract: Estimates of the branching ratios for $B_{d,s} \to J/\psi\,\eta$ and $B_{d,s} \to \eta\,\ell^+\ell^-$ are obtained by SU(3) relation to $B_d \to J/\psi\,K$ and $B_d \to K\ell^+\ell^-$, respectively, as functions of the $\eta_1 - \eta_8$ mixing parameter, θ_P. Based on these estimates, a discussion of the prospects for HERA-B, CDF-II, and ATLAS on these processes is given. The CP violation in $B_{d,s} \to J/\psi\,\eta$ is analyzed in depth and a method to extract the angle γ of the unitarity triangle is discussed. Finally, a possible method to constrain the Wilson coefficient C_{10A} from measurements on semileptonic B decays such as $B_d \to K\ell^+\ell^-$ and $B_s \to \eta\,\ell^+\ell^-$ is proposed along with a discussion of the prospects for future experiments and form factor calculations to reach the precision required for this method to be interesting

Keywords: Rare Decays, B-Physics, CP violation.

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Contents

1. Introduction

One of the most important processes relevant to the study of CP violation (see [1] for an excellent review of this topic) is the well-known "gold-plated mode" $B_d \to J/\psi K_S$. Besides its usefulness in extracting $\sin(2\beta)$ (see e.g. [2]), it has relatively recently been argued that this process can also be used to extract the angle γ when combined with its U-spin partner $B_s \to J/\psi\,K_S$ [3].

In this paper, an alternative to the gold-plated mode for the extraction of γ is presented: $B_{d,s} \to J/\psi\,\eta$. Due to the inherent problems in adapting factorization to non-leptonic decays we use the completely general and model independent method of quark topologies to analyze the structure of $B_{d,s} \to J/\psi\,\eta$, obtaining a relation between the B_d and B_s amplitudes through SU(3) flavour symmetry (the approximate symmetry of u, d, s). The structure of the two processes is such that the phase, $e^{i\gamma}$, is CKM suppressed in $B_s \to J/\psi\,\eta$ relative to $B_d \to J/\psi\,\eta$, and so effects of CP violation will be more easily visible in the $B_d \to J/\psi\,\eta$ decay. However, the extraction of γ from this process is plagued by the appearance of a normalization factor which cannot be determined directly. Through U spin relation, the CP averaged rates of $B_s \to J/\psi\,\eta$ and $B_d \to J/\psi\,\eta$ can be combined to fix this normalization. Thus, γ can be determined with a theoretical uncertainty depending only on SU(3) breaking corrections and the $\eta_1 - \eta_8$ mixing angle, θ_P. Furthermore, the B_d amplitude itself is suppressed relative to the B_s amplitude, the inverse of the situation in the $B_{d,s} \to J/\psi\, K_S$ case. The difference in production rates of B_d and B_s mesons in the experiment compensates to some extent for this difference in the case $B_{d,s} \to J/\psi\, \eta$ whereas it worsens the situation for $B_{d,s} \to J/\psi\, K_S$.

As an addendum, a method to constrain C_{10A} is presented. In semileptonic decays, C_{10A} describes the effective coupling of the axial OPE operator $\mathcal{O}_{10A}$ [1]. In models of new physics, its value generally deviates from the SM prediction, due to virtual particle contributions from New Physics particles present in such models. In this paper, a method is proposed which allows the elimination of large hadronic uncertainties caused by the presence of intermediate ψ resonances by measuring distributions in semileptonic decays and using relations between them. It is, however, doubtful whether this method can find immediate application due to the high precision required both for theoretical and experimental input.

The outline of the paper is as follows: in section 2, we use the method of quark topologies to analyze the CKM structure of the contributions to the $B_{d,s} \to J/\psi\, \eta$ decay amplitude (see [4] for a recent update on this method). Noting the Zweig and SU(3) suppressions of certain topologies and assuming SU(3) symmetry of the strong interaction dynamics, a simple relation can be obtained between the $B_d \to J/\psi\, \eta$ and $B_s \to J/\psi\, \eta$ amplitudes. In section 3, the procedure proposed in [3] for extracting γ from $B_{d,s} \to J/\psi\, K_S$ is adapted for $B_{d,s} \to J/\psi\, \eta$, and the amplitude relation obtained in the previous section is used to obtain the normalization of the CP averaged $B_d \to J/\psi\, \eta$ rate.

In section 4, a simplified picture of the $B_{d,s} \to J/\psi \eta$ process is adopted and SU(3) symmetry is invoked to obtain a relation between the amplitudes of $B_d \to J/\psi\, \eta$, $B_s \to J/\psi\, \eta$ and $B_s \to J/\psi\, K_S$ to the already measured $B_d \to J/\psi\, K_S$. In section 5, essentially the same is employed to obtain $B_{d,s} \to \eta$ form factors from those for $B_d \to K$ calculated by Light Cone Sum Rules (LCSR) in [5, 6].[1] Using these, an estimate for the $B_{d,s} \to \eta\, \ell^+\ell^-$ branching ratios can be obtained. Both the $B_{d,s} \to J/\psi\, \eta$ and the $B_{d,s} \to \eta\, \ell^+\ell^-$ branchings depend on the degree of octet-singlet mixing in the η system, expressed through the mixing angle, θ_P.

In section 6, a simple method is proposed to constrain C_{10A}, the axial semileptonic Wilson coefficient. We replace the theoretically poorly known quantity $C_{9V}^{\rm eff}$ in the $B_d \to K\tau^+\tau^-$ amplitude by measurable decay distributions for the $B_d \to K\mu^+\mu^-$ process, yielding C_{10A} as a function of the total branching to $K\tau\tau$, the differential branching to $K\mu\mu$, $|V_{ts}^* V_{tb}|^2$, and the form factors f_+ and f_-. Estimates in various SUGRA models and the 2 Higgs doublet model are considered, and the application of the same procedure to the case $B_s \to \eta\, \ell^+\ell^-$ is considered. Concluding remarks and outlook are given in section 7.

[1] For a more recent calculation, see [7].

2. Analysis of $B_{d,s} \to J/\psi\,\eta$

The time-independent transition amplitudes for B^0 and $\overline{B}^0$ states into the final CP eigenstate, f_{CP}, are parametrized as in [3]:

$$\langle\, f_{CP}\,|H|\,B_q^0\,\rangle \equiv \mathcal{A}_q = N_q\left[1 - a_q e^{i\theta_q} e^{i\gamma}\right] \equiv N_q z_q\,, \tag{2.1}$$

$$\langle\, f_{CP}\,|H|\,\overline{B}_q^0\,\rangle \equiv \overline{\mathcal{A}}_q = \eta N_q\left[1 - a_q e^{i\theta_q} e^{-i\gamma}\right] \equiv \eta N_q \overline{z}_q\,, \tag{2.2}$$

where η is the CP eigenvalue of f_{CP}, θ is a strong phase, and $q \in \{d, s\}$. The amplitude at time t for an initial $B/\overline{B}$ meson to decay then becomes:

$$\begin{aligned}|A_q(t)|^2 = \frac{|N_q|^2}{2}\Big[&(R_L + \varepsilon_{pq} R_L^1)e^{-\Gamma_L t} + (R_H + \varepsilon_{pq} R_H^1)e^{-\Gamma_H t} + \\ &+ 2e^{-\Gamma t}\left[(A_D + \varepsilon_{pq} A_D^1)\cos(\Delta M t) + A_M\left(1 - \frac{\varepsilon_{pq}}{2}\right)\sin(\Delta M t)\right]\Big],\end{aligned} \tag{2.3}$$

$$\begin{aligned}|\overline{A}_q(t)|^2 = \frac{|N_q|^2}{2}\Big[&(R_L + \varepsilon_{pq} \overline{R}_L^1)e^{-\Gamma_L t} + (R_H + \varepsilon_{pq} \overline{R}_H^1)e^{-\Gamma_H t} - \\ &- 2e^{-\Gamma t}\left[(A_D + \varepsilon_{pq} \overline{A}_D^1)\cos(\Delta M t) + A_M\left(1 + \frac{\varepsilon_{pq}}{2}\right)\sin(\Delta M t)\right]\Big],\end{aligned} \tag{2.4}$$

where we define rate functions as in [3], except that we here give the formulae to first order in the small parameter $\varepsilon_{pq} \equiv |p/q|^2 - 1$, where p and q are the standard factors parametrizing the B meson mass eigenstates in terms of flavour states [1]. As there is, however, only small possibility for new physics to result in a large ε_{pq} we henceforth ignore this parameter. The rate functions entering the above are defined by (ϕ_q is the $B_q - \overline{B}_q$ mixing phase):

$$\begin{aligned}
R_L &\equiv \frac{1}{2}\left(|z_q|^2 + |\overline{z}_q|^2 + 2\eta\,\mathrm{Re}\{z_q^*\overline{z}_q e^{i\phi_q}\}\right)\\
R_H &\equiv \frac{1}{2}\left(|z_q|^2 + |\overline{z}_q|^2 - 2\eta\,\mathrm{Re}\{z_q^*\overline{z}_q e^{i\phi_q}\}\right)\\
A_D &\equiv \frac{1}{2}(|z_q|^2 - |\overline{z}_q|^2)\\
A_M &\equiv -\eta\,\mathrm{Im}\{z_q^*\overline{z}_q e^{i\phi_q}\}\\
R_L^1 &\equiv -\frac{1}{2}\left(|\overline{z}_q|^2 + \eta\,\mathrm{Re}\{z_q^*\overline{z}_q e^{i\phi_q}\}\right) = -\frac{1}{2}(R_L + A_D)\\
R_H^1 &\equiv -\frac{1}{2}\left(|\overline{z}_q|^2 - \eta\,\mathrm{Re}\{z_q^*\overline{z}_q e^{i\phi_q}\}\right) = -\frac{1}{2}(R_H - A_D)\\
A_D^1 &\equiv \frac{1}{2}|\overline{z}_q|^2 = \frac{1}{2}\left(\frac{1}{2}(R_H + R_L) - A_D\right)\\
\overline{R}_L^1 &\equiv \frac{1}{2}\left(|z_q|^2 + \eta\,\mathrm{Re}\{z_q^*\overline{z}_q e^{i\phi_q}\}\right) = \frac{1}{2}(R_L + A_D)\\
\overline{R}_H^1 &\equiv \frac{1}{2}\left(|z_q|^2 - \eta\,\mathrm{Re}\{z_q^*\overline{z}_q e^{i\phi_q}\}\right) = \frac{1}{2}(R_H - A_D)\\
\overline{A}_D^1 &\equiv \frac{1}{2}|z_q|^2 = \frac{1}{2}\left(\frac{1}{2}(R_H + R_L) A_D\right).
\end{aligned} \tag{2.5}$$

With regard to the final state itself, a few comments are necessary regarding the

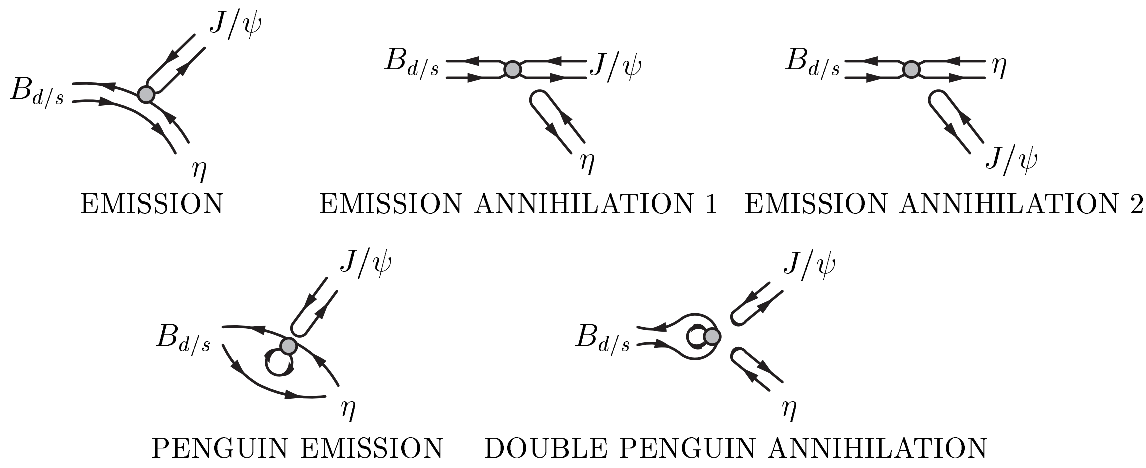

Figure 1: Quark topologies in $B_{d,s} \to J/\psi\,\eta$. The gray blobs denote the contracted short-distance parts.

octet-singlet mixing in the η system. This mixing, parametrized by the mixing angle θ_P, is still a controversial issue (for current experimental values, see e.g. [8]), so rather than using some specific value for θ_P, we rewrite the η wavefunction in the following way:

$$\begin{aligned}\eta &= \frac{1}{\sqrt{6}}(u\bar{u} + d\bar{d} - 2s\bar{s})\cos(\theta_P) + \frac{1}{\sqrt{3}}(u\bar{u} + d\bar{d} + s\bar{s})\sin(\theta_P) \\ &\equiv N_\eta(u\bar{u} + d\bar{d}) + S_\eta(s\bar{s})\,, \qquad (2.6)\end{aligned}$$

$$N_\eta \equiv \frac{\cos(\theta_P)}{\sqrt{6}} + \frac{\sin(\theta_P)}{\sqrt{3}}\,, \qquad S_\eta \equiv \frac{\sin(\theta_P)}{\sqrt{3}} - \frac{2\cos(\theta_P)}{\sqrt{6}} \qquad (2.7)$$

which is the definition we shall use in the following.

2.1 The time-independent amplitudes

The expressions given above for $|A_q(t)|^2$ and $|\overline{A}_q(t)|^2$ depend on the time-independent amplitudes parametrized by eqs. (2.1) and (2.2). In non-leptonic processes, these amplitudes can generally not be evaluated by any method relying on the factorization approach due to final state interaction (FSI) effects. In the present work, however, we do not need an explicit calculation. Rather, we wish to obtain a parametrization of the B_d and B_s amplitudes that will allow us to relate them by SU(3) flavour symmetry. To arrive at such a parametrization, it is sufficient to use the method of quark topologies which does not require the ability to solve the full theory and which allows a systematic classification of long-distance contributions (for a recent update on this method, see [4]). With the topologies listed in figure 1, we obtain the following:[2]

[2]The computational details can be found in an unpublished project. Please contact the author if a copy is needed.

$$A(B_d \to J/\psi\,\eta\,) = N_\eta\,(\lambda^c_{bd} A_d + \lambda^u_{bd} B_d) + \lambda^u_{bd} N_\eta \left[D_d + \frac{\lambda^c_{bd}}{\lambda^u_{bd}}\zeta_d + \xi_d\right],$$

$$A(B_s \to J/\psi\,\eta\,) = S_\eta\,(\lambda^c_{bs} A_s + \lambda^u_{bs} B_s) + \lambda^u_{bs} N_\eta \left[D_s + \frac{\lambda^c_{bs}}{\lambda^u_{bs}}\zeta_s + \xi_s\right] \tag{2.8}$$

with the definitions:

$$\begin{aligned}
A_x &\equiv \langle \mathcal{Q}^{ccx}_{1,2}\rangle_{\mathsf{E}} + \langle \mathcal{Q}^{ccx}_{1,2}\rangle_{\mathsf{PE}} + \langle \mathcal{Q}^{(c)x}_{\text{pen}}\rangle_{\mathsf{E}} + \sum_q \langle \mathcal{Q}^{(q)x}_{\text{pen}}\rangle_{\mathsf{PE}}\,,\\
B_x &\equiv \langle \mathcal{Q}^{uux}_{1,2}\rangle_{\mathsf{PE}} + \langle \mathcal{Q}^{(c)x}_{\text{pen}}\rangle_{\mathsf{E}} + \sum_q \langle \mathcal{Q}^{(q)x}_{\text{pen}}\rangle_{\mathsf{PE}}\,,\\
D_x &\equiv \langle \mathcal{Q}^{uux}_{1,2}\rangle_{\mathsf{EA2}}\,,\\
\zeta_x &\equiv \frac{2N_\eta + S_\eta}{N_\eta}\left(\langle \mathcal{Q}^{ccx}_{1,2}\rangle_{\mathsf{EA1}} + \langle \mathcal{Q}^{ccx}_{1,2}\rangle_{\mathsf{DPA}} + \langle \mathcal{Q}^{(c)x}_{\text{pen}}\rangle_{\mathsf{EA1}} + \sum_q \langle \mathcal{Q}^{(q)x}_{\text{pen}}\rangle_{\mathsf{DPA}}\right) +\\
&\quad + \langle \mathcal{Q}^{(u)x}_{\text{pen}}\rangle_{\mathsf{EA2}} + \langle \mathcal{Q}^{(d)x}_{\text{pen}}\rangle_{\mathsf{EA2}} + \frac{S_\eta}{N_\eta}\langle \mathcal{Q}^{(s)x}_{\text{pen}}\rangle_{\mathsf{EA2}}\,,\\
\xi_x &\equiv \frac{2N_\eta + S_\eta}{N_\eta}\left(\langle \mathcal{Q}^{uux}_{1,2}\rangle_{\mathsf{DPA}} + \langle \mathcal{Q}^{(c)x}_{\text{pen}}\rangle_{\mathsf{EA1}} + \sum_q \langle \mathcal{Q}^{(q)x}_{\text{pen}}\rangle_{\mathsf{DPA}}\right) +\\
&\quad + \langle \mathcal{Q}^{(u)x}_{\text{pen}}\rangle_{\mathsf{EA2}} + \langle \mathcal{Q}^{(d)x}_{\text{pen}}\rangle_{\mathsf{EA2}} + \frac{S_\eta}{N_\eta}\langle \mathcal{Q}^{(s)x}_{\text{pen}}\rangle_{\mathsf{EA2}}\,,
\end{aligned} \tag{2.9}$$

where $\lambda^{q'}_{bq} \equiv V^*_{q'b} V_{q'q}$, $\langle \mathcal{Q}^{uud}_i\rangle_{\mathsf{X}}$ denotes the insertion of the OPE operator $\mathcal{Q}_i$ having external quark lines $buud$ into topology X, and $\langle \mathcal{Q}^{(u)d}_{\text{pen}}\rangle_{\mathsf{X}}$ denotes the combined contribution of the QCD penguin operators $\mathcal{Q}_{3-6}$. The quark content is denoted $(q)d, (q)s$ where q is the flavour of the $q\bar{q}$ pair coming from the gluon and d or s are from the flavour-changing b-quark transition. It should be mentioned that, relative to the current-current operators $\mathcal{Q}_{1,2}$, the electroweak (QCD) penguins contain an extra power of α_{EM} (α_s). In neither the B_d nor the B_s case are the dominant current-current contributions CKM suppressed relative to the penguins, and so we expect $|A_{EWP}|/|A_{CC}| = \mathcal{O}(10^{-2}) = |A_{EWP}|/|A_{QCDP}|$. I have therefore neglected elelectroweak penguin contributions in the above analysis.

Noting that the ζ and ξ terms are SU(3) suppressed and that the D terms are OZI suppressed (see figure 1), we neglect these and obtain:

$$\begin{aligned}
\mathcal{A}_d &= A(B_d \to J/\psi\,\eta\,) = N_\eta\,(V^*_{cb} V_{cd} A_d + V^*_{ub} V_{ud} B_d)\\
&\equiv N_d\left[1 - a_d e^{i\theta_d} e^{i\gamma}\right],\\
\mathcal{A}_s &= A(B_s \to J/\psi\,\eta\,) = S_\eta\,(V^*_{cb} V_{cs} A_s + V^*_{ub} V_{us} B_s)\\
&\equiv N_s\left[1 - a_s \frac{\lambda^2}{1-\lambda^2} e^{i\theta_s} e^{i\gamma}\right]
\end{aligned} \tag{2.10}$$

with

$$N_d \equiv -N_\eta A\lambda^3 A_d\,, \qquad a_d \equiv R_b\left(1 - \frac{\lambda^2}{2}\right)\left|\frac{B_d}{A_d}\right|, \qquad \theta_d \equiv \text{Arg}\left\{\frac{B_d}{A_d}\right\},$$

$$
\begin{aligned}
N_s &\equiv S_\eta A\lambda^2 A_s\left(1-\frac{\lambda^2}{2}\right), \qquad a_s \equiv R_b\left(1-\frac{\lambda^2}{2}\right)\left|\frac{B_s}{A_s}\right|, \\
\theta_s &\equiv -\operatorname{Arg}\left\{\frac{B_s}{A_s}\right\}
\end{aligned}
\tag{2.11}
$$

and the Wolfenstein parameters [3]:

$$
\begin{aligned}
\lambda &\equiv |V_{us}| = 0.22\,, \qquad A \equiv \frac{1}{\lambda^2}|V_{cb}| = 0.81 \pm 0.06\,, \\
R_b &\equiv \frac{1}{\lambda}\left|\frac{V_{ub}}{V_{cb}}\right| = 0.41 \pm 0.07\,.
\end{aligned}
\tag{2.12}
$$

With eqs. (2.10) we have recovered exactly the form of eq. (2.1) by which we parametrized the time-independent amplitudes, but in a form that explicitly separates the CKM structure from the strong amplitudes. This is of essential use below.

3. Extracting γ from $B_{d,s} \to J/\psi\,\eta$

We now adapt the method proposed in [3] for the decay considered here. Defining the time-dependent CP asymmetry by:

$$
a_{CP} = \frac{|A_q(t)|^2 - |\overline{A}_q(t)|^2}{|A_q(t)|^2 + |\overline{A}_q(t)|^2}
\tag{3.1}
$$

and inserting the above expressions, one obtains [3]:

$$
a_{CP}(t) = 2e^{-\Gamma t}\left[\frac{\mathcal{A}_{CP}^{\text{dir}}\cos(\Delta M\; t) + \mathcal{A}_{CP}^{\text{mix}}\sin(\Delta M\; t)}{e^{-\Gamma_H t} + e^{-\Gamma_L t} + \mathcal{A}_{\Delta\Gamma}(e^{-\Gamma_H t} - e^{-\Gamma_L t})}\right]
\tag{3.2}
$$

with

$$
\mathcal{A}_{CP}^{\text{dir}} \equiv \frac{2A_D}{R_H + R_L} = \frac{2\tilde{a}_q \sin\gamma \sin\tilde{\theta}_q}{1 - 2\tilde{a}_q\cos\gamma\cos\tilde{\theta}_q + \tilde{a}_q^2}\,,
\tag{3.3}
$$

$$
\mathcal{A}_{CP}^{\text{mix}} \equiv \frac{2A_M}{R_H + R_L} = \eta\frac{\sin\phi - 2\tilde{a}_q\sin(\gamma+\phi)\cos\tilde{\theta}_q + \tilde{a}_q^2\sin(2\gamma+\phi)}{1 - 2\tilde{a}_q\cos\gamma\cos\tilde{\theta}_q + \tilde{a}_q^2}\,,
\tag{3.4}
$$

$$
\mathcal{A}_{\Delta\Gamma} \equiv \frac{R_H - R_L}{R_H + R_L} = -\eta\frac{\cos\phi_q - 2\tilde{a}_q\cos(\gamma+\phi_q)\cos\tilde{\theta}_q + \tilde{a}_q^2\cos(2\gamma+\phi_q)}{1 - 2\tilde{a}_q\cos\gamma\cos\tilde{\theta}_q + \tilde{a}_q^2}\,,
\tag{3.5}
$$

$$
\tilde{a}_q = \begin{cases} a_d\,; & \text{for } B_d \to J/\psi\,\eta \\ a_s\frac{\lambda^2}{1-\lambda^2}\,; & \text{for } B_s \to J/\psi\,\eta\,, \end{cases}
\tag{3.6}
$$

$$
\tilde{\theta}_q = \begin{cases} \theta_d\,; & \text{for} B_d \to J/\psi\,\eta \\ \theta_s + 180^\circ\,; & \text{for } B_s \to J/\psi\,\eta\,. \end{cases}
\tag{3.7}
$$

As a_s is suppressed by $\frac{\lambda^2}{1-\lambda^2}$, we expect a very small direct CP asymmetry in the B_s process, thus we shall use the asymmetries in $B_d \to J/\psi\,\eta$ combined with R_d and

R_s for extracting γ ($R_q \equiv \frac{1}{2}(R_H^q + R_L^q)$). It should be mentioned that, due to the smallness of $\Delta\Gamma_d$, the "observable" $\mathcal{A}_{\Delta\Gamma}$ will only be measurable for the B_s system. This has no direct consequence on the analysis presented here. Only two of the three asymmetries are independent, and so a measurement of $\mathcal{A}_{CP}^{\rm dir}$ and $\mathcal{A}_{CP}^{\rm mix}$ is sufficient.

$$(\mathcal{A}_{CP}^{\rm dir})^2 + (\mathcal{A}_{CP}^{\rm mix})^2 + (\mathcal{A}_{\Delta\Gamma})^2 = 1 \tag{3.8}$$

which has been checked also to be valid when going to $\varepsilon_{pq} \neq 0$.

The observables $\mathcal{A}_{CP}^{\rm dir}$ and $\mathcal{A}_{CP}^{\rm mix}$ do not depend on the normalization, $|N_d|^2$, and can in principle be obtained by fitting to the CP asymmetry. This yields two equations (3.3)–(3.4) in the three "unknowns": $\tilde{a}_d$, $\tilde{\theta}_d$, and γ (taking the mixing angle, 2β, to be known beforehand). Thus, we need one more observable. Measuring the CP averaged rate yields

$$\langle\Gamma_q\rangle \equiv \Pi_2 \times |N_q|^2 \times R_q\,, \tag{3.9}$$

where Π_2 is the 2-body phase space, and the normalization factors for the B_d and B_s modes are given by eqs. (2.11). A priori, we cannot determine this normalization, but assuming now that the strong interaction dynamics is the same for the B_d and B_s modes (U-spin symmetry), we have $a_d = a_s \equiv a$, $\theta_d = \theta_s \equiv \theta$, and taking kinematics into account:

$$\left|\frac{N_s}{N_d}\right|^2 = \frac{\lambda^2}{1-\lambda^2}\left|\frac{A_s}{A_d}\right|^2 = \frac{\lambda^2}{1-\lambda^2}\frac{\lambda(m_{B_s}^2, m_{J/\psi}^2, m_\eta^2)}{\lambda(m_{B_d}^2, m_{J/\psi}^2, m_\eta^2)} \tag{3.10}$$

with the λ function being the standard one used in kinematics. Thus, forming the ratio $R_s/R_d = (R_H^s + R_L^s)/(R_H^d + R_L^d)$, the strong dynamics cancels out, and we find by combining eq. (3.9) with $R = \frac{1}{2}(R_H + R_L) = \frac{1}{2}(1 - 2\tilde{a}_q \cos\gamma\cos\tilde{\theta}_q + \tilde{a}_q^2)$:

$$\begin{aligned}
\frac{R_s}{R_d} \equiv H &= \frac{1 - 2\frac{\lambda^2}{1-\lambda^2} a_s \cos\gamma\cos\theta_s + (\frac{\lambda^2}{1-\lambda^2})^2 a_s^2}{1 - 2a_d\cos\gamma\cos\theta_d + a_d^2} \\
&= \frac{\langle\Gamma_s\rangle}{\langle\Gamma_d\rangle}\frac{|N_d|^2}{|N_s|^2}\frac{m_{B_s}^3}{m_{B_d}^3}\left(\frac{\lambda(m_{B_d}^2, m_{J/\psi}^2, m_\eta^2)}{\lambda(m_{B_s}^2, m_{J/\psi}^2, m_\eta^2)}\right)^{1/2} \\
&= \frac{\lambda^2}{1-\lambda^2}\frac{\langle\Gamma_s\rangle}{\langle\Gamma_d\rangle}\frac{|N_\eta|^2}{|S_\eta|^2}\frac{m_{B_s}^3}{m_{B_d}^3}\left(\frac{\lambda(m_{B_d}^2, m_{J/\psi}^2, m_\eta^2)}{\lambda(m_{B_s}^2, m_{J/\psi}^2, m_\eta^2)}\right)^{3/2}.
\end{aligned} \tag{3.11}$$

This relation furnishes us with the last observable we need for determining a, θ, and γ as functions of the mixing phase, ϕ. In the Standard Model, this angle is negligibly small for the $B_s - \overline{B}_s$ system whereas it is 2β for the $B_d - \overline{B}_d$ system. Thus, the extraction of γ requires β as an input parameter.

To summarize: the two observables in eqs. (3.3) and (3.4) are to be determined for the decay $B_d \to J/\psi\,\eta$ by measuring the CP asymmetry and fitting to the time-dependent decay amplitudes. This will require tagging and will fix a contour in the

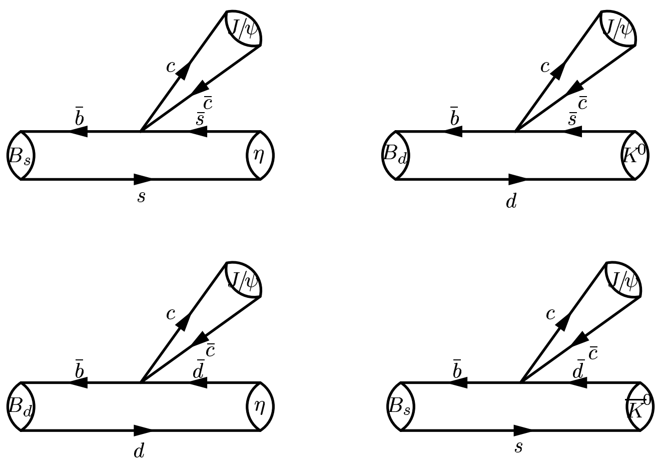

Figure 2: A simplified picture of the four SU(3) related decays $B_{d,s} \to J/\psi\,\eta$ and $B_{d,s} \to J/\psi\, K^0$.

$\gamma - a$ plane. The last observable is provided by $H \equiv R_s/R_d$ which does not require tagging as it depends only on the CP averaged rates. Together with e.g. $\mathcal{A}_{CP}^{\rm mix}$, we fix another contour in the $\gamma - a$ plane. The intersections of these two contours will fix both a and γ, to a theoretical precision depending on SU(3)-breaking corrections and the accuracies on the η and B_d mixing angles, θ_P and 2β.

As we have chosen the same terminology and parametrizations as [3] the contour equations in that paper are directly applicable, and we do not repeat them here.

4. Branchings for $B_{d,s} \to J/\psi\,\eta$

We now present a simplified picture of the $B_{d,s} \to J/\psi\,\eta$ transitions in order to obtain an estimate for the branching ratios. We again take the strong interaction dynamics symmetric under SU(3) transformations and further make the simplifying ansatz of figure 2 that the $\mathcal{Q}_{1,2}$ insertions into emission topologies represent the dominant contributions, resulting in the parameters a_d and a_s of the previous sections being negligibly small (we really only have to make this crude assumption for the $B_d \to J/\psi\eta$ and $B_s \to J/\psi K_S$ processes. In $B_s \to J/\psi\eta$ and $B_d \to J/\psi K_S$ (see [3]), we can justify neglecting these terms because of the λ^2 suppression). With these assumptions, the amplitudes in figure 2 can differ only by CKM factors, kinematics and factors coming from the hadronic wavefunctions. Indeed, eqs. (2.10) become (separating kinematics and dynamics):

$$
\begin{aligned}
A(B_d \to J/\psi\,\eta\,) &= N_\eta \lambda_{bd}^c A_d = N_\eta \lambda_{bd}^c (p_{B_d} + p_\eta)^\mu \epsilon_\mu F(B_d \to J/\psi\,\eta\,)\,,\\
A(B_s \to J/\psi\,\eta\,) &= S_\eta \lambda_{bs}^c A_s = S_\eta \lambda_{bs}^c (p_{B_s} + p_\eta)^\mu \epsilon_\mu F(B_s \to J/\psi\,\eta\,)\,,
\end{aligned}
\tag{4.1}
$$

where p and ϵ are momenta and polarization vectors, respectively, and the form factors, F, parametrize the strong dynamics.

For the $B_{d,s} \to J/\psi\, K_S$ I use the amplitudes in [3] with the same assumptions as above and with a slight modification of the author's notation.

$$\begin{aligned}
A(B_d \to J/\psi\, K_S) &= \lambda^c_{bs} A_d = \lambda^c_{bs}(p_{B_d} + p_{K_S})^\mu \epsilon_\mu F(B_d \to J/\psi\, K_S)\,, \\
A(B_s \to J/\psi\, K_S) &= \lambda^c_{bd} A_s = \lambda^c_{bd}(p_{B_s} + p_{K_S})^\mu \epsilon_\mu F(B_s \to J/\psi\, K_S)\,.
\end{aligned} \tag{4.2}$$

With the assumption that the strong dynamics is SU(3) symmetric, the form factors cancel out when forming ratios, and we arrive at:

$$\begin{aligned}
\frac{A(B_d \to J/\psi\, \eta)}{A(B_d \to J/\psi\, K_S)} &= N_\eta \sqrt{2} \frac{\lambda^c_{bd}}{\lambda^c_{bs}} \frac{(p_{B_d} + p_\eta)^\mu \epsilon_\mu}{(p_{B_d} + p_{K_S})^\mu \epsilon_\mu}\,, \\
\frac{A(B_s \to J/\psi\, \eta)}{A(B_d \to J/\psi\, K_S)} &= S_\eta \sqrt{2} \frac{(p_{B_s} + p_\eta)^\mu \epsilon_\mu}{(p_{B_d} + p_{K_S})^\mu \epsilon_\mu}\,, \\
\frac{A(B_s \to J/\psi\, K_S)}{A(B_d \to J/\psi\, K_S)} &= \frac{\lambda^c_{bd}}{\lambda^c_{bs}} \frac{(p_{B_s} + p_{K_S})^\mu \epsilon_\mu}{(p_{B_d} + p_{K_S})^\mu \epsilon_\mu}\,,
\end{aligned} \tag{4.3}$$

where the $\sqrt{2}$ comes from the translation from K^0 to K_S. In addition to the modes we are interested in, we have written down $A(B_s \to J/\psi\, K_S)$ as a bonus.

By using these relations, it is straightforward to obtain the branching ratios for $B_{d,s} \to J/\psi\eta$ and $B_s \to J/\psi K_S$ in terms of that for $B_d \to J/\psi K_S$ with the measured value $BR(B_d \to J/\psi\, K^0) = 2BR(B_d \to J/\psi\, K_S) = (8.9 \pm 1.2) \times 10^{-4}$ [8]. Inserting this value yields:

$$\begin{aligned}
BR_{B_d \to J/\psi\, \eta} &= 9 \times 10^{-4} |N_\eta|^2 \left|\frac{V_{cd}}{V_{cs}}\right|^2 \left(\frac{\lambda(m^2_{B_d}, m^2_{J/\psi}, m^2_\eta)}{\lambda(m^2_{B_d}, m^2_{J/\psi}, m^2_K)}\right)^{3/2}\,, \\
BR_{B_s \to J/\psi\, \eta} &= 9 \times 10^{-4} |S_\eta|^2 \frac{m^3_{B_d}}{m^3_{B_s}} \left(\frac{\lambda(m^2_{B_s}, m^2_{J/\psi}, m^2_\eta)}{\lambda(m^2_{B_d}, m^2_{J/\psi}, m^2_K)}\right)^{3/2}\,, \\
BR_{B_s \to J/\psi\, K^0} &= 9 \times 10^{-4} \left|\frac{V_{cd}}{V_{cs}}\right|^2 \frac{m^3_{B_d}}{m^3_{B_s}} \left(\frac{\lambda(m^2_{B_s}, m^2_{J/\psi}, m^2_K)}{\lambda(m^2_{B_d}, m^2_{J/\psi}, m^2_K)}\right)^{3/2}\,.
\end{aligned} \tag{4.4}$$

As mentioned in section 2, there is some controversy as to the precise value of the η mixing angle, θ_P, which determines N_η and S_η. Rather than adopting some specific value, we have varied the parameter between $-20°$ $< \theta_P < -10°$, producing the results shown in table 1. The θ_P-dependence between these limits is linear to a good approximation. The uncertainty on these branching ratios is roughly 40%, slightly more for $B_d \to J/\psi\, \eta$ and $B_s \to J/\psi\, K_S$.

	$\theta_P = -10°$	$\theta_P = -20°$
$BR(B_d \to J/\psi\, \eta)$	4.1×10^{-6}	1.6×10^{-6}
$BR(B_s \to J/\psi\, \eta)$	8.3×10^{-4}	9.5×10^{-4}
$BR(B_s \to J/\psi \overline{K}^0)$	5.28×10^{-5}	

Table 1: Estimated branching ratios for $B_{d,s} \to J/\psi\, \eta$ and $B_d \to J/\psi\, \overline{K}^0$.

		$B_d \to J/\psi\,\eta$	$B_s \to J/\psi\,\eta$
HERA-B	$\theta_P = -10°$	35	1.4×10^3
(untagged, /yr)	$\theta_P = -20°$	15	1.6×10^3
CDF II	$\theta_P = -10°$	250	9.5×10^3
(untagged, $2\,\mathrm{fb}^{-1}$)	$\theta_P = -20°$	100	1.1×10^4
ATLAS	$\theta_P = -10°$	1.5×10^4	6.0×10^5
(tagged, $30\,\mathrm{fb}^{-1}$)	$\theta_P = -20°$	5.5×10^3	7.5×10^5
η not reconstructed ($\to$ factor 10% – 20%)			

Table 2: Estimated number of reconstructed $B_{d,s} \to J/\psi\,\eta$ events at HERA-B, CDFII, and ATLAS.

According to the PYTHIA simulation for HERA-B [9], the production rate of B_d mesons is about five times greater than that of B_s mesons, and so we will expect to see about 100 times more B_s decays than B_d decays in the experiment for $\theta_P = -20°$. For $\theta_P = -10°$, we expect about 40 times more B_s decays than B_d decays. This situation is slightly better than for the $B_{d,s} \to J/\psi\,K_S$ decays [3] where we expect to see around 250 B_d events for each B_s event, and so the statistical error on the B_s events is going to have a larger influence on the precision with which γ can be extracted. In both strategies, the asymmetries are to be determined for the mode with *least* statistics, so a few more factors are definitely of use. Given the branching ratios, it has been estimated how many events will be seen by HERA-B, CDF-II, and ATLAS, based on their simulations for $B_d \to J/\psi\,K_S$ [10, 11, 12]. Results are listed in table 2. As the η reconstruction efficiency is not known at present, the numbers presented here are without η reconstruction included. A loose estimate of this efficiency is 10–20%. Taking the η reconstruction efficiency into account, it is certain that HERA-B will not be able to access this mode (the number given is for the machine running at full luminosity), it is an open question whether CDF-II will be able to get it (depending on how much luminosity they get before LHC, and whether they improve their trigger efficiency [11, 13], and it is certain that ATLAS will access it within the first three years of operation.

5. Branchings for $B_{d,s} \to \eta\,\ell^+\ell^-$

The branchings for $B_s \to \eta\,\ell^+\ell^-$ can be obtained by using that the process is related by the approximate SU(3) flavour symmetry to $B_d \to K\ell^+\ell^-$ whose spectrum has been calculated in [6]. Due to the close similarities between η and K mesons (pseudoscalars, similar masses), and going to the SU(3) symmetric limit, we merely need to replace $B \to K$ by $B \to \eta$ form factors and to take CKM factors into account. At present, no reliable form factor calculations for $B_{d,s} \to \eta$ exist.

In the following, we estimate the form factors for $B_s \to \eta$ by SU(3) relation to the $B \to K$ form factors presented in [6], effectively resulting in a multiplication of the form factors by S_η (see eq. (2.6)). The calculation of $B_d \to \eta$ form factors is essentially identical, and so we omit an explicit calculation.

$$\begin{aligned} \langle\, \eta(p_\eta)\, |\bar{s}\gamma_\mu b|\, B_s(p_B)\,\rangle &= S_\eta \Big[f_+^K(s)(p_B+p_\eta)_\mu + f_-^K(s)(p_B-p_\eta)_\mu \Big]\,, \\ \langle\, \eta(p_\eta)\, |\bar{s}\sigma_{\mu\nu}q^\nu b|\, B_s(p_B)\,\rangle &= S_\eta i \Big[s(p_B+p_\eta)_\mu - (m_B^2-m_\eta^2)q_\mu \Big] \frac{f_T^K(s)}{m_B+m_\eta}\,. \end{aligned} \tag{5.1}$$

With these form factors, the total branching ratios for $B_{d,s} \to \eta\,\ell^+\ell^-$ as a function of θ_P has been calculated using eq. (6.1) integrated over the kinematical region without intermediate ψ resonant states included in $C_9^{\rm eff}$.

	$\theta_P = -10°$	$\theta_P = -20°$
$B_s \to \eta\,\mu^+\mu^-$	4.6×10^{-7}	5.2×10^{-7}
$B_d \to \eta\,\mu^+\mu^-$	2.8×10^{-9}	1.1×10^{-9}

Table 3: Branching ratios for $B_{d,s} \to \eta\,\ell^+\ell^-$.

A conservative theoretical uncertainty on this calculation is given by $\approx 40\%$ from SU(3) breaking corrections, and $\approx 25\%$ from uncertainties on the initial $B \to K$ form factors, yielding a total uncertainty around 50%. Results are listed in table 3 for the case $\ell = \mu$. Based on simulations from the two experiments of $B \to K^*\ell^+\ell^-$, the number of reconstructed $B_s \to \eta\,\mu^+\mu^-$ and $B_d \to \eta\,\mu^+\mu^-$ events have been estimated, again without taking the η reconstruction efficiency into account (table 4).

6. Constraining C_{10A} in semileptonic B decays

The differential branching ratios of $B_{d,s} \to K\ell^+\ell^-$ and $B_{s,d} \to \eta\ell^+\ell^-$ can be written as [6]:

$$\begin{aligned} \frac{d\Gamma}{d\hat{s}} = \frac{G_F^2\alpha^2 m_B^5}{2^{10}\pi^5}|V_{ts,td}^* V_{tb}|^2 \hat{u}(\hat{s}) \Bigg[& \left(|A'|^2+|C'|^2\right)\left(\lambda - \frac{\hat{u}(\hat{s})^2}{3}\right) + \\ & + |C'|^2 4\hat{m}_\ell^2 \left(2+2\hat{m}_K^2-\hat{s}\right) + \\ & + \mathrm{Re}(C'D'^*)8\hat{m}_\ell^2(1-\hat{m}_K^2)|D'|^2 4\hat{m}_\ell^2\hat{s} \Bigg] \end{aligned} \tag{6.1}$$

		$B_d \to \eta\,\mu^+\mu^-$	$B_s \to \eta\,\mu^+\mu^-$
CDF-II	$\theta_P = -10°$	< 5	90
($2\,\mathrm{fb}^{-1}$)	$\theta_P = -20°$	< 2	110
ATLAS	$\theta_P = -10°$	< 10	250
($30\,\mathrm{fb}^{-1}$)	$\theta_P = -20°$	< 5	275
η not reconstructed ($\to$ factor 10% – 20%)			

Table 4: Estimated number of reconstructed $B_{d,s} \to \eta\,\ell^+\ell^-$ events at CDF-II and ATLAS.

with

$$\lambda_{K,\eta} \equiv \lambda\left(1, \hat{m}_{K,\eta}^2, \hat{s}\right), \qquad \hat{u}(\hat{s}) \equiv \sqrt{\lambda\left(1 - 4\frac{\hat{m}_\ell^2}{\hat{s}}\right)} \tag{6.2}$$

and the coefficients A', C', D' given in terms of Wilson coefficients and form factors as [6]:

$$\begin{aligned} A'(s) &= C_9^{\rm eff}(s) f_+^{K,\eta}(s) + \frac{2m_b}{m_B + m_{K,\eta}} C_7^{\rm eff} f_T^{K,\eta}(s)\,, \\ C'(s) &= C_{10A} f_+^{K,\eta}(s)\,, \qquad D'(s) = C_{10A} f_-^{K,\eta}(s)\,. \end{aligned} \tag{6.3}$$

Variables with ˆ have been normalized by m_B^2, λ is the standard kinematical function, and s is the energy in the CM of the lepton pair. The C_i entering these expressions are the Wilson coefficients of the OPE operators contributing to semileptonic B decays:

$$\begin{aligned} \mathcal{O}_{7\gamma} &= \frac{e}{8\pi^2} m_b \bar{s}_\alpha \sigma^{\mu\nu}(1+\gamma_5) b^\alpha F_{\mu\nu}\,, \qquad \mathcal{O}_{9V}^\ell = (\bar{s}_\alpha b^\alpha)_{V-A}(\bar{\ell}\ell)_V\,, \\ \mathcal{O}_{10A}^\ell &= (\bar{s}_\alpha b^\alpha)_{V-A}(\bar{\ell}\ell)_A \end{aligned} \tag{6.4}$$

$C_7^{\rm eff} = C_7 + C_5/3 - C_6$ and $C_9^{\rm eff}$ contains corrections due to intermediate quark loops in the decay (see e.g. [6]).

For zero lepton mass ($\ell = e, \mu$) only the $(|A'|^2 + |C'|^2)$ term in (6.1) remains. Measuring $\mathrm{d}\Gamma/d\hat{s}$ for e.g. the $B \to K\mu^+\mu^-$ thus yields a way of experimentally determining this sum. This is attractive since A' is excessively difficult to calculate theoretically, due to large hadronic uncertainties caused by the existence of intermediate ψ resonances entering $C_9^{\rm eff}$.

In contrast to this, the terms in (6.1) containing C' and D' depend only on C_{10A} and form factors. Eliminating A' from the τ distribution by isolating the term $(|A'|^2 + |C'|^2)$ in the muon distribution and inserting it in eq. (6.1) for $\ell = \tau$ yields the following relation,[3] where we have integrated over the kinematical region of the τ spectrum:

$$\begin{aligned} \int_{4\hat{m}_\tau^2}^{(\hat{m}_B - \hat{m}_K)^2} \frac{dB_\tau}{d\hat{s}} d\hat{s} &= \int_{4\hat{m}_\tau^2}^{(\hat{m}_B - \hat{m}_K)^2} \sqrt{1 - \frac{4\hat{m}_\tau^2}{\hat{s}}} \frac{dB_\mu}{d\hat{s}} \left(1 + \frac{2\hat{m}_\tau^2}{\hat{s}}\right) d\hat{s} + \\ &\quad + |C_{10A}|^2 \frac{G_F^2 \alpha^2 m_B^5}{2^{10}\pi^5} |V_{ts}^* V_{tb}|^2 \tau_B \int_{4\hat{m}_\tau^2}^{(\hat{m}_B - \hat{m}_K)^2} F(\hat{s}) d\hat{s}\,, \end{aligned} \tag{6.5}$$

where τ_B is the B meson lifetime and

$$F(\hat{s}) = \sqrt{\lambda\left(1 - \frac{4\hat{m}_\tau^2}{\hat{s}}\right)} \left(f_+^2 4\hat{m}_\tau^2 \left(2 + 2\hat{m}_K^2 - \hat{s}\right) + f_+ f_- 8\hat{m}_\tau^2 \left(1 - \hat{m}_K^2\right) + f_-^2 4\hat{m}_\tau^2 \hat{s}\right).$$

[3]The result for $B_s \to \eta\,\mu^+\mu^-$ is completely analogous.

To illustrate the order of magnitude of the terms entering eq. (6.5), we give their SM values below (with the form factors presented in [6]).

$$\begin{aligned} \int_{4\hat{m}_\tau^2}^{(\hat{m}_B-\hat{m}_K)^2} \frac{dB_\tau}{d\hat{s}} d\hat{s} &= 1.35\times 10^{-7}\,, \\ \int_{4\hat{m}_\tau^2}^{(\hat{m}_B-\hat{m}_K)^2} d\hat{s}\sqrt{1-\frac{4\hat{m}_\tau^2}{\hat{s}}}\,\frac{dB_\mu}{d\hat{s}}\left(1+\frac{2\hat{m}_\tau^2}{\hat{s}}\right) &= 0.90\times 10^{-7}\,, \\ |C_{10A}|^2\frac{G_F^2\alpha^2\hat{m}_B^5}{2^{10}\pi^5}|V_{ts}^*V_{tb}|^2\tau_B\int_{4\hat{m}_\tau^2}^{(\hat{m}_B-\hat{m}_K)^2} d\hat{s}F(\hat{s}) &= 0.45\times 10^{-7}\,. \end{aligned} \tag{6.6}$$

Due to the close similarity between K and η, these results are not changed substantially when considering $B_s \to \eta\,\mu^+\mu^-$. Using very precise form factor calculations, $F(s)$ can be evaluated theoretically to within perhaps $\pm 20\%$ uncertainty. At second generation machines, it is reasonable to expect measurements of $BR(B\to K\tau^+\tau^-)$ and the muon spectrum to about 10% accuracy. Isolating C_{10A} in the above formula would then yield an overall uncertainty of approximately $\pm 20\%$. In the SUGRA models investigated in [14] (minimal and non minimal, $\tan\beta = 2$ and $\tan\beta = 30$), C_{10A} lies within $\pm 10\%$ of the SM value, and so a distinction is not yet possible for these cases. In the 2HDM,[4] C_{10A} receives contributions from charged Higgs bosons, and we have used [15] for values of the Yukawa coupling of the charged Higgs to the top quark, $0 < \lambda_{tt} < 0.3$ and charged Higgs masses $0.2\,\mathrm{TeV} < M_{H^\pm} < 1\,\mathrm{TeV}$. The case of maximal deviation ($M_{H^\pm} = 1\,\mathrm{TeV}$, $\lambda_{tt} = 0.3$) produces only about 5% deviation. Thus we see that some of the most common models cannot be ruled out by a measurement such as the one proposed here in the near future. Both form factor calculations and experimental precision must be improved before this method becomes viable. Again, this result is not substantially altered by having a B_S initial state and an η in the final state.

7. Conclusion

Using the method of quark topologies and invoking SU(3)-suppression of certain topologies as well as SU(3) symmetry of the strong interaction dynamics, it is possible to obtain a parametrization of the $B_{d,s}\to J/\psi\,\eta$ amplitude which allows the extraction of the angle γ of the Unitarity Triangle through measurement of CP asymmetries in the $B_d \to J/\psi\,\eta$ process and measurement of the CP averaged widths of the $B_{d,s}\to J/\psi\eta$ processes together. Electroweak penguins cannot lead to any problems in this extraction. The measurement of asymmetries require tagging whereas the CP averaged rates do not.

[4]Two Higgs Doublet Model

Additionally, an estimate of the branching ratios for $B_d \to J/\psi\,\eta$ has been presented, yielding an expected branching ratio of around 2×10^{-6} for $B_d \to J/\psi\,\eta$ and around 9×10^{-4} for $B_s \to J/\psi\,\eta$ (taking $\theta_P \approx -20°$). Due to the difference in production ratios of B_d and B_s, this means a factor 40-100 more B_s decays than B_d decays in the experiment, depending on the value of θ_P. This is an improved situation relative to the $B_{d,s} \to J/\psi\, K_S$ decays, where a factor of around 250 difference is expected. It has been estimated that HERA-B will probably not be able to see $B_d \to J/\psi\,\eta$ decays whereas the Tevatron experiments CDF-II/D0 perhaps stand a chance of extracting of γ from $B_{d,s} \to J/\psi\eta$. Finally, the ATLAS experiment should have sufficient statistics to allow the extraction of γ even taking the most pessimistic approach to the uncertainties in the estimates of branching ratios and production rates presented in this paper.

As an addendum, a strategy to constrain C_{10A} in semileptonic B decays with small hadronic uncertainties has been proposed. It has been found to require both high theoretical and experimental precision, of a kind that will not be available in the near future.

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Erratum

To comply with the PDG convention for the sign of θ_P [8], equations (2.6) and (2.7) should read:

$$\eta = \frac{1}{\sqrt{6}}(u\bar{u} + d\bar{d} - 2s\bar{s})\cos(\theta_P) - \frac{1}{\sqrt{3}}(u\bar{u} + d\bar{d} + s\bar{s})\sin(\theta_P)$$
$$\equiv N_\eta(u\bar{u} + d\bar{d}) + S_\eta(s\bar{s}) \tag{2.6}$$
$$N_\eta \equiv \frac{\cos(\theta_P)}{\sqrt{6}} - \frac{\sin(\theta_P)}{\sqrt{3}} \qquad S_\eta \equiv \frac{-\sin(\theta_P)}{\sqrt{3}} - \frac{2\cos(\theta_P)}{\sqrt{6}} \tag{2.7}$$

While this change of sign does not affect the algebra in the article, it does produce changes whereever numerical values of N_η and S_η are inserted. Hence, the branching fractions and expected event numbers for $B_{d,s} \to J/\psi\,\eta$ (tables 1 and 2) and $B_{d,s} \to \eta\ell^+\ell^-$ (tables 3 and 4) are replaced as follows:

	$\theta_P = -10^\circ$	$\theta_P = -20^\circ$
$BR(B_d \to J/\psi\,\eta)$	1.1×10^{-5}	1.5×10^{-5}
$BR(B_s \to J/\psi\,\eta)$	5.0×10^{-4}	3.3×10^{-4}

Table 1: Estimated branching ratios for $B_{d,s} \to J/\psi\,\eta$ and $B_d \to J/\psi\,\overline{K}^0$.

		$B_d \to J/\psi\,\eta$	$B_s \to J/\psi\,\eta$
HERA-B	$\theta_P = -10^\circ$	100	850
(untagged, /yr)	$\theta_P = -20^\circ$	125	560
CDF II	$\theta_P = -10^\circ$	700	5.8×10^3
(untagged, 2fb^{-1})	$\theta_P = -20^\circ$	1000	3.8×10^3
ATLAS	$\theta_P = -10^\circ$	4×10^4	3.5×10^5
(tagged, 30fb^{-1})	$\theta_P = -20^\circ$	5×10^4	2.5×10^5
η not reconstructed ($\to$ factor 10% – 20%)			

Table 2: Estimated number of reconstructed $B_{d,s} \to J/\psi\,\eta$ events at HERA-B, CDFII, and ATLAS.

	$\theta_P = -10^\circ$	$\theta_P = -20^\circ$
$B_s \to \eta\,\mu^+\mu^-$	3.0×10^{-7}	2.0×10^{-7}
$B_d \to \eta\,\mu^+\mu^-$	7.7×10^{-9}	1.0×10^{-8}

Table 3: Branching ratios for $B_{d,s} \to \eta\,\ell^+\ell^-$.

		$B_d \to \eta\, \mu^+ \mu^-$	$B_s \to \eta\, \mu^+ \mu^-$
CDF-II	$\theta_P = -10^\circ$	8	60
(2 fb^{-1})	$\theta_P = -20^\circ$	10	40
ATLAS	$\theta_P = -10^\circ$	20	160
(30 fb^{-1})	$\theta_P = -20^\circ$	25	105
η not reconstructed ($\to$ factor 10% – 20%)			

Table 4: Estimated number of reconstructed $B_{d,s} \to \eta\, \ell^+ \ell^-$ events at CDF-II and ATLAS.