Extracting $\beta$ and the new $D_{sJ}$ resonances\footnote{talk given at MRST 2004: From Quarks to Cosmology, Concordia University, Montreal, May 2004.}

Alakabha Datta\footnote{datta@physics.utoronto.ca}

Department of Physics,  
University of Toronto,  
60 St George St,  
Toronto, Ontario, Canada

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Abstract

The three body decays $B \to D^{(*)} \bar{D}^{(*)} K_s$ may be used to measure both $\sin 2\beta$ and $\cos 2\beta$. Crucial to the $\cos 2\beta$ measurement is the resonant contribution to the three body decay from p-wave excited $D_s$ states. If these p-wave states are the newly discovered $D_s(2317)$ and $D_s(2460)$ then they are below the $D^{(*)}K$ threshold and hence do not contribute to $B \to D^{(*)} \bar{D}^{(*)} K_s$. The three body decays can then be used to measure $\sin 2\beta$ without resonant dilution and to look for new physics in $b \to c \bar{c} s$ transition.
1 Introduction

The decay $B^0 \rightarrow J/\psi K_s$ provides a clean measurement of the angle $\sin(2\beta)$ in the unitarity triangle. Both BaBar and Belle have measured this CP phase, with the world average being

$$\sin 2\beta = 0.736 \pm 0.049.$$  \hspace{1cm} (1)

Other modes can also provide relevant information on the angle $\beta$, an example being the decay $B^0 \rightarrow D^{(*)}\overline{D}^{(*)}$. The possibility of extracting $\cos 2\beta$ from the decay $B^0 \rightarrow D^{(*)}\overline{D}^{(*)}$ was mentioned in Ref. [3]. These modes are enhanced relative to $B^0 \rightarrow D^{(*)}\overline{D}^{(*)}$ by the factor $|V_{cs}/V_{cd}|^2 \sim 20$. As in the case of $B^0 \rightarrow J/\psi K_s$ decay, the penguin contamination is expected to be small in these decays. Moreover these decays can be used to probe both $\sin 2\beta$ and $\cos 2\beta$ which can resolve $\beta \rightarrow \pi/2 - \beta$ ambiguity [4].

2 $\beta$ from $B \rightarrow D^{(*)}\overline{D}^{(*)}K_s$

The amplitude for the decay $B^0 \rightarrow D^{*}\overline{D}^{*}K_s$ can have a resonant contribution and a non-resonant contribution. For the resonant contribution the $D^{*}K_s$ in the final state comes dominantly from an excited $D_s(1^+)$ state. In the approximation of treating $D^{*}\overline{D}^{*}K_s$ as $D^{*}D_s(\text{excited})$, there are four possible excited p-wave $D_s$ states which might contribute. These are the two states with the light degrees of freedom in a $j^P = 3/2^+$ state and the two states with light degrees of freedom in a $j^P = 1/2^+$ state. Since the states with $j^P = 3/2^+$ decay via d-wave to $D^{*}K_s$, they are suppressed. Of the states with light degrees of freedom in $j^P = 1/2^+$ states, only the $1^+$ state contributes. The $0^+$ state is forbidden to decay to the final state $D^{*}K_s$.

To estimate the above contribution and to calculate the non-resonant amplitude, we use heavy hadron chiral perturbation theory (HHCHPT)[5]. The momentum $p_k$ of $K_s$ can have a maximum value of about 1 GeV for $B^0 \rightarrow D^{*}\overline{D}^{*}K_s$. This is of the same order as $\Lambda_\chi$ which sets the scale below which we expect HHCHPT to be valid. It follows that in the present case it is reasonable to apply HHCHPT to calculate the three body decays.

In the lowest order in the HHCHPT expansion, contributions to the decay amplitude come from the contact interaction terms and the pole diagrams which give rise to the non-resonant and resonant contributions respectively. The pole diagrams get contributions from the various multiplets involving $D_s$ type resonances as mentioned above. In the framework of HHCHPT, the ground state heavy meson has the light degrees of freedom in a spin-parity state $j^P = \frac{1}{2}^-$, corresponding to the usual pseudoscalar-vector meson doublet with $J^P = (0^-,1^-)$. The first excited state involves a p-wave excitation, in which the light degrees of freedom have $j^P = \frac{1}{2}^+$ or $\frac{3}{2}^+$. In the latter case we have a heavy doublet with $J^P = (1^+,2^+)$. These states can probably be identified with $D_{s1}(2536)$ and $D_{sJ}(2573)$. Heavy quark symmetry

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rules out any pseudoscalar coupling of this doublet to the ground state at lowest order in the chiral expansion \[ \mathcal{O} \]; hence the effects of these states will be suppressed and we will ignore them in our analysis.

The other excited doublet has \( J^P = (0^+, 1^+) \). These states are expected to decay rapidly through s-wave pion emission and have large widths \[ \mathcal{O} \]. The \( 1^+ \) state in the \( D \) system has already been seen. Only the \( 1^+ \) can contribute in this case. For later reference, we denote this state by \( D^\pi_{s1} \). However, quark model estimates suggest \[ \mathcal{O} \] that these states should have masses near \( m + \delta m \) with \( \delta m = 500 \text{ MeV} \), where \( m \) is the mass of the lowest multiplet.

Assuming that the leading order terms in HHCHPT give the dominant contribution to the decay amplitude we will neglect all sub-leading effects suppressed by \( 1/\Lambda \) and \( 1/m \), where \( m \) is the heavy quark mass. It can be shown that from the time dependent analysis of \( B^0(t) \rightarrow D^{**}D^{*-}K_s \) one can extract \( \sin(2\beta) \) and \( \cos(2\beta) \). Measurement of both \( \sin(2\beta) \) and \( \cos(2\beta) \) can resolve the \( \beta \rightarrow \pi/2 - \beta \) ambiguity as already mentioned. The measurement of \( \sin(2\beta) \) can be made from the time dependent partial rate asymmetry while a fit to the time dependent rate for \( \Gamma[B^0(t) \rightarrow D^{**}D^{*-}K_s] + \Gamma[B^0(t) \rightarrow D^{*+}D^{*-}K_s] \) may be used for the extraction of \( \cos(2\beta) \). The \( \cos(2\beta) \) term measures the overlap of the imaginary part of the amplitudes for \( B \rightarrow D^{**}D^{*-}K_s \) and \( B \rightarrow D^{*+}D^{*-}K_s \) decays and is non zero only if there is a resonance contribution.

As in the case for \( B \rightarrow D^{**}D^{*-} \) the asymmetry in \( B \rightarrow D^{**}D^{*-}K_s \) is also diluted. For the non resonant contribution to \( B \rightarrow D^{**}D^{*-}K_s \) the final state is an admixture of CP states with different CP parities. This leads to the dilution of the asymmetry and this is the same dilution of the asymmetry as in the case for \( B \rightarrow D^{**}D^{*-} \). When the resonant contribution is included there is further dilution of the asymmetry from the additional mismatch of the amplitudes for \( B \) and \( B^- \) decays. One can reduce the additional dilution of the CP asymmetry by imposing cuts to remove the resonance. A narrow resonance is preferable as it can be more effectively removed from the signal region than a broad resonance. When we include the resonance contribution it turns out that a broader resonance leads to a larger value of \( D \) and is a more useful probe of \( \cos(2\beta) \) because of the the larger overlap of the amplitudes for \( B \rightarrow D^{**}D^{*-}K_s \) and \( B^- \rightarrow D^{**}D^{*-}K_s \) decays.

The differential decay distribution of the time independent process \( B^0 \to D^{**}D^{*-}K_s \) can be used to discover the \( 1^+ \) resonance \( D^{**}_{s1} \) if it is above the \( D^{(*)}K \) threshold. The differential decay distribution for small values of \( E_k \), the kaon energy, shows a clear resonant structure which comes from the pole contribution to the amplitude with the excited \( J^P = 1^+ \) intermediate state. Therefore, examination of the \( D^{*}K_s \) mass spectrum may be the best experimental way to find the broad \( 1^+ \) p-wave \( D_s \) meson and a fit to the decay distribution will measure its mass and the coupling.

The extraction of \( \sin 2\beta \) and \( \cos 2\beta \) from the time dependent rate for \( B(t) \rightarrow D^{**}D^{*-}K_s \) can be done in the following manner: We define the following amplitudes

\[
a^{\lambda_1,\lambda_2} = A(B^0(p) \rightarrow D^{**}_{\lambda_1}(p_+)D^{*-}_{\lambda_2}(p_-)K_s(p_k)) \quad (2)
\]
\[\bar{a}^{\lambda_1,\lambda_2} \equiv A(B^0(p) \to D_{\lambda_1}^+(p_+)D_{\lambda_2}^- (p_-)K_s(p_k)), \tag{3}\]

where \(B^0\) and \(\bar{B}^0\) represent unmixed neutral \(B\) and \(\lambda_1\) and \(\lambda_2\) are the polarization indices of the \(D^+\) and \(D^-\) respectively.

The time-dependent amplitudes for an oscillating state \(B^0(t)\) which has been tagged as a \(B^0\) meson at time \(t = 0\) is given by,

\[A^{\lambda_1,\lambda_2}(t) = a^{\lambda_1,\lambda_2} \cos \left(\frac{\Delta m t}{2}\right) + ie^{-2i\beta} \bar{a}^{\lambda_1,\lambda_2} \sin \left(\frac{\Delta m t}{2}\right), \tag{4}\]

and the time-dependent amplitude squared summed over polarizations and integrated over the phase space angles is:

\[|A(s^+, s^-; t)|^2 = \frac{1}{2} \left[G_0(s^+, s^-) + G_c(s^+, s^-) \cos \Delta m t - G_s(s^+, s^-) \sin \Delta m t\right] \]

with

\[G_0(s^+, s^-) = |a(s^+, s^-)|^2 + |ar{a}(s^+, s^-)|^2, \]
\[G_c(s^+, s^-) = |a(s^+, s^-)|^2 - |ar{a}(s^+, s^-)|^2, \]
\[G_s(s^+, s^-) = 2\Im(e^{-2i\beta} \bar{a}(s^+, s^-)a^*(s^+, s^-)) = -2\sin(2\beta) \Re(\bar{a}a^*) + 2\cos(2\beta) \Im(\bar{a}a^*).\]

The variables \(s^+\) and \(s^-\) are the Dalitz plot variable

\[s^+ = (p_+ + p_k)^2, \quad s^- = (p_- + p_k)^2\]

The transformation defining the CP-conjugate channel \(\bar{B}^0(t) \to D^{*-}D^{**}K_s\) is \(s^+ \leftrightarrow s^-, \ a \leftrightarrow \bar{a}\) and \(\beta \to -\beta\). Then:

\[|\bar{A}(s^-, s^+; t)|^2 = \frac{1}{2} \left[G_0(s^-, s^+) - G_c(s^-, s^+) \cos \Delta m t + G_s(s^-, s^+) \sin \Delta m t\right]. \]

Note that for simplicity the \(e^{-i\gamma}\) and constant phase space factors have been omitted in the above equations.

It is convenient in our case to replace the variables \(s^+\) and \(s^-\) by the variables \(y\) and \(E_k\) where \(E_k\) is the \(K_s\) energy in the rest frame of the \(B\) and \(y = \cos \theta\) with \(\theta\) being the angle between the momentum of \(K_s\) and \(D^{**}\) in a frame where the two \(D^*\) are moving back to back along the \(z\)-axis. This frame is boosted with respect to the rest frame of the \(B\) with \(\beta = -(\vec{p}_k/m_B)\left(1/(1 - E_k/m_B)\right)\). Note \(s^+ \leftrightarrow s^-\) corresponds to \(y \leftrightarrow -y\). The variable \(y\) can be expressed in terms of variables in the rest frame of \(B\). For instance

\[E_+ = \frac{E'_B E'_p - p'_B p'_y}{m_B} \tag{5}\]
where $E_+$ and $E'_+$ are the energy of the $D^{*+}$ in the rest frame of the $B$ and in the boosted frame while $E'_B$ is the energy of the $B$ in the boosted frame. The magnitudes of the momentum of the $B$ and the $D^{*+}$ in the boosted frame are given by $p'_B$ and $p'_+$ respectively.

In the approximation of neglecting the penguin contributions, proportional to the small CKM elements $V_{ub}V_{us}^*$, to the amplitude there is no direct CP violation. This leads to the relation

$$a^{\lambda_1,\lambda_2}(\vec{p}_{k_1}, E_k) = \bar{a}^{-\lambda_1,-\lambda_2}(-\vec{p}_{k_1}, E_k)$$

(6)

where $\vec{p}_{k_1}$ is the momentum of the of the $K_s$ in the boosted frame. The above relations then leads to

$$G_0(-y, E_k) = G_0(y, E_k)$$

(7)

$$G_c(-y, E_k) = -G_c(y, E_k)$$

(8)

$$G_{s1}(-y, E_k) = G_{s1}(y, E_k)$$

(9)

$$G_{s2}(-y, E_k) = -G_{s2}(y, E_k)$$

(10)

where we have defined

$$G_{s1}(y, E_k) = \Re(\bar{a}a^*)$$

(11)

$$G_{s2}(-y, E_k) = \Im(\bar{a}a^*)$$

(12)

Carrying out the integration over the phase space variables $y$ and $E_k$ one gets the following expressions for the time-dependent total rates for $B^0(t) \to D^{*+} D^{*-} K_s$ and the CP conjugate process

$$\Gamma(t) = \frac{1}{2}[I_0 + 2\sin(2\beta)\sin(\Delta m t)I_{s1}]$$

(13)

$$\Gamma(t) = \frac{1}{2}[I_0 - 2\sin(2\beta)\sin(\Delta m t)I_{s1}]$$

(14)

where $I_0$ and $I_{s1}$ are the integrated $G_0(y, E_k)$ and $G_{s1}(y, E_k)$ functions. One can then extract $\sin(2\beta)$ from the rate asymmetry

$$\frac{\Gamma(t) - \Gamma(t)}{\Gamma(t) + \Gamma(t)} = D \sin(2\beta) \sin(\Delta m t)$$

(15)

where

$$D = \frac{2I_{s1}}{I_0}$$

(16)

is the dilution factor. The quantities $I_{s1,0}$ can be calculated in HHCHPT[4].
The cos(2\beta) term can be probed by integrating over half the range of the variable \( y \) which can be taken for instance to be \( y \geq 0 \). In this case we have

\[
\Gamma(t) = \frac{1}{2} [J_0 + J_c \cos(\Delta mt) + 2 \sin(2\beta) \sin(\Delta mt) J_{s1} - 2 \cos(2\beta) \sin(\Delta mt) J_{s2}]
\]

\[
\bar{\Gamma}(t) = \frac{1}{2} [J_0 + J_c \cos(\Delta mt) - 2 \sin(2\beta) \sin(\Delta mt) J_{s1} - 2 \cos(2\beta) \sin(\Delta mt) J_{s2}]
\]

where \( J_0, J_c, J_{s1} \) and \( J_{s2} \), are the integrated \( G_0(y, E_k) \), \( G_c(y, E_k) \), \( G_{s1}(y, E_k) \) and \( G_{s2}(y, E_k) \) functions integrated over the range \( y \geq 0 \). The quantities \( J_{0,c,s1,s2} \) can be calculated in HHCHPT\(^4\). One can measure cos(2\beta) by fitting to the time distribution of \( \Gamma(t) + \bar{\Gamma}(t) \). Measurement of the cos(2\beta) then resolves the \( \beta \rightarrow \frac{\pi}{2} - \beta \) ambiguity. In passing we note that only the sign of cos(2\beta) is required to resolve the \( \beta \rightarrow \frac{\pi}{2} - \beta \) ambiguity.

3 The new \( D_{sJ} \) resonances

The previous year saw the discovery of an unexpectedly light narrow resonance in \( D_s^+ \pi^0 \) with a mass of 2317\,MeV/c\(^2\) first reported by the BaBar collaboration\(^9\), together with another second narrow resonance in \( D_s \pi^0 \gamma \) with a mass 2460\,MeV/c\(^2\)\(^1\). The smaller than expected masses and narrow widths of these states have led, among other explanations\(^1\), to a multi-quark anti-quark or a \( DK \) molecule interpretation of these states\(^1\), or to an interpretation as p-wave states where the light degrees of freedom are in an angular momentum state \( j_q = \frac{1}{2} \)\(^1\), or even some combination of these\(^1\). There are also conflicting lattice interpretations of these states\(^1\). The mass difference between the \( D_s(2317) \) and the well established lightest charm-strange meson, \( D_s \), is \( \Delta M = 350 \, \text{MeV/c}^2 \). This is less than the kaon mass, thus kinematically forbidding the decay \( D_s(2317) \rightarrow D_{u,d} + K \). The possible resonance at 2460\,MeV/c\(^2\) also has such a mass difference when taken with the lighter \( D^* \) state. The interpretation of these states as bound \( D^{(*)} K \) molecules just below the \( D^{(*)} K \) threshold is particularly interesting in the light of the recent discovery of a narrow resonance in the decay \( J/\psi \rightarrow \gamma pp \)\(^1\) which has been interpreted as a zero baryon number, “deuteron-like singlet \(^1S_0\)” bound state of \( p \) and \( \bar{p} \)\(^1\).

The production of these new \( D_{sJ} \) states in non leptonic \( B \) decays appear to be in conflict with theory predictions based on factorization and heavy quark symmetry\(^1\).\(^1\). Following Ref.\(^1\) let us first assume that we can identify the the newly discovered states \( D_s(2317) \) with \( D_{s0} \) and \( D_s(2460) \) with \( D_{s1}^* \). In the Standard Model (SM) the amplitudes for \( B \rightarrow D^{(*)} D_{s0}(D_{s1}^*) \), are generated by the following effective Hamiltonian\(^2\):

\[
H_{eff}^q = \frac{G_F}{\sqrt{2}} V_{fb} V_{fq}^* (c_1 O_{1f}^q + c_2 O_{2f}^q)
\]
\[ - \sum_{i=3}^{10} (V_{ub} V_{ub}^* c_i^u + V_{cb} V_{cb}^* c_i^c + V_{tb} V_{tb}^* c_i^t) O_i^q + H.C., \]  

(17)

where the superscript \( u, c, t \) indicates the internal quark, \( f \) can be \( u \) or \( c \) quark, \( q \) can be either a \( d \) or a \( s \) quark depending on whether the decay is a \( \Delta S = 0 \) or \( \Delta S = -1 \) process. The operators \( O_i^q \) are defined as [21]

\[ \begin{align*}
O_{1f}^q &= \bar{q}_a \gamma_\mu L_f \bar{f}_\beta \gamma^\mu L_b \alpha, \\
O_{2f}^q &= \bar{q}_a \gamma_\mu L_f \bar{f}_\beta \gamma^\mu L_b, \\
O_{3,5}^q &= \bar{q}_a \gamma_\mu L_b \bar{q}_\beta (\gamma_\mu L(R) q') \\
O_{4,6}^q &= \bar{q}_a \gamma_\mu L_b \bar{q}_\beta \gamma_\mu L(R) q', \\
O_{7,9}^q &= \frac{3}{2} \bar{q}_a \gamma_\mu L_b \bar{q}_\beta (\gamma^\mu R(L) q') \\
O_{8,10}^q &= \frac{3}{2} \bar{q}_a \gamma_\mu L_b \bar{q}_\beta (\gamma^\mu R(L) q'),
\end{align*} \]

(18)

where \( R(L) = 1 \pm \gamma_5 \), and \( q' \) is summed over all flavors except \( t \). \( O_{1f,2f} \) are the current-current operators that represent tree level processes. \( O_{3-6} \) are the strong gluon induced penguin operators, and operators \( O_{7-10} \) are due to \( \gamma \) and \( Z \) exchange (electroweak penguins), and “box” diagrams at loop level. The values of the Wilson coefficients can be found in Ref. [20].

In the factorization assumption the amplitude for \( B \to D^{(*)} D_{s0} D_{s1}^* \), can now be written as

\[ M = M_1 + M_2 \]

(19)

where

\[ \begin{align*}
M_1 &= \frac{G_F}{\sqrt{2}} X_1 < D_{s0}(D_{s1}^*) | \bar{s} \gamma_\mu (1 - \gamma_5) c | 0 > < D^{(*)} | \bar{c} \gamma_\mu (1 - \gamma_5) b | B > \\
M_2 &= \frac{G_F}{\sqrt{2}} X_2 < D_{s0}(D_{s1}^*) | \bar{s} (1 + \gamma_5) c | 0 > < D^{(*)} | \bar{c} (1 - \gamma_5) b | B >
\end{align*} \]

(20)

(21)

where

\[ \begin{align*}
X_1 &= V_c \left( \frac{c_1}{N_c} + c_2 \right) + \frac{B_3}{N_c} + B_4 + \frac{B_9}{N_c} + B_{10} \\
X_2 &= -2 \left( \frac{1}{N_c} B_5 + B_6 + \frac{1}{N_c} B_7 + B_8 \right)
\end{align*} \]

(22)

We have defined

\[ B_i = - \sum_{q=u,c,t} c_i^q V_q \]

(23)

with

\[ V_q = V_{qs}^* V_{qb} \]

(24)

In the above equations \( N_c \) represents the number of colors. To simplify matters we neglect the small penguin contributions and so as a first approximation we will neglect \( M_2 \).
We can now define the following ratios

\[
R_{D0} = \frac{BR[B \to DD_{s0}]}{BR[B \to DD_s]}
\]
\[
R_{D^*0} = \frac{BR[B \to D^*D_{s0}]}{BR[B \to D^*D_s]}
\]
\[
R_{D1} = \frac{BR[B \to DD^*_{s1}]}{BR[B \to DD^*_s]}
\]
\[
R_{D^*1} = \frac{BR[B \to D^*D^*_{s1}]}{BR[B \to D^*D^*_s]}
\]

(25)

Let us focus on the ratio \( R_{D0} \) which within factorization and the heavy quark limit can be written as

\[
R_{D0} = \left| \frac{f_{D_{s0}}}{f_{D_s}} \right|^2
\]

(26)

where we have neglected phase space (and other) effects that are subleading in the heavy quark expansion. Similarly we have

\[
R_{D1} = \left| \frac{f_{D^*_{s1}}}{f_{D^*_s}} \right|^2
\]

(27)

Now in the heavy quark limit \( f_{D_{s0}} = f_{D^*_{s1}} \) and \( f_{D_s} = f_{D^*_s} \) and so one would predict \( R_{D0} \approx R_{D1} \). There have been various estimates of the decay constant \( f_{D_{s0}} \) in quark models and in QCD sum rule calculations; these typically find the p-wave, \( j_q = \frac{1}{2} \) states to have the similar decay constants as the ground state mesons. We therefore expect \( f_{D_{s0}} \sim f_{D_s} \) giving in addition to the heavy quark predictions

\[
R_{D0} \approx R_{D1} \approx 1
\]

(28)

Experimentally Belle measures [22]

\[
BR[B \to DD_s(2317)]BR[D_s(2317) \to D_s\pi^0] = (9.9^{+2.5}_{-2.0} \pm 3.0) \times 10^{-4}
\]

The dominant decay of the \( D_s(2317) \) is expected to be through the \( D_s\pi \) mode [23] and so

\[
BR[D \to DD_s(2317)] \approx 10^{-3}
\]

(29)

Now using the measured branching ratio [24]

\[
BR[B^+ \to \bar{D}^0D^+_s] = (1.3 \pm 0.4) \times 10^{-2}
\]
\[
BR[B_d \to D^-D^*_s] = (8 \pm 3) \times 10^{-3}
\]

(30)
one obtains a combined branching ratio

$$BR[B \to DD_s] \approx 10^{-2}$$

(31)

This leads to $R_{D0} \approx \frac{1}{10}$ (or, $f_{D_{10}} \sim \frac{1}{3} f_{D_s}$) which is a factor 10 smaller than theoretical expectations.

One might argue that factorization is not applicable to $B \to D^{(*)}D^{(*)}$ decays. However, recent analysis in Ref. [25] find that factorization works well for these decays. Moreover, the quantities in Eq. (25) are ratios of nonleptonic decay amplitudes and so nonfactorizable effects may cancel. So what one really requires is significantly different nonfactorizable corrections between decays with the p-wave states in the final state and decays with the ground state mesons in the final state. It is possible that the discrepancies between experiments and theory may arise from a combination of incorrect model prediction of p-wave state properties and nonfactorizable effects.

If these new $D_{sJ}$ states are indeed the p-wave $c\bar{s}$ meson then they cannot contribute to $B \to D^{(*)}\bar{D}^{(*)}K_s$ as they are below the $D^{(*)}K$ threshold. Hence to the extent that the corrections to predictions of HHCHPT are small the three body decays $B \to D^{(*)}\bar{D}^{(*)}K_s$ should be dominated by the non resonant contribution. This would then imply that $\sin 2\beta$ can be cleanly measured in this decay. On the other hand if the new $D_{sJ}$ states are not the p-wave $c\bar{s}$ mesons but something exotic like four quark states or molecules then experimentally the real p-wave states should show up.

4 New Physics

There are many reasons to believe that the Standard Model is not a complete theory as it leaves several puzzles unresolved, specially in the flavour sector. There are several hints of possible deviations from the SM [20, 21] and several interesting methods have been proposed to measure the parameters of the underlying new physics [27]. If there is new physics in $b \to s$ transition then this should show up in decays with the underlying quark transition $b \to c\bar{c}s$ which can then be probed in $B \to D^{(*)}\bar{D}^{(*)}K_s$. One can probe for this new physics also in other decays with the same quark level transition such as $B \to J/\psi K_s$. However, it is possible that the matrix element of the new physics may be suppressed in some decays and enhanced in other making it mandatory to study as many different decays as possible. If the new physics involves a new effective $b \to sg$ vertex its effect in $B \to J/\psi K_s$ may be negligible because of OZI suppression. Such a suppression would not apply to $B \to D^{(*)}\bar{D}^{(*)}K_s$ decays.
5 Conclusion

In this talk we have pointed out that the three body decays \( B \to D^{(*)} \bar{D}^{(*)} K_s \) may be used to measure both \( \sin 2\beta \) and \( \cos 2\beta \). The \( \cos 2\beta \) measurement requires resonant contribution to the three body decay from p-wave excited \( D_s \) states. We discussed the newly discovered \( D_{s,J} \) states and pointed out that their properties are sometimes inconsistent with an interpretation as p-wave \( c\bar{s} \) states. If these p-wave states are the newly discovered \( D_s(2317) \) and \( D_s(2460) \), then they are below the \( D^{(*)}K \) threshold and hence do not contribute to \( B \to D^{(*)} \bar{D}^{(*)} K_s \). The three body decays can then be used to measure \( \sin 2\beta \) without resonant dilution and to look for new physics in \( b \to c\bar{c}s \) transition.

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