Flavor changing processes in quarkonium decays

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We study the flavor changing processes $\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$ in B factories and tau-charm factories. In the standard model, these processes are predicted to be unobservable; so they serve as a probe of new physics. We first perform a model-independent analysis, then examine the predictions of the models, such as the top-color models, the MSSM with R-parity violation, and the two-Higgs-doublet model, for the branching ratios of $\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$. We find that these branching ratios could be as large as 10^{-6} and 10^{-5} in the presence of new physics. [S0556-2821(99)00513-5]

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I. INTRODUCTION

The possibility of observing large *CP* violating asymmetries in the decay of *B* mesons has motivated the construction of high luminosity *B* factories at several of the world's high energy physics laboratories. These *B* factories will be producing roughly about 10^8 Υ's. Meanwhile, BES has already accumulated 9×10^6 J/ψ and plans to increase the number to 5×10^7 in the near future. An interesting question that we investigate in this paper is whether the large sample of the Υ and J/ψ can be used to probe flavor changing processes in the decays of Υ and J/ψ. In particular we look at the flavor changing processes $\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$, from the underlying $b\to s$ and $c\to u$ quark transitions. For the quarkonium system, these flavor changing processes are expected to be much smaller than in the case of decays of the *B* or *D* meson because of the larger decay widths of the bottonium and charmonium systems which decay via the strong interactions. Indeed the standard model contributions to $\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$ are tiny. However, new physics may enhance the branching ratios for these processes. Whether this enhancement will be sufficient for these processes to be observable in the next round of experiments is the subject of this work. The invisible decays of Υ and J/ψ resonances in the standard model and beyond have been studied recently [1].

Nonleptonic decays of heavy quarkonium systems can be more reliably calculated than nonleptonic decays of heavy mesons. A consistent and systematic formalism to handle heavy quarkonium decays is available in nonrelativistic QCD (NRQCD) [2] which is missing for heavy mesons. As in the meson system [3] it is more fruitful to concentrate on quasi-inclusive processes such as $\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$ because they can be calculated with less theoretical uncertainty and have larger branching ratios than purely exclusive quarkonium nonleptonic decays. The branching ratios of exclusive flavor changing nonleptonic decays of Υ and J/ψ in the standard model have been calculated and found to be very small [4].

We begin with a model-independent description of the processes $\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$. In the standard model these decays can proceed through tree and penguin processes. For new physics contributions to these processes we concentrate on four-quark operators of the type $\bar{s}b\bar{b}b$ and $\bar{u}c\bar{c}c$. We choose the currents in the four-quark operators to be scalars and so these operators may arise through exchange of a heavy scalar for, e.g., a Higgs boson or a leptoquark in some model of new physics. These four-quark operators, at the one-loop level, generate effective $\bar{s}b\{g,\gamma,Z\}$ and $\bar{u}c\{g,\gamma,Z\}$ vertices which would effect the flavor changing decays of the *B* and *D* mesons. The effective vertices for an on-shell g and γ vanish and so there is no contribution to $b\to s\gamma$ or $c\to u\gamma$. We can, however, put constraints on these operators by considering the processes $b\to sl^+l^-$ and $c\to ul^+l^-$. The constrained operators can then be used to calculate the branching ratios for $\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$.

We then consider some models that may generate the kind of four-quark operators described above. A few examples of models where these operators can be generated are top color models, minimal supersymmetric standard model (MSSM) with *R*-parity violation and a general two-Higgs-doublet model without any discrete symmetry. In some cases constraints on the parameters that appear in the prediction for the branching ratios for $\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$ are already available. In other cases the parameters are constrained, as in our model-independent analysis, from the processes $b\to sl^+l^-$ and $c\to ul^+l^-$.

In the sections which follow, we describe the effective Hamiltonian for the $\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$. Next we describe the calculation of the matrix elements and decay rates for these processes. We then discuss the calculation of the effective $\bar{s}b\{g,\gamma,Z\}$ and $\bar{u}c\{g,\gamma,Z\}$ vertices and constraints from the processes $b\to sl^+l^-$ and $c\to ul^+l^-$. This is followed by a description of some models that can generate the new four-quark operators in the effective Hamiltonian for

$\Upsilon \to B/\bar{B}X_s$ and $J/\psi \to D/\bar{D}X_u$. Finally we present our results and conclusions.

II. EFFECTIVE HAMILTONIAN

In this section we present the effective Hamiltonian for Υ decays. The effective Hamiltonian for charmonium decays can be written down by making obvious changes. In the standard model (SM) the amplitudes for hadronic Υ decays of the type $b\bar{b} \to s\bar{b} + \bar{s}b$ are generated by the following effective Hamiltonian [5,6]:

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

where the superscripts u,c,t indicate the internal quark, f can be a u or c quark, and 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

$$\begin{aligned}
O_{1f}^q &= \bar{q}_\alpha \gamma_\mu L f_\beta \bar{f}_\beta \gamma^\mu L b_\alpha, \\
O_{2f}^q &= \bar{q} \gamma_\mu L f \bar{f} \gamma^\mu L b, \\
O_{3,5}^q &= \bar{q} \gamma_\mu L b \bar{q}' \gamma_\mu L(R) q', \\
O_{4,6}^q &= \bar{q}_\alpha \gamma_\mu L b_\beta \bar{q}'_\beta \gamma_\mu L(R) q'_\alpha, \\
O_{7,9}^q &= \frac{3}{2} \bar{q} \gamma_\mu L b e_{q'} \bar{q}' \gamma^\mu R(L) q', \\
O_{8,10}^q &= \frac{3}{2} \bar{q}_\alpha \gamma_\mu L b_\beta e_{q'} \bar{q}'_\beta \gamma_\mu R(L) q'_\alpha,
\end{aligned} \tag{2}$$

where $R(L) = 1 \pm \gamma_5$, and q' is summed over all flavors except t. $O_{1f,2f}$ are the tree level and QCD-corrected operators. O_{3-6} are the strong gluon-induced-penguin operators, and the operators O_{7-10} are due to γ and Z exchange (electroweak penguins) and "box" diagrams at the loop level. The Wilson coefficients c_i^f are defined at the scale $\mu \approx m_b$ and have been evaluated to next-to-leading order in QCD. The c_i^t are the regularization scheme independent values obtained in Ref. [7]. We give the nonzero c_i^f below for $m_t = 176$ GeV, $\alpha_s(m_Z) = 0.117$, and $\mu = m_b = 5$ GeV,

$$\begin{aligned}
&c_1 = -0.307, \; c_2 = 1.147, \; c_3^t = 0.017, \\
&c_4^t = -0.037, \; c_5^t = 0.010, \; c_6^t = -0.045, \\
&c_7^t = -1.24 \times 10^{-5}, \; c_8^t = 3.77 \times 10^{-4}, \\
&c_9^t = -0.010, \; c_{10}^t = 2.06 \times 10^{-3}, \\
&c_{3,5}^{u,c} = -c_{4,6}^{u,c}/N_c = P_s^{u,c}/N_c, \; c_{7,9}^{u,c} = P_e^{u,c}, \; c_{8,10}^{u,c} = 0,
\end{aligned} \tag{3}$$

where N_c is the number of colors. The leading contributions to $P_{s,e}^i$ are given by $P_s^i = (\alpha_s/8\pi) c_2 [\frac{10}{9} + G(m_i, \mu, q^2)]$ and $P_e^i = (\alpha_{em}/9\pi)(N_c c_1 + c_2)[\frac{10}{9} + G(m_i, \mu, q^2)]$. The function $G(m, \mu, q^2)$ is given by

$$G(m, \mu, q^2) = 4 \int_0^1 x(1-x) \ln \frac{m^2 - x(1-x)q^2}{\mu^2} \mathrm{d}x. \tag{4}$$

All the above coefficients are obtained up to one-loop order in electroweak interactions. The momentum q is the momentum carried by the virtual gluon in the penguin diagram. When $q^2 > 4m^2$, $G(m, \mu, q^2)$ becomes imaginary. In our calculation, we use $m_u = 5$ MeV, $m_d = 7$ MeV, $m_s = 200$ MeV, and $m_c = 1.35$ GeV [8,9]. For $\Upsilon \to \bar{B}X_s$, the operators O_{1f}^q and O_{2f}^q do not contribute and the gluon momentum is fixed at $q^2 = M_\Upsilon^2$.

A similar expression for the standard model contribution to the flavor changing decays J/ψ can be written down.

To the standard model contribution we add higher-dimensional four-quark operators generated by physics beyond the standard model. In this paper, we consider the four-quark operators with two scalar currents:

$$L_{new} = \frac{R_1}{\Lambda^2} \bar{s}(1-\gamma^5) b \bar{b}(1+\gamma^5) b + \frac{R_2}{\Lambda^2} \bar{s}(1+\gamma^5) b \bar{b}(1-\gamma^5) b + \text{H.c.} \tag{5}$$

The four-quark operators in L_{new} are the product of two scalar currents. In Eq. (5), Λ represents the new physics scale and R_1 and R_2 are two free parameters which describe the strength of the contribution of the underlying new physics to the effective operators. In our analysis we will only keep dimension-6 operators suppressed by $1/\Lambda^2$ and neglect all higher dimension operators.

Using a Fierz transformation one can express the scalar-scalar combination in terms of vector-vector combination. For instance we can write

$$\begin{aligned}
&\bar{s}_\alpha (1-\gamma^5) b_\alpha \bar{b}_\beta (1+\gamma^5) b_\beta \\
&\quad = -\frac{1}{2} \bar{s}_\alpha \gamma_\mu (1+\gamma^5) b_\beta \bar{b}_\beta \gamma^\mu (1-\gamma^5) b_\alpha \\
&\quad = -\frac{1}{2N_c} \bar{s}_\alpha \gamma_\mu (1+\gamma^5) b_\alpha \bar{b}_\beta \gamma^\mu (1-\gamma^5) b_\beta \\
&\qquad - \bar{s}_\alpha T^a_{\alpha\beta} \gamma_\mu (1-\gamma^5) b_\beta \bar{b}_{\beta'} T^a_{\beta'\alpha'} \gamma^\mu (1-\gamma^5) b_{\alpha'},
\end{aligned} \tag{6}$$

where T^a are the SU(3) color matrices with the normalization $\mathrm{Tr}[T^aT^b]=\delta_{ab}/2$ and N_c is the number of colors. In the quarkonium system the leading component in the Fock space expansion involves the quark-antiquark pair being in a $^3S_1$ state, probed by the operator $\bar{s}\gamma_\mu(1+\gamma^5)b\bar{b}\gamma^\mu(1-\gamma^5)b$. Note the general Fock space expansion of quarkonium in NRQCD is [10]

$$|\psi_Q\rangle=O(1)|\bar{Q}Q\,[^3S_1^{(1)}]\rangle+O(v)|\bar{Q}Q[^3P_J^{(8)}]g\rangle+O(v^2)|\bar{Q}Q[^3S_1^{(1,8)}]gg\rangle+O(v^2)|\bar{Q}Q[^1S_0^{(8)}]g\rangle+O(v^2)|\bar{Q}Q[^3D_J^{(1,8)}]gg\rangle+\cdots, \tag{7}$$

where v is the velocity of the constituents in the quarkonium and g represents a dynamical gluon, i.e., one whose effects cannot be incorporated into an instantaneous potential and whose typical momentum is m_Qv^2. The low energy hadronization of the leading component in the Fock space expansion of the quarkonium takes place at $O(v^3)$. As to the other Fock states notice that the $|\bar{Q}Q[^3P_J^{(8)}]g\rangle$ configuration arises when the predominant state radiates a soft dynamical gluon. Such a process is mediated principally by the electric dipole operator, for which the selection rule is $L'=L\pm1$, $S'=S$, and which involves a single power of heavy quark three-momentum. Thus, the coefficient associated with this state is of order v. The electric dipole emission of yet another gluon involves a change from the P-wave state to the S- and D-wave states $|\bar{Q}Q[^3S_1^{(1,8)}]gg\rangle$, $|\bar{Q}Q\rangle,[^3D_J^{(1,8)}]gg\rangle$ and so the associated coefficients are of order v^2. Finally, the coefficient of the state $|\bar{Q}Q[^1S_0^{(8)}]g\rangle$ results from fluctuations into this spin-singlet state from the predominant spin-triplet state with the emission of a soft gluon via a spin flipping magnetic dipole transition. Such transitions involve the gluon three-momentum ($\sim m_Qv^2$) rather than the heavy quark three-momentum ($\sim m_Qv$), and therefore the associated coefficient is of order v^2. The low energy hadronization of these component in the Fock space expansion of the quarkonium takes place at $O(v^7)$. Also for the P wave, an additional factor of v comes from the derivative of the wave function.

Before concluding this section, we point out that in addition to the operator in Eq. (5) there are additional four-quark operators in the effective Lagrangian [11], such as those with vector-vector current structure, which contribute also to the processes we consider. We focus on operators in Eq. (5) because, as mentioned earlier, they can be generated by the exchanges of the new scalar bosons in models we consider below in Sec. V.

III. MATRIX ELEMENTS FOR $\Upsilon\to\bar{B}X_S$

We proceed to calculate the matrix elements of the form $\langle\bar{B}X_s|H_{eff}|\Upsilon\rangle$ which represents the process $\Upsilon\to\bar{B}X_s$ and where H_{eff} has been described above. The effective Hamiltonian consists of operators with a current $\times$ current structure. Pairs of such operators can be expressed in terms of color-singlet and color-octet structures. The factorization formalism based on NRQCD [2], which allows a systematic and consistent probe of the complete quarkonium Fock space, can then be used to calculate the Υ decay rate.

The matrix element of $\Upsilon\to\bar{B}X_s$ decay can be expressed as

$$\begin{aligned}M=&\frac{G_F}{\sqrt{2}}W_1\langle\bar{B}X_s|\bar{s}\gamma^\mu(1-\gamma^5)b|0\rangle\langle0|\bar{b}\gamma_\mu b|\Upsilon\rangle\\&+\frac{G_F}{\sqrt{2}}W_1'\langle\bar{B}X_s|\bar{s}\gamma^\mu(1+\gamma^5)b|0\rangle\langle0|\bar{b}\gamma_\mu b|\Upsilon\rangle\\&+\frac{G_F}{\sqrt{2}}W_8\langle\bar{B}X_s|\bar{s}\gamma^\mu(1-\gamma^5)T^ab|0\rangle\langle0|\bar{b}\gamma_\mu T^ab|\Upsilon\rangle\\&+\frac{G_F}{\sqrt{2}}W_8'\langle\bar{B}X_s|\bar{s}\gamma^\mu(1+\gamma^5)T^ab|0\rangle\langle0|\bar{b}\gamma_\mu T^ab|\Upsilon\rangle\\&+\frac{G_F}{\sqrt{2}}U_8\langle\bar{B}X_s|\bar{s}\gamma^\mu(1-\gamma^5)T^ab|0\rangle\langle0|\bar{b}\gamma_\mu\gamma^5T^ab|\Upsilon\rangle\\&+\frac{G_F}{\sqrt{2}}U_8'\langle\bar{B}X_s|\bar{s}\gamma^\mu(1+\gamma^5)T^ab|0\rangle\langle0|\bar{b}\gamma_\mu\gamma^5T^ab|\Upsilon\rangle\\&+\frac{G_F}{\sqrt{2}}V_8\langle\bar{B}X_s|\bar{s}(1+\gamma^5)T^ab|0\rangle\langle0|\bar{b}(1-\gamma^5)T^ab|\Upsilon\rangle\\&+\frac{G_F}{\sqrt{2}}V_8'\langle\bar{B}X_s|\bar{s}(1-\gamma^5)T^ab|0\rangle\langle0|\bar{b}(1+\gamma^5)T^ab|\Upsilon\rangle,\end{aligned} \tag{8}$$

where

$$\begin{aligned}W_1&=W_{1std}+W_{1new},\\W_1'&=W_{1std}'+W_{1new}',\\W_8&=W_{8std}+W_{8new},\\W_8'&=W_{8std}'+W_{8new}',\\U_8&=U_{8std}+U_{8new},\\U_8'&=U_{8std}'+U_{8new}',\\V_8&=V_{8std}+V_{8new},\\V_8'&=V_{8std}'+V_{8new}',\end{aligned} \tag{9}$$

with

$$W_{1std}=\left[\left\{A_3+A_4-\frac{1}{2}(A_9+A_{10})\right\}\left(1+\frac{1}{N_c}\right)+\left(A_5+\frac{A_6}{N_c}\right)-\frac{1}{2}\left(A_7+\frac{A_8}{N_c}\right)\right],$$

$$W_{1std}'=0,$$

$$W_{8std}=[2(A_3+A_4+A_6)-(A_8+A_9+A_{10})],$$

$$W'_{8std}=0,$$
$$U_{8std}=[2(-A_3-A_4+A_6)-(A_8-A_9-A_{10})],$$
$$U'_{8std}=0,$$
$$V_{8std}=[-4A_5+2A_7],$$
$$V'_{8std}=0 \tag{10}$$

and

$$W_{1new}=-\frac{1}{2N_c}\frac{\sqrt{2}}{G_F}\frac{R_2}{\Lambda^2},$$
$$W'_{1new}=-\frac{1}{2N_c}\frac{\sqrt{2}}{G_F}\frac{R_1}{\Lambda^2},$$
$$W_{8new}=-\frac{\sqrt{2}}{G_F}\frac{R_2}{\Lambda^2},$$
$$W'_{8new}=-\frac{\sqrt{2}}{G_F}\frac{R_1}{\Lambda^2},$$
$$U_{8new}=-\frac{\sqrt{2}}{G_F}\frac{R_2}{\Lambda^2},$$
$$U'_{8new}=\frac{\sqrt{2}}{G_F}\frac{R_1}{\Lambda^2},$$
$$V'_8=0. \tag{11}$$

We have defined

$$A_i=-\sum_{q=u,c,t}c_i^q V_q, \tag{12}$$

with

$$V_q=V_{qs}^* V_{qb}. \tag{13}$$

Similar expressions can be written for the matrix elements describing the J/ψ decay.

To calculate the decay rate we use the parton model to write the process $\Upsilon\to\bar{B}X_s$ as $\Upsilon(P)\to\bar{b}(p_1)s(p_2)$. The squared matrix element is then given by

$$|M|^2=2MZ_1\bigg[\langle\Upsilon|O_1(^3S_1)|\Upsilon\rangle(|W_1|^2+|W_1'^2|)+\left(\left\langle\Upsilon\left|O_8(^3S_1)+2\frac{O_8(^3P_1)}{m_b^2}\right|\Upsilon\right\rangle\right)(|W_8|^2+|W_8'^2|)\bigg]+6MZ_2[\langle\Upsilon|O_8(^1S_0)|\Upsilon\rangle(|U_8|^2+|U_8'^2|)]+6MZ_3[\langle\Upsilon|O_8(^1S_0)|\Upsilon\rangle(|V_8|^2+|V_8'^2|)] \tag{14}$$

where

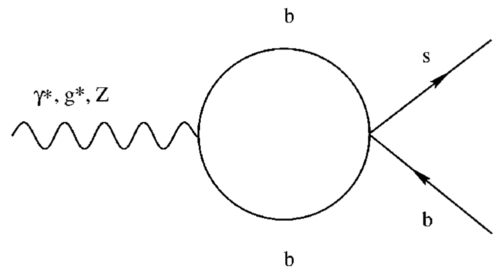

FIG. 1. Effective $\bar{s}b\gamma^*$, $\bar{s}bg^*$, $\bar{s}bZ$ vertices generated by L_{new}.

$$Z_1=8\left[p_1\cdot p_2+\frac{2p_1\cdot Pp_2\cdot P}{M^2}\right],$$
$$Z_2=8\left[-p_1\cdot p_2+\frac{2p_1\cdot Pp_2\cdot P}{M^2}\right],$$
$$Z_3=8[p_1\cdot p_2], \tag{15}$$

with M being the quarkonium mass. The matrix elements of the various color-singlet and color-octet operators $O_1(^{(2S+1)}L_J)$ and $O_8(^{(2S+1)}L_J)$ encode the nonperturbative long distance effects in the evolution of $\bar{Q}Q(^{(2S+1)}L_J)_{1,8}$ to Υ.

Along with the *CP* violating phases present in the standard model contribution there can be additional phases from the new contribution. We can then construct the *CP* violating rate asymmetry as

$$a_{CP}=\frac{\Gamma(\Upsilon\to\bar{B}X_s)-\Gamma(\Upsilon\to B\bar{X}_s)}{\Gamma(\Upsilon\to\bar{B}X_s)+\Gamma(\Upsilon\to B\bar{X}_s)}. \tag{16}$$

IV. LOW ENERGY CONSTRAINTS

The Lagrangian L_{new} generates, at the one-loop level, the effective $\bar{s}b\gamma^*$, $\bar{s}bg^*$, $\bar{s}bZ$ vertices as shown in Fig. 1, where γ^* and g^* indicate an off-shell photon and a gluon. Similar vertices involving $c\to u$ transitions are generated in the charmonium sector also. These vertices, with a γ and Z, will contribute to $b\to sl^+l^-$ and $c\to ul^+l^-$. Note that there is no contribution to $b\to s\gamma$. The vertex $b\to sg^*$ can give rise to the process $b\to sq\bar{q}$ which will contribute to nonleptonic B decays. We expect the constraints from $b\to sl^+l^-$ to be better than from nonleptonic B decays because of the theoretical uncertainties in calculating nonleptonic decays. The effective $\bar{s}b\gamma^*$, $\bar{s}bg^*$, $\bar{s}bZ$ Lagrangian can be written as

$$\delta L_{\bar{s}b\gamma^*}=\frac{e_b}{\Lambda^2}\bar{s}[R_+C_+^\mu+R_-C_-^\mu]b, \tag{17}$$

where

$$R_+=\frac{R_1+R_2}{2},$$

$$R_-=\frac{R_1-R_2}{2},$$

$$C_+=\frac{1}{16\pi^2}\int_0^1 dx\,\log\left(\frac{\Lambda^2}{B^2}\right)\times[8q^\mu\gamma\cdot qx(x-1)+8\gamma^\mu q^2x(1-x)],$$

$$C_-=\frac{1}{16\pi^2}\int_0^1 dx\,\log\left(\frac{\Lambda^2}{B^2}\right)\times[8q^\mu\gamma\cdot q\gamma^5x(x-1)+8\gamma^\mu\gamma^5q^2x(1-x)], \quad (18)$$

with

$$B^2=m_b^2-q^2x(1-x),$$

where e_b is the b-quark electric charge and q is the photon momentum. The effective $\bar{s}bg^*$ Lagrangian can be obtained by replacing e_b by g_s, the strong coupling constant. The effective $\bar{s}bZ$ Lagrangian can be written as

$$\delta L_{\bar{s}bZ}=\frac{g}{2c_w\Lambda^2}\bar{s}[R_+(g_LF_{1+}+g_RF_{2+})+R_-(g_LF_{1-}+g_RF_{2-})], \quad (19)$$

where

$$F_{1+}=\frac{1}{16\pi^2}\int_0^1 dx\,\log\left(\frac{\Lambda^2}{B^2}\right)[8q^\mu\gamma\cdot qx(x-1)(1+\gamma^5)+8\gamma^\mu q^2x(1-x)(1+\gamma^5)-8\gamma^\mu\gamma^5m_b^2],$$

$$F_{1-}=\frac{1}{16\pi^2}\int_0^1 dx\,\log\left(\frac{\Lambda^2}{B^2}\right)[8q^\mu\gamma\cdot qx(x-1)\times(1+\gamma^5)8\gamma^\mu q^2x(1-x)(1+\gamma^5)-8\gamma^\mu m_b^2],$$

$$F_{2+}=\frac{1}{16\pi^2}\int_0^1 dx\,\log\left(\frac{\Lambda^2}{B^2}\right)[8q^\mu\gamma\cdot qx(x-1)\times(1-\gamma^5)8\gamma^\mu q^2x(1-x)(1-\gamma^5)+8\gamma^\mu\gamma^5m_b^2],$$

$$F_{2-}=\frac{1}{16\pi^2}\int_0^1 dx\,\log\left(\frac{\Lambda^2}{B^2}\right)[8q^\mu\gamma\cdot qx(1-x)\times(1-\gamma^5)8\gamma^\mu q^2x(x-1)(1-\gamma^5)+8\gamma^\mu m_b^2], \quad (20)$$

with

$$g_L=-\frac{1}{2}+\frac{1}{3}s_w^2,$$

$$g_R=\frac{1}{3}s_w^2.$$

For $b\to sl^+l^-$ the q^μ terms in the equations above do not contribute if we neglect the lepton masses. Furthermore, the contribution from Z exchange is suppressed with respect to γ exchange by a factor of q^2/M_Z^2 and so we do not include the Z contribution. The additional contribution to the effective Hamiltonian for $b\to sl^+l^-$ can be written as

$$\delta H_{b\to sl^+l^-}=-\frac{e^2}{16\pi^2}\frac{e_b}{\Lambda^2}\int_0^1 dx8x(1-x)\log\left(\frac{\Lambda^2}{B^2}\right)\times[R_1\bar{s}\gamma^\mu b_L\bar{l}\gamma_\mu l+R_2\bar{s}\gamma^\mu b_R\bar{l}\gamma_\mu l], \quad (21)$$

which has to be added to the standard model contribution [6]. Similar results can also be written for the charm sector.

V. MODELS

In this section we look at various models that can give rise to L_{new} given in Eq. (5). As a first example we consider a recent version of top color models [12]. In such models the top quark participates in a new strong interaction which is broken at some high energy scale Λ. The strong interaction, though not confining, leads to the formation of a top condensate $\langle\bar{t}_Lt_R\rangle$ resulting in a large dynamical mass for the top quark. The scale Λ is chosen to be of the order of 1 TeV to avoid the naturalness problem which implies that the electroweak symmetry cannot be broken solely by the top condensate. In the low energy sector of the theory, scalar bound states are formed that couple strongly to the b quark [13,14]:

$$L_b=\frac{m_t}{f_{\tilde{\pi}}\sqrt{2}}\bar{b}_L(H+iA^0)b_R+\text{H.c.}, \quad (22)$$

where $f_{\tilde{\pi}}\sim 50$ GeV is the top pion decay constant. On integrating out the Higgs fields H and A^0 we have an effective four-fermion operator

$$L_{eff}=\frac{m_t^2}{f_{\tilde{\pi}}^2m_H^2}\bar{b}_Lb_R\bar{b}_Rb_L. \quad (23)$$

Since the b quark in Eq. (22) is in the weak eigenstate, L_{eff} in Eq. (22) will induce flavor changing neutral current (FCNC) four-quark operators in Eq. (5) after diagonalizing the quark mass matrix [14], with coefficients

$$R_1=\frac{1}{4}\frac{m_t^2}{f_{\tilde{\pi}}^2m_H^2}|D_{Lbb}|^2D_{Rbb}D_{Rbs}^*,$$

$$R_2=\frac{1}{4}\frac{m_t^2}{f_{\tilde{\pi}}^2m_H^2}|D_{Rbb}|^2D_{Lbb}D_{Lbs}^*, \quad (24)$$

where D_L and D_R are the mixing matrices in the left and right handed down sector. In the charm sector similar interactions can arise due to the strong couplings of the top quark to top pions. The effective operators generated by integrating out the top pions are similar to Eq. (5) with the replacement of b by c and s by u. In top color II models [14,15], where there can be strong top pion couplings of the top quark with the charm quark, we have

$$R_1=\frac{1}{4}\frac{m_t^2}{f_{\tilde{\pi}}^2 m_{\tilde{\pi}}^2}|U_{Lcc}|^2 U_{Rtc}U_{Rtu}^*,$$

$$R_2=\frac{1}{4}\frac{m_t^2}{f_{\tilde{\pi}}^2 m_{\tilde{\pi}}^2}|U_{Rtc}|^2 U_{Lcu}U_{Lcc}^*. \tag{25}$$

In supersymmetric standard models without R parity, the most general superpotential of the MSSM, consistent with $SU(3)\times SU(2)\times U(1)$ gauge symmetry and supersymmetry, can be written as

$$\mathcal{W}=\mathcal{W}_R+\mathcal{W}_{\not R}, \tag{26}$$

where $\mathcal{W}_R$ is the R-parity conserving part while $\mathcal{W}_{\not R}$ violates the R parity. They are given by

$$\mathcal{W}_R=h_{ij}L_iH_2E_j^c+h'_{ij}Q_iH_2D_j^c+h''_{ij}Q_iH_1U_j^c, \tag{27}$$

$$\mathcal{W}_{\not R}=\lambda_{ijk}L_iL_jE_k^c+\lambda'_{ijk}L_iQ_jD_k^c+\lambda''_{ijk}U_i^cD_j^cD_k^c + \mu_iL_iH_2. \tag{28}$$

Here $L_i(Q_i)$ and $E_i(U_i,D_i)$ are the left-handed lepton (quark) doublet and lepton (quark) singlet chiral superfields, with i,j,k being generation indices and c denoting a charge conjugate field. $H_{1,2}$ are the chiral superfields representing the two Higgs doublets. In the R-parity violating superpotential above, the λ and λ' couplings violate lepton-number conservation, while the λ'' couplings violate baryon-number conservation. λ_{ijk} is antisymmetric in the first two indices and λ''_{ijk} is antisymmetric in the last two indices. While it is theoretically possible to have both baryon-number and lepton-number violating terms in the Lagrangian, the nonobservation of proton decay imposes very stringent conditions on their simultaneous presence [16]. We therefore assume the existence of either L violating couplings or B violating couplings, but not the coexistence of both. We calculate the effects of both types of couplings.

In terms of the four-component Dirac notation, the Lagrangian involving the λ' and λ'' couplings is given by

$$\mathcal{L}_{\lambda'}=-\lambda'_{ijk}[\tilde{\nu}_L^i\overline{d}_R^k d_L^j+\tilde{d}_L^j\overline{d}_R^k\nu_L^i+(\tilde{d}_R^k)^*(\overline{\nu}_L^i)^c d_L^j-\tilde{e}_L^i\overline{d}_R^k u_L^j - \tilde{u}_L^j\overline{d}_R^k e_L^i-(\tilde{d}_R^k)^*(\overline{e}_L^i)^c u_L^j]+\text{H.c.}, \tag{29}$$

$$\mathcal{L}_{\lambda''}=-\lambda''_{ijk}[\tilde{d}_R^k(\overline{u}_L^i)^c d_L^j+\tilde{d}_R^j(\overline{d}_L^k)^c u_L^i+\tilde{u}_R^i(\overline{d}_L^j)^c d_L^k]+\text{H.c.} \tag{30}$$

The terms proportional to λ are not relevant to our present discussion and will not be considered here. The exchange of sneutrinos with the λ' coupling will generate L_{new} for $\Upsilon\to\overline{B}X_s$ with

$$R_1=\frac{1}{4}\Sigma_i\frac{\lambda'_{i32}\lambda'^*_{i33}}{m_{\tilde{\nu}_i}^2},$$

$$R_2=\frac{1}{4}\Sigma_i\frac{\lambda'^*_{i23}\lambda'_{i33}}{m_{\tilde{\nu}_i}^2}. \tag{31}$$

For the case of $J/\psi\to DX_u$ the operators in L_{new} cannot be generated at the tree level.

Another model of interest is an extension of the SM with additional scalar SU(2) doublets; the simplest of these would be the two-Higgs-doublet model (2HDM). In general, when the quarks couple to more than one scalar doublet, there are inevitably FCNC couplings to the neutral scalars. When the up-type quarks and down-type quarks are allowed simultaneously to couple to more than one scalar doublet, the diagonalization of the up-type and down-type mass matrices does not automatically ensure the diagonalization of the couplings with each single scalar doublet. Frequently, as in the Weinberg model for *CP* violation [17] or in supersymmetry (SUSY), the 2HDM scalar potential and Yukawa Lagrangian are constrained by an *ad hoc* discrete symmetry [18], whose only role is to protect the model from FCNC at the tree level. Let us consider a Yukawa Lagrangian of the form

$$\mathcal{L}_Y^{(A)}=\eta_{ij}^U\overline{Q}_{i,L}\tilde{\phi}_1U_{j,R}+\eta_{ij}^D\overline{Q}_{i,L}\phi_1D_{j,R}+\xi_{ij}^U\overline{Q}_{i,L}\tilde{\phi}_2U_{j,R}+\xi_{ij}^D\overline{Q}_{i,L}\phi_2D_{j,R}+\text{H.c.}, \tag{32}$$

where ϕ_i, for $i=1,2$, are the two scalar doublets of a 2HDM, while $\eta_{ij}^{U,D}$ and $\xi_{ij}^{U,D}$ are the nondiagonal matrices of the Yukawa couplings.

When no discrete symmetry is imposed, then both up-type and down-type quarks can have FC couplings [19]. Such models were called class A in [20], to be contrasted with models with a forced absence of FCNC, called class B.

In the notation and basis of Ref. [21] the flavor changing part of the Lagrangian can be written as

$$\mathcal{L}_{Y,FC}^{(III)}=\hat{\xi}_{ij}^U\overline{Q}_{i,L}\tilde{\phi}_2U_{j,R}+\hat{\xi}_{ij}^D\overline{Q}_{i,L}\phi_2D_{j,R}+\text{H.c.}, \tag{33}$$

where $Q_{i,L}$, $U_{j,R}$, and $D_{j,R}$ denote now the quark mass eigenstates and $\hat{\xi}_{ij}^{U,D}$ are the rotated couplings, in general not diagonal. If we define $V_{L,R}^{U,D}$ to be the rotation matrices acting on the up- and down-type quarks, with left or right chirality, respectively, then the neutral FC couplings will be

$$\hat{\xi}_{\text{neutral}}^{U,D}=(V_L^{U,D})^{-1}\cdot\xi^{U,D}\cdot V_R^{U,D}. \tag{34}$$

On the other hand, for the charged FC couplings we will have

$$\hat{\xi}_{\text{charged}}^U=\hat{\xi}_{\text{neutral}}^U\cdot V_{\text{CKM}},$$

$$\hat{\xi}_{\text{charged}}^D=V_{\text{CKM}}\cdot\hat{\xi}_{\text{neutral}}^D, \tag{35}$$

where V_{CKM} denotes the Cabibbo-Kobayashi-Maskawa matrix.

The phenomenology for the 2HDM, for the quarkonium processes under study, is not very different from the other two models considered above and so we will concentrate mainly on the top color models and supersymmetry models with R-parity violation.

TABLE I. Input parameters used in our calculations.

NRQCD matrix elements	Value
$\langle O_1^{\psi}(^3S_1)\rangle \approx 3\langle\psi\|O_1(^3S_1)\|\psi\rangle$	$0.73\ \mathrm{GeV}^3$
$\langle\Upsilon\|O_1(^3S_1)\|\Upsilon\rangle$	$2.3\ \mathrm{GeV}^3$
$\langle O_8^{\psi}(^3S_1)\rangle$	$0.014\ \mathrm{GeV}^3$
$\langle O_8^{\psi}(^1S_0)\rangle \approx \langle O_8^{\psi}(^3P_0)\rangle/m_c^2$	$10^{-2}\ \mathrm{GeV}^3$
$\langle\Upsilon\|O_8(^3S_1)\|\Upsilon\rangle$	$5\times10^{-4}\ \mathrm{GeV}^3$
$\langle O_8^{\Upsilon}(^1S_0)\rangle \approx \langle O_8^{\Upsilon}(^3P_0)\rangle/m_b^2$	$7\times10^{-3}\ \mathrm{GeV}^3$

VI. RESULTS AND DISCUSSION OF THEORETICAL UNCERTAINTIES

In this section we discuss the results of our calculations. First let us look at the Υ decays. The inputs to our calculation are the various well-known NRQCD matrix elements given in Table I [22,23]. The standard model contribution to the branching ratio is 5.2×10^{-11} from the penguin-induced $b\to s$ transition. The process $\Upsilon\to\bar{B}X_s$ can also have a contribution in the standard model from tree level processes. The effective Hamiltonian, suppressing the Dirac structure of the currents,

$$H_W=\frac{G_F}{\sqrt{2}}V_{ub}V_{us}^{*}[a_1(\bar{u}b)(\bar{s}u)+a_2(\bar{s}b)(\bar{u}u)], \tag{36}$$

where a_1 and a_2 are the QCD coefficients can generate the process $\Upsilon\to B^+K^-$. We can estimate the branching ratio for this process as

$$\mathcal{B}[\Upsilon\to B^+K^-]\approx\left|\frac{V_{ub}}{V_{cb}}\right|^2\mathcal{B}[\Upsilon\to B_c^+K^-].$$

Using $\mathcal{B}[\Upsilon\to B_c^+K^-]$ calculated in Ref. [4] one obtains $\mathcal{B}[\Upsilon\to B^+K^-]\sim1.5\times10^{-14}$. For a rough estimate of $\mathcal{B}[\Upsilon\to B^+X_s]$ we can scale $\mathcal{B}[\Upsilon\to B^+K^-]$ by the factor $\mathcal{B}[B\to D^0X]/\mathcal{B}[B\to D\pi]$. The measured value of $\mathcal{B}[B\to D^0X]$ [9] includes D^0 coming from the decay of D^{0*} and D^{+*}. From the spin phase factors $\mathcal{B}[B\to D^*X]\sim3\mathcal{B}[B\to DX]$. Hence $\mathcal{B}[B\to D^0X]/\mathcal{B}[B\to D\pi]\sim20$, leading to $\mathcal{B}[\Upsilon\to B^+X_s]\sim3\times10^{-13}$.

So far we have not considered R_1 and R_2 in L_{new}. In our model-independent analysis we vary R_1/Λ^2, R_2/Λ^2 one at a time and then use the constraint from measurements of $b\to se^+e^-$ and $b\to s\mu^+\mu^-$. We identify Λ with the masses of the exchange particles which we take to be between 100 and 200 GeV. We also take the cutoff for the integral in Eqs. (17)–(20) as 200 GeV. The allowed values of R_1/Λ^2, R_2/Λ^2 are then used to calculate $\Upsilon\to\bar{B}X_s$. The constraint from $b\to sl^+l^-$ gives

$$|R_{1,2}|/\Lambda^2\langle(6-9)\times10^{-6}(1/\mathrm{GeV})^2.$$

Using the upper bounds on $|R_{1,2}|/\Lambda^2$ we find the branching ratio for the process $\Upsilon(1S)\to\bar{B}X_s$ to be in the range $(1-2)\times10^{-6}$. Branching ratios of similar order are also obtained for $\Upsilon(2S)$ and $\Upsilon(3S)$. For $\Upsilon(4S)$ the branching ratio is smaller by a factor of 100 because of the larger width of $\Upsilon(4S)$ which decays predominantly to two B mesons.

Turning now to models, we find that for the top color model from Eq. (24) we can write

$$D_{Rbs}^{*}=4\frac{R_1}{\Lambda^2}\frac{f_{\tilde{\pi}}^2m_H^2}{m_t^2|D_{Lbb}|^2D_{Rbb}}. \tag{37}$$

We can identify $\Lambda=m_H$ and use the constraint from $b\to se^+e^-$ for a typical value of $|R_1|/\Lambda^2\sim6\times10^{-6}(1/\mathrm{GeV})^2$. Assuming $|D_{Lbb}|\approx|D_{Rbb}|\approx1$, and $f_{\tilde{\pi}}=50$ GeV we obtain $|D_{Rbs}|\sim2m_H^2\times10^{-6}$. With typical values of $m_H\sim100$–200 GeV we get $|D_{Rbs}|\sim0.02$–0.08. Similar values have been obtained for $|D_{Rbs}|$ in Ref. [14] by considering the contributions of the charged Higgs boson and top pion to $b\to s\gamma$. A similar exercise can be carried out with $|D_{Lbs}|$. Note that B_s mixing probes the combination $D_{Lbs}^{*}D_{Rbb}D_{Rbs}^{*}D_{Lbb}$ and so by either choosing $R_1\sim0$ or $R_2\sim0$ we can satisfy the constraint on B_s mixing by choosing the appropriate mixing elements to be small. Note that in top color models we can have operators $\bar{s}(1-\gamma^5)b\bar{d}(1+\gamma^5)d$ and $\bar{s}(1+\gamma^5)b\bar{d}(1-\gamma^5)d$ that can contribute to $\Upsilon\to\bar{B}s\bar{d}\to\bar{B}X_s$ after Fierz reordering. However, these operators will be suppressed by form factor effects and also from mixing effects. We have checked that the contribution to $\Upsilon\to\bar{B}X_s$ from these operators is much suppressed relative to the contribution of the operators in L_{new}. We will therefore not consider the above operators in our analysis.

Turning to R-parity violating SUSY we first collect the constraints on the relevant couplings. The upper limits of the L violating couplings for the squark mass of 100 GeV are given by

$$|\lambda'_{kij}|<0.012\quad(k,j=1,2,3;i=1,2), \tag{38}$$

$$|\lambda'_{13j}|<0.16\quad(j=1,2), \tag{39}$$

$$|\lambda'_{133}|<0.001, \tag{40}$$

$$|\lambda'_{23j}|<0.16\quad(j=1,2,3), \tag{41}$$

$$|\lambda'_{33j}|<0.26\quad(j=1,2,3). \tag{42}$$

The first set of constraints in Eq. (38) comes from the decay $K\to\pi\nu\nu$ with FCNC processes in the down quark sector [24]. The set of constraints in Eq. (39) and Eq. (41) is obtained from the semileptonic decays of B mesons [25]. The constraint on the coupling λ'_{133} in Eq. (40) is obtained from the Majorana mass that the coupling can generate for the electron-type neutrino [26]. The last set of limits in Eq. (42) is derived from the leptonic decay modes of the Z [27]. Assuming all the couplings to be positive we find the branching ratio for $\Upsilon\to\bar{B}X_s$ to be around 2×10^{-6} for $m_{\tilde{\nu}}=100$ GeV.

Turning next to $J/\psi\to\bar{D}X_u$, we first make an estimate for this process in the standard model. Since the penguin $c\to u$ transition is small in the standard model, we neglect its contribution. As in the case for the Υ system, for a rough estimate, can write

$$\mathcal{B}[J/\psi\to D^0X_u] \sim\mathcal{B}[J/\psi\to D^0\pi^0]\mathcal{B}[D^0\to K^-X]/\mathcal{B}[D^0\to K^-\pi^+].$$

We obtain $\mathcal{B}[J/\psi\to D^0\pi^0]$ from [4] and keeping in mind that $\mathcal{B}[D^0\to K^-X]$ contains contributions from states decaying to K^- we obtain $\mathcal{B}[J/\psi\to D^0X_u]\sim 10^{-10}$. A similar exercise gives $\mathcal{B}[J/\psi\to D^+X_u]\sim 10^{-9}$.

Considering new physics effects we can constrain R_1 and R_2 from $c\to ul^+l^-$. We get an estimate of the constraint on $c\to ue^+e^-$ by adding up the exclusive modes:

$$\mathcal{B}[D\to ue^+e^-]\geqslant\mathcal{B}[D\to(\pi^0+\eta+\rho^0+\omega)e^+e^-].$$

From $c\to ul^+l^-$ one obtains

$$|R_{1,2}|/\Lambda^2\leqslant 3.7\times 10^{-4}(1/\mathrm{GeV})^2.$$

We find that the branching fraction for the process $J/\psi\to\bar{D}X_u$ using the constraint from $c\to ul^+l^-$ can be (3–4) $\times 10^{-5}$.

In top color models taking R_1 and R_2 one at a time, one obtains

$$\frac{2.1\times 10^3}{m_{\tilde{\pi}}^4}||U_{Lcc}|^2U_{Rtc}U_{Rtu}^*|^2$$

or

$$\frac{2.1\times 10^3}{m_{\tilde{\pi}}^4}||U_{Rtc}|^2U_{Lcc}U_{Lcu}^*|^2$$

as the branching fraction for $J/\psi\to\bar{D}X_u$. For $m_{\tilde{\pi}}$ between 100 and 200 GeV this rate can be in the range (0.1–2.0) $\times 10^{-5}$ if all the mixing angles are $\sim 1°$. It has been shown in Ref. [28] that our choices for $f_{\tilde{\pi}}$ and $m_{\tilde{\pi}}$ give unacceptably large corrections to $Z\to b\bar{b}$ from one-loop contributions of the top pions. However, in a strongly coupled theory higher loop terms can have significant contributions. Nonetheless, if we change $f_{\tilde{\pi}}$ to ~ 100 GeV for better agreement with $Z\to b\bar{b}$ data, then the effect in $J/\psi\to\bar{D}X_u$ is reduced by a factor of 16. As in the case of the Υ system we can satisfy the constraint from D mixing by choosing $R_1\sim 0$ or $R_2\sim 0$.

For R-parity violating SUSY, the contribution to $J/\psi\to\bar{D}X_u$ can only occur at the loop level, with both λ' or λ'' contributing, through the box diagram, and so is suppressed. However, in the general 2HDM, from Eq. (33), the operator $\bar{u}c\bar{c}c$ can be generated by the tree level exchange of the field ϕ_2. The contribution to $J/\psi\to\bar{D}X_u$ will be proportional to the combination of couplings $R_2=\hat{\xi}_{12}^U\hat{\xi}_{22}^{U*}$ and $R_1=\hat{\xi}_{21}^{U*}\hat{\xi}_{22}^U$. Note the D^0-$\bar{D}^0$ mixing probes $\hat{\xi}_{12}^U\hat{\xi}_{21}^{U*}$. So we can satisfy the constraint on D^0-$\bar{D}^0$ mixing by choosing either $R_1\sim 0$ or $R_2\sim 0$. One then obtains

$$\frac{14}{m_H^4}|\hat{\xi}_{12}^U\hat{\xi}_{22}^{U*}|^2$$

or

$$\frac{14}{m_H^4}|\hat{\xi}_{21}^{U*}\hat{\xi}_{22}^U|^2$$

as the branching ratio for $J/\psi\to\bar{D}X_u$. For m_H between 100 and 200 GeV this rate can be in the range $(0.1–1.4)\times 10^{-7}$ if all the couplings are ~ 1. In a strongly interacting theory these couplings can be >1 as in the example of top color models discussed above.

We note in passing that the four-quark operators we have considered can also give rise to mixing operators that are $1/\Lambda^4$ suppressed. Taking one operator at a time one can generate the following operators that contribute to mixing:

$$O_1=-\frac{3R_1^2}{2\pi^2\Lambda^2}\frac{m_b^2}{\Lambda^2}\log\left(\frac{\Lambda^2}{m_b^2}\right)\bar{s}(1-\gamma^5)b\bar{s}(1-\gamma^5)b,$$

$$O_2=-\frac{3R_2^2}{2\pi^2\Lambda^2}\frac{m_b^2}{\Lambda^2}\log\left(\frac{\Lambda^2}{m_b^2}\right)\bar{s}(1+\gamma^5)b\bar{s}(1+\gamma^5)b. \tag{43}$$

We then have, for B_s mixing,

$$\Delta B_s=\frac{5}{3}f_{B_s}{}^2\eta BM_{B_s}\delta, \tag{44}$$

where η is the QCD correction factor and B and δ are defined through

$$\langle B_s^0|\bar{s}(1-\gamma^5)b\bar{s}(1-\gamma^5)|\bar{B}_s^0\rangle=-\frac{5}{3}f_{B_s}{}^2BM_{B_s}{}^2$$

and

$$\delta=-\frac{3R_i^2}{2\pi^2\Lambda^2}\frac{m_b^2}{\Lambda^2}\log\left(\frac{\Lambda^2}{m_b^2}\right),$$

where $i=1,2$. A similar result can be written for D mixing.

For the Υ system, if we include the constraint from B_s mixing generated by only the operators $O_{1,2}$, then we obtain $|R_{1,2}|/\Lambda^2>1–2\times 10^{-6}(1/\mathrm{GeV})^2$. This is of the same order as the constraint obtained from $b\to sl^+l^-$. In the case of J/ψ, if we include the constraint from the D mixing generated by only the operators $O_{1,2}$, then we obtain $|R_{1,2}|/\Lambda^2\sim 10^{-6}(1/\mathrm{GeV})^2$. This will lower the branching fraction for $J/\psi\to\bar{D}X_u$ to $\sim 10^{-9}$. However, to evaluate consistently the effects of these operators at order of $(1/\Lambda^2)^2$ would require the addition of other operators with dimension $\leqslant 8$. As noted earlier we restrict our analysis to only dimension-6 operators in the effective Lagrangian; thus the conservative result for $J/\psi\to\bar{D}X_u$ is of order of 10^{-5}, which, as we shown above, lies also in the region predicted by the top color model and 2HDM.

Direct CP violation is possible in these decays through the interference of the standard model and new physics contributions. The CP conserving phase is generated at the quark level from the penguin diagrams. However, the standard model contribution is small and so the CP asymmetry a_{CP} is also small with a typical value of 0.1% for $\Upsilon \to \bar{B}X_s$.

In summary we have calculated branching ratios for the flavor changing processes $\Upsilon \to \bar{B}X_s$ and $J/\psi \to \bar{D}X_u$. In a model-independent description of new physics [29], constrained by low energy data from $b \to s l^+ l^-$ and $c \to u l^+ l^-$, we found branching fractions for these processes can be $\sim 10^{-6}$ and $\sim 10^{-5}$ for Υ and J/ψ decays, respectively. We also discussed several models of new physics that can allow these processes to occur with branching ratios that may be measurable in the next round of experiments.

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