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Decay Widths of Excited Heavy Mesons in The Leading Order of $1/M_Q$ Expansion¹

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ABSTRACT

General formulas for decay widths of heavy mesons with arbitrary spin and parity to the heavy meson ground state plus a π or K are derived in the leading order of the $1/M$ expansion in the relativistic B-S equation. Numerical results for various partial decay widths of 1^+ , 2^+ doublets in the D , D_s , B , and B_s series are obtained with a kernel used previously to obtain the spectra of low-lying states in these series in good agreement with experiments. The results are in satisfactory agreement with the experimental data when they are available.

In recent years the physics of hadrons containing a heavy quark has attracted much attention. This followed from the discovery of the heavy spin-flavor symmetry[1] that emerges from QCD as the heavy quark mass $M_Q \rightarrow \infty$. Systematic expansion in $1/M_Q$ can be carried out with the heavy quark effective theory (HQET)[2]. Such systems are interesting also for other physical reasons. A heavy meson composed of a heavy quark and a light anti-quark is the simplest system for studying the idea of constituent quark. Furthermore, with the B factories being available within several years, the theoretical predictions for properties of such mesons are useful for future experiments.

HQET provides a general frame-work for studying such systems. The expansions obtained from this approach contains a series of parameters and form-factors. When available experimental data are sufficient for determining the relevant parameters, they can be used to make reliable predictions. Many interesting results have been

obtained by this method. When experimental data are insufficient for the problem under investigation, one must calculate the parameters by some non-perturbative method.

A formalism by the $1/M_Q$ expansion in the relativistic B-S equation for the heavy meson has been developed in [3-6]. With some assumed kernel, it can be used to calculate some parameters and form-factors in the $1/M_Q$ expansion. It has been shown in [4] in the leading order of $1/M_Q$ expansion that, if the interaction becomes independent of the heavy quark spin as $M_Q \rightarrow \infty$, the number of independent components of the most general form of the B-S wave functions for heavy mesons of arbitrary spin-parity is reduced from eight to two. This greatly simplifies the calculations in the leading order. The simple form of the leading order wave functions also makes the calculations for corrections in the order $O(1/M_Q)$ easier. In this approach, the light component in the heavy meson is treated as relativistic. In [5], spectra of low-lying states in the D , D_s , B and B_s series are calculated with a kernel which is reduced to the one-gluon exchange in the Coulomb gauge and a static confinement term in the rest-frame of the meson. Numerical results up to next-to leading order in the $1/M_Q$ expansion in good agreement with experimental data for 12 observed states are obtained with 7 parameters, the values of which are close to that used in the quarkonium system and other problems. In [6] the decay widths for the processes $D^* \rightarrow D \pi$ and $D^* \rightarrow D \gamma$ were calculated to the order $1/M_Q$. The ratio $\text{Br}(D^{*0} \rightarrow D^0 \pi^0)/\text{Br}(D^{*0} \rightarrow D^0 \gamma)$ was used to determine the axial-vector coupling constant of pion to the light quark.

In the present article, the general formulas for partial decay widths of heavy mesons with arbitrary spin-parity to the ground state plus one π or K will be given in the leading order of $1/M_Q$ expansion in B-S equation. The parameters determined in [5,6] will be used to calculate various decay widths of 1^+ , 2^+ doublets in the D , D_s , B and B_s series in the leading order of $1/M_Q$. These quantities have been calculated in [7] under the assumption that s quark satisfies the heavy quark symmetry and some other assumptions. As pointed out in [7], these decay widths are interesting for CP violation experiments on B factories.

The covariant B-S equation for a heavy meson is of the following form.

$$\chi_P(p) = S_h(\lambda_1 P + p) \int G(P, p, q) \chi_P(q) \frac{d^4 q}{(2\pi)^4} S_l(-\lambda_2 P + p) \quad (1)$$

where $\lambda_1 = \frac{M_Q}{M_Q + m}$, $\lambda_2 = \frac{m}{M_Q + m}$, M_Q is the mass of the heavy quark, m is a mass parameter of the light quark. $P = Mv$, M is the meson mass, v is the meson velocity, p is the relative momentum. Let $p_l = v \cdot p$, $p_t = p - p_l v$. For convenience, we shall shift the p_t variable by defining $p'_t = p_t - \frac{m}{M_Q + m} M$ and omit the " l " in the following.

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We expand the quantities in B-S equation in series of $1/M$,

$$\begin{aligned}\chi(p) &= \chi_0 + \frac{1}{M_Q} \chi_1 + O\left(\frac{1}{M_Q^2}\right), \\ G &= G_0 + \frac{1}{M_Q} G_1 + O\left(\frac{1}{M_Q^2}\right), \\ M - M_Q - m &= E_0 + \frac{1}{M_Q} E_1 + O\left(\frac{1}{M_Q^2}\right).\end{aligned}\quad (2)$$

In the case $M_Q \gg p_t = O(\Lambda_{QCD})$ and $M_Q \gg m$, the residue of $\chi p(p)$ in (1) at the pole $p_t \approx -2M_Q$ of the heavy quark propagator in the upper half p_t plane is of the order $1/M_Q^2$. Neglecting this pole and consider $p_t = (\Lambda_{QCD})$, we obtain

$$S_h(\lambda_1 P + p) = \frac{Z(1 + \hat{v})}{2(p_t + E_0 + m + i\epsilon)} + O\left(\frac{1}{M_Q}\right), \quad (3)$$

where Z is a function of p_t and p_t^2 . It depends on M_Q only logarithmically. Let $G_0 = G_+ + G_-$, where the γ matrix on the heavy quark line of G_+ and G_- commutes and anti-commutes with $\hat{v}$ respectively. From equation (1)-(3), we obtain the equation for the leading wave function

$$\chi_0(p) = \frac{iZ(1 + \hat{v})}{2(p_t + E_0 + m + i\epsilon)} \int G_+(P, p, q) \chi_0(q) \frac{d^4 q}{(2\pi)^4} S_l(v, p) \quad (4)$$

It was shown in [4] that if G_+ does not contain terms proportional to $\gamma_5 \gamma_{\mu\nu}$ or $\frac{1}{2}(\gamma_{\mu\nu} \gamma_{\mu\nu} - \gamma_{\nu\mu} \gamma_{\nu\mu})$ at the heavy quark vertex, the general solutions of equation (4) consist of two series of degenerate doublets of the following form

$$\begin{aligned}\chi_{0,j(-1)^{j+1},1}(p) &= \sqrt{\frac{2j+1}{2}} \frac{1+\hat{v}}{2} \gamma^5 \eta_{\alpha_1 \dots \alpha_j} p_t^{\alpha_1} \dots p_t^{\alpha_j} (p_t^{\alpha_{j+1}} - \frac{j}{2j+1} \gamma^{\alpha_1} \hat{p}_t) (\psi_{1j} - \hat{p}_t \psi_{2j}), \\ \chi_{0,j+1(-1)^{j+1},1}(p) &= \frac{1}{\sqrt{2}} \frac{1+\hat{v}}{2} \eta_{\alpha_1 \dots \alpha_{j+1}} \gamma^{\alpha_1} p_t^{\alpha_2} \dots p_t^{\alpha_{j+1}} (\psi_{1j} - \hat{p}_t \psi_{2j})\end{aligned}\quad (5)$$

and

$$\begin{aligned}\chi_{0,j+1(-1)^j,2}(p) &= \frac{1}{\sqrt{2}} \frac{1+\hat{v}}{2} \gamma^5 \eta_{\alpha_1 \dots \alpha_{j+1}} \gamma^{\alpha_1} p_t^{\alpha_2} \dots p_t^{\alpha_{j+1}} (\psi_{1j} - \hat{p}_t \psi_{2j}), \\ \chi_{0,j(-1)^j,2}(p) &= \sqrt{\frac{2j+1}{2}} \frac{1+\hat{v}}{2} \eta_{\alpha_1 \dots \alpha_j} p_t^{\alpha_2} \dots p_t^{\alpha_j} (p_t^{\alpha_1} - \frac{j}{2j+1} \gamma^{\alpha_1} \hat{p}_t) (\psi_{1j} - \hat{p}_t \psi_{2j}).\end{aligned}\quad (6)$$

They correspond to meson doublets of the spin-parity $[j^{(-1)^{j+1}}, (j+1)^{(-1)^{j+1}}]$ and $[j^{(-1)^j}, (j+1)^{(-1)^j}]$ respectively.

Since we shall confine ourselves to the results in the leading order in this article, we shall not write down the equations for eigen-values and wave functions in the order $O(1/M)$. They can be found in [5,6].

The above results are essentially model independent. For calculations of values of physical quantities we must use a specific model. We shall assume that the light quark propagator is that of the free quark with the mass m , the parameter $Z=1$ and that the kernel is the sum of a static confinement term and a gluon exchange term in the Coulomb gauge in the rest frame of the meson. In order to obtain wave functions of moving mesons we generalize the kernel to the following covariant form

$$\begin{aligned}-iG_+ &= I \otimes I \left[\frac{8\pi\kappa}{(|p_t - q_t|^2 + \mu^2)^2} - (2\pi)^3 \delta^3(p_t - q_t) \right. \\ &\quad \left. \times \int \frac{8\pi\kappa}{[k^2 + \mu^2]^2} \frac{d^3 k}{(2\pi)^3} \right] - I \otimes \hat{v} \frac{16\pi\alpha_{eff}}{3(|p_t - q_t|^2)},\end{aligned}\quad (7)$$

μ is taken to approach zero after solving the equation numerically. The running coupling constant α_{eff} is taken as

$$\begin{aligned}\alpha_{eff}(p_t, q_t) &= \frac{12\pi}{27 \ln \frac{\Lambda_{QCD}^2 (|p_t|^2 + |q_t|^2)}{\Lambda_{QCD}^2}}, \quad \text{when } \alpha_{eff} < 1, \\ \alpha_{eff}(p_t, q_t) &= 1 \quad \text{otherwise.}\end{aligned}\quad (8)$$

G_- contains the propagating transverse gluon term written in the covariant form. It does not enter the equation of leading order.

The kernel G_+ in (7) is independent of p_t and q_t . Defining $\tilde{f}(p) = \int f(p) \frac{d^3 p}{2\pi}$ for arbitrary function $f(p)$, we find from (3) the following 3-dim equation

$$\tilde{\chi}_0(p_t) = \frac{-i}{E_0 + m - W_p} \int G_+ \tilde{\chi}_0 \frac{d^3 q_t}{(2\pi)^3} \frac{W_p \hat{v} - \hat{p}_t - m}{2W_p}. \quad (9)$$

where $W_p = \sqrt{p_t^2 + m^2}$, $p_t^n = |p_t|^n$. Introduce the projection operators

$$\Lambda_{\pm}(-p_t) = W_p \pm \hat{v}(\hat{p}_t + m). \quad (10)$$

Using the property of $\Lambda_{\pm}$, we obtain from (9)

$$\tilde{\phi}_{1j} = -(W_p + m)\tilde{\phi}_{2j}, \quad \tilde{\psi}_{1j} = (W_p - m)\tilde{\psi}_{2j}. \quad (11)$$

Therefore, only one independent component remains in the zero-th order wave function. Equations satisfied by $\tilde{\phi}_{2j}$ and $\tilde{\psi}_{1j}$ can be written as [4,5]

$$\begin{aligned}\tilde{\phi}_{2j} &= \frac{1}{2W_p(E + m - W_p)} \{ -A_-^j [(m + W_q)\tilde{\phi}_{2j}] - (W_p - m)A_+^{j+1}[\tilde{\phi}_{2j}] \}, \\ \tilde{\psi}_{1j} &= \frac{1}{2W_p(E' + m - W_p)} \{ p_t^2 A_-^j [(W_q - m)^{-1}\tilde{\psi}_{1j}] + (W_p - m)A_+^{j-1}[\tilde{\psi}_{1j}] \},\end{aligned}\quad (12)$$

where $A_{\pm}^j[\tilde{f}]$ can be defined as

$$A_{\pm}^j[\tilde{f}] = \frac{1}{|p_l|^j} \int \frac{d^3 q_l}{(2\pi)^3} |q_l|^j P_j(\cos\theta) (\tilde{V}_1 \pm \tilde{V}_2) \tilde{f}(q_l^2), \quad (13)$$

and $\tilde{V}_1, \tilde{V}_2$ are the coefficients of $I \otimes I$ and $I \otimes \hat{v}$ in (7) respectively.

The 4-dim wave function χ_0 can be expressed in term of the 3-dim wave function as the following

$$\chi_0 = \frac{-i}{p_l + E + m + i\epsilon} \left[\frac{E + M - W_p}{p_l + W_p - i\epsilon} \tilde{\chi}_0 + \frac{\Omega}{p_l - W_p + i\epsilon} \right], \quad (14)$$

where Ω is defined as

$$\Omega = \int -iG_+ \tilde{\chi}_0 \frac{d^3 q_l}{(2\pi)^3} \hat{v} \Lambda_+(-p_l), \quad (15)$$

The function Ω for the doublets (5) can be written as

$$\begin{aligned} \Omega_{j(-1)^{j+1},1}(p) &= \sqrt{\frac{2j+1}{2j+2}} \frac{1+\hat{v}}{2} \gamma^5 \eta_{\alpha_1 \dots \alpha_j} p_l^{\alpha_1} \dots p_l^{\alpha_j} (p_l^{\alpha_{j+1}} - \frac{j}{2j+1} \gamma^{\alpha_{j+1}} \hat{p}_l) (\Omega_{1j,1} - \hat{p}_l \Omega_{2j,1}), \\ \Omega_{j+1(-1)^{j+1},1}(p) &= \frac{1}{\sqrt{2}} \frac{1+\hat{v}}{2} \eta_{\alpha_1 \dots \alpha_{j+1}} \gamma^{\alpha_1} p_l^{\alpha_2} \dots p_l^{\alpha_{j+1}} (\Omega_{1j,1} - \hat{p}_l \Omega_{2j,1}) \end{aligned} \quad (16)$$

Using (12), we obtain

$$\begin{aligned} \Omega_{1j,1} &= (-W_p + m) A_-^j[\tilde{\phi}_{1j}] + p_l^2 A_+^{j+1}[\tilde{\phi}_{2j}], \\ \Omega_{2j,1} &= \frac{1}{W_p - m} \Omega_{1j,1}. \end{aligned} \quad (17)$$

Similarly, the function Ω for the doublets (6) are written as

$$\begin{aligned} \Omega_{j+1(-1)^j,2}(p) &= \frac{1}{\sqrt{2}} \frac{1+\hat{v}}{2} \gamma^5 \eta_{\alpha_1 \dots \alpha_{j+1}} \gamma^{\alpha_1} p_l^{\alpha_2} \dots p_l^{\alpha_{j+1}} (\Omega_{1j,2} - \hat{p}_l \Omega_{2j,2}), \\ \Omega_{j(-1)^j,2}(p) &= \sqrt{\frac{2j+1}{2j+2}} \frac{1+\hat{v}}{2} \eta_{\alpha_1 \dots \alpha_j} p_l^{\alpha_1} \dots p_l^{\alpha_j} (p_l^{\alpha_{j+1}} - \frac{j}{2j+1} \gamma^{\alpha_{j+1}} \hat{p}_l) (\Omega_{1j,2} - \hat{p}_l \Omega_{2j,2}). \end{aligned} \quad (18)$$

where

$$\begin{aligned} \Omega_{1j,2} &= p_l^2 A_-^j[\tilde{\psi}_{2j}] + (W_p + m) A_+^{j-1}[\tilde{\psi}_{1j}], \\ \Omega_{2j,2} &= -\frac{1}{W_p + m} \Omega_{1j,2}. \end{aligned} \quad (19)$$

Let us introduce the interaction vertex of the pion to the light quark of the following form

$$\frac{g_q}{f_\pi} \gamma^5 \hat{q} \frac{\tau_i}{\sqrt{2}}, \quad (20)$$

where q is the momentum of the pion. From (20) and (14), the amplitude for the decay process $D^j(p) \rightarrow D(p') + \pi(q)$ of the excited heavy meson state with spin j is given by

$$\begin{aligned} \langle D^0 \pi | D^j \rangle &= \frac{g_q I}{f_\pi} \int (-i)(p_l + E + m) \text{tr}[\bar{\chi}(p')^0 \chi^j(p) \gamma_5 \hat{q}] \frac{d^4 p}{(2\pi)^4} \\ &= \frac{g_q I}{f_\pi} \int \left\{ -\text{tr}[\bar{\chi}^0(p_l') \tilde{\chi}^j(p_l) \gamma_5 \hat{q}] + \frac{1}{q_l - W_p' - W_p} \text{tr}[\bar{\Omega}^0(p_l') \tilde{\chi}^j(p_l) \gamma_5 \hat{q}] \right. \\ &\quad \left. - \frac{1}{q_l + W_p' + W_p} \text{tr}[\bar{\chi}^0(p_l') \Omega^j(p_l) \gamma_5 \hat{q}] \right\} \frac{d^3 p_l}{(2\pi)^3}, \end{aligned} \quad (21)$$

where $I = 1$ for π^+ , $I = 1/\sqrt{2}$ for π^0 . The amplitude for the decay process $D_2^j(p) \rightarrow D(p') + K(q)$ is obtained by replacing f_π by f_K and let $I = 1$ for both K^+ and K^0 in the r.h.s of (21). The formula for the decay processes of B series has the same form.

From (21) we obtain for the decay of 1^+ states D_1 to $D^* \pi$

$$\langle D^*(e) \pi | D_1(\eta) \rangle = \frac{g_q I}{f_\pi} [\eta \cdot q_l e \cdot q_l + \frac{1}{3} q_l^2 e \cdot \eta] \sqrt{\frac{3}{2}} F_1(q) \quad (22)$$

$$\Gamma(D_1 \rightarrow D^* + \pi) = \frac{g_q^2 I^2}{6\pi f_\pi^2} q_l^5 |F_1|^2 \quad (23)$$

where η and e are polarization vectors of D_1 and D^* respectively, and

$$\begin{aligned} F_1(q) &= \int \frac{1}{2q_l^4} \left[-2\phi_1' \phi_{1j} q_l^2 p_l \cdot q_l \right. \\ &\quad - \phi_2' \phi_{2j} [6(p_l \cdot q_l)^3 - 3(p_l \cdot q_l)^2 q_l^2 - 4p_l \cdot q_l p_l^2 q_l^2 + p_l^2 q_l^4] \\ &\quad - q_l \{ \phi_2' \phi_{1j} [3(p_l \cdot q_l)^2 - 2p_l \cdot q_l q_l^2 - p_l^2 q_l^2] + \phi_1' \phi_{2j} [3(p_l \cdot q_l)^2 - p_l^2 q_l^2] \} \\ &\quad \left. + (\chi^0 \rightarrow \Omega^0) \frac{1}{q_l - W_p' - W_p} - (\chi^j \rightarrow \Omega^j) \frac{1}{q_l + W_p' + W_p} \right] \frac{d^3 p_l}{(2\pi)^3} \end{aligned} \quad (24)$$

For the process of 2^+ state $D_2^* \rightarrow D\pi$, one have

$$\langle D\pi | D_2^*(\eta^{\mu\nu}) \rangle = \frac{g_q I}{f_\pi} \eta_{\mu\nu} q_l^\mu q_l^\nu F_1(q), \quad (25)$$

$$\Gamma(D_2 \rightarrow D + \pi) = \frac{g_q^2 I^2}{15\pi f_\pi^2} q_l^5 |F_1|^2, \quad (26)$$

where $\eta_{\mu\nu}$ is the polarization tensor of D_2^* . For the decay process $D_2^* \rightarrow D^* \pi$

$$\langle D^*(e^*) \pi | D_2^*(\eta^{\mu\nu}) \rangle = \frac{g_q I}{f_\pi} i \epsilon_{\alpha\beta\sigma\rho} v^\alpha e^\beta \eta^{\sigma\mu} q_l^\rho q_{l\sigma} F_1(q), \quad (27)$$

$$\Gamma(D2 \rightarrow D^* + \pi) = \frac{g_q^2 I^2}{10\pi f_\pi^2} q_t^5 |F_1|^2, \quad (28)$$

Notice that the amplitudes for the above three processes are determined by the same form-factor $F_1(q)$. We also obtain for the decay $D_1 \rightarrow D(0^+) \pi$

$$\langle D^{0+} \pi | D_1(\eta) \rangle = \frac{g_q I}{f_\pi} \eta \cdot q_t F_2, \quad (29)$$

$$\Gamma(D1 \rightarrow D^{0+} + \pi) = \frac{g_q^2 I^2}{6\pi f_\pi^2} q_t^3 |F_2|^2, \quad (30)$$

where

$$F_2(q) = \frac{1}{3q_t} \int \sqrt{\frac{3}{2}} \left[\psi'_1 \phi_{2j} [3(p_t \cdot q_t)^2 - p_t^2 q_t^2] \right. \\ + \psi'_2 \phi_{1j} [-3(p_t \cdot q_t)^2 + 2p_t \cdot q_t q_t^2 + p_t^2 q_t^2] \\ + q_t \{ \psi'_2 \phi_{2j} [3(p_t \cdot q_t)^2 - 2p_t \cdot q_t p_t^2 - p_t^2 q_t^2] + 2\psi'_1 \phi_{1j} p_t \cdot q_t \} \\ \left. + (\chi^0 \rightarrow \Omega^0) \frac{1}{q_t - W'_p - W_p} - (\chi^j \rightarrow \Omega^j) \frac{1}{q_t + W'_p + W_p} \right] \frac{d^3 p_t}{(2\pi)^3}, \quad (31)$$

For numerical calculations, we use the following values of the parameters in B-S equation determined in [5] from fitting the experimental data for the masses of observed heavy meson states²

$$\begin{aligned} M_b &= 5.024 \text{ GeV}, & M_c &= 1.58 \text{ GeV}, & m_\mu &= 0.35 \text{ GeV}, \\ m_s &= 0.465 \text{ GeV}, & \Lambda_{QCD} &= 0.38 \text{ GeV}, & \kappa &= 0.2 \text{ GeV}^2. \end{aligned} \quad (32)$$

and value of g_q

$$g_q^2 = 0.16 \quad (33)$$

determined in [6]³ from the branching ratio $\text{Br}(D^{*0} \rightarrow D^0 \pi^0) / \text{Br}(D^{*0} \rightarrow D^0 \gamma)$. Numerical results for the widths of various decay modes are listed in the table together with available experimental data [8-10]. The masses of B_1 , B_2^* , B_{s1} and B_{s2}^* used for determining the momentum q_t are taken from the theoretical results obtained in [6]. The results for the decay widths of these particles are sensitive to their mass values. We give in the table the corrections to the widths for $\pm 10 \text{ MeV}$ corrections to the masses of these particles.

The widths for the processes emitting a ρ instead of π were calculated in [7]. The processes with a large branching ratio in their results are $D_1 \rightarrow D\rho$, $D^* \rho$ which has the widths 6 MeV and $D_2^* \rightarrow D\rho$, $D^* \rho$ which has 3 MeV. These facts should be kept in the mind when one compares our results with experimental data on the total width of D_1 and D_2^* .

²the value of m_s was misprinted in [5]

³The value of g_q^2 listed in the table of [6] has an error of factor 1/2

The values (33) for g_q was obtained in [6] under the assumption that the anomalous magnetic moment μ_q of the quark is negligible. This assumption is usually made in the constituent quark model and is supported by the result from the approach of chiral quark Lagrangian[11]. For non-vanishing μ_q the values of g_q^2 will be larger. A more accurate measurement of the total width of D^* or the branching ratio for the process $D^{*+} \rightarrow D^+ + \gamma$ will be given more information on g_q^2 . The theoretical values in the table can be considered as $g_q^2/0.16$ times the decay width. We have neglected SU(3) breaking for g_q in the calculations. In non-relativistic quark model, the experimental data for g_A/g_V for β decay of neutron and Λ decay imply that the SU(3) breaking correction for g_q in $\Delta S = 1$ transition is small and negative ($\sim -5\%$). Assuming that this qualitative behavior does not change in the relativistic model, this approximation together with replacing f_π by f_K may be an improvement to the SU(3) symmetry approximation for the emission of Goldstone mesons.

In conclusion, we have obtained general formulas for the widths of the decay processes of excited heavy meson states with arbitrary spin-parity to the ground state heavy meson plus a pion or kaon in the leading order of $1/M$ expansion in the relativistic B-S equation developed previously. Numerical results for various decay widths of 1^+ and 2^+ states are obtained with a model similar to that used successfully for quarkonium systems and giving results for heavy meson spectra in good agreement with experimental data. The calculated values of widths are in satisfactory agreement with experimental data when they are available. No new parameter has been introduced in these calculations. Corrections in the next to leading order in $1/M$ expansion are expected to be small at least for B and B_s series. Though tedious, the calculations for these corrections can be carried out. They will be done in a later work.

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Table-4 The decay widths of low-lying states of D and B systems(Mev)

Decay Process	calculation	CLEO[9]	PDG[8]	E687[10]
$\Gamma(D1 \rightarrow D^* + \pi)$	8.0			
$\Gamma(D1 \rightarrow D(0^+) + \pi)$	0.5			
$\Gamma(D1 \rightarrow all)$		20^{+9}_{-8}	20^{+9}_{-5}	15^{+8+5}_{-8-5}
$\Gamma(D2 \rightarrow D + \pi)$	15.			
$\Gamma(D2 \rightarrow D^* + \pi)$	6.8			
$D2 \rightarrow D\pi/D^*\pi$	2.2	$2.1 \pm 0.6 \pm 0.6$	2.4 ± 0.7	
$\Gamma(D2 \rightarrow all)$		28^{+14}_{-13}	19^{+7}_{-7}	24^{+7+5}_{-7-5}
$\Gamma(Ds2 \rightarrow D + K)$	7.5			
$\Gamma(Ds2 \rightarrow D^* + K)$	0.67			
$\Gamma(Ds2 \rightarrow all)$			16^{+5+3}_{-4-3}	
$\Gamma(Ds1 \rightarrow D^* + K)$	0.16			
$\Gamma(Ds1 \rightarrow all)$		< 2.3	< 4.6	< 3.2
$\Gamma(B2 \rightarrow B + \pi)$	$6.0^{+0.5}_{-0.5}$			
$\Gamma(B2 \rightarrow B^* + \pi)$	$5.5^{+0.6}_{-0.6}$			
$\Gamma(B1 \rightarrow B^* + \pi)$	$7.9^{+0.5}_{-0.9}$			
$\Gamma(Bs2 \rightarrow B + K)$	1.25			
$\Gamma(Bs2 \rightarrow B^* + K)$	$0.4^{+0.2}_{-0.2}$			
$\Gamma(Bs1 \rightarrow B^* + K)$	$0.3^{+0.2}_{-0.2}$			

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