Inclusive decays of heavy quark hybrids

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Abstract

In order to understand the nature of the XYZ particles, theoretical predictions of the various XYZ decay modes are essential. In this work, we focus on the semi-inclusive decays of heavy quarkonium hybrids into traditional quarkonium in the EFT framework. We begin with weakly coupled potential NRQCD effective theory that describes systems with two heavy quarks and incorporates multipole expansions and use it to develop a Born-Oppenheimer effective theory (BOEFT) to describe the hybrids and compute the semi-inclusive decay rates. We compute both the spin-conserving and spin-flipping decay rates and find that our numerical results of the decay rates are different from the previous studies. We also develop a systematic framework in which the theoretical uncertainty can be systematically improved.

1 Introduction

The Standard Model (SM) describes hadrons as bound states of quarks and gluons bounded by the strong interactions. Traditionally, the hadrons were classified as mesons that are bound state of quark-antiquark pair or baryons that are bound state of 3-quarks using the quark model. However, the underlying theory of strong interactions, Quantum Chromodynamics (QCD) also allows for existence of complex hadron structures beyond mesons and baryons such as tetraquark (4-quark states), pentaquark (5-quark states), hybrids (hadrons with active gluons) and glueballs (bound state of gluons), which are known as exotic hadrons or exotics. The so called XYZ states are the exotic hadrons in the heavy-quark sector. The XYZ states do not fit the usual charmonium (c\bar{c}) or bottomonium (b\bar{b}) spectrum and in some cases have exotic quantum numbers which cannot be reproduced by the ordinary hadrons such as charged Zc and Zb states. In 2003, the Belle experimental collaboration observed the first exotic state X(3872) [1] and since then, several of the new exotic hadrons in the heavy-quark sector have

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been observed by the different experimental groups: B-factories (BaBar, Belle and CLEO), τ-charm facilities (CLEO-c, BESIII) and also proton-(anti)proton colliders (CDF, D0, LHCb, ATLAS, CMS) (see the reviews [2, 3] for details on experimental observation).

There are several theoretical models and proposals to understand the nature of the XYZ exotics. One viable and attractive interpretation for at-least some of the XYZ mesons is the quarkonium hybrids, which is a bound state of heavy quark and a heavy antiquark together with gluonic excitation. The other proposals are hadroquarkonium, heavy meson molecule, tetraquark, and diquark-diquark model (see Ref. [2, 3] for review). However, no single proposal can theoretically explain the complete spectrum of the XYZ exotic states. On the other hand, several new exotic states have been observed in experiments for which the masses and the decay rates has been measured (see Ref. [4]). Specifically, several of these exotic states has been discovered from their decays to standard quarkonium. Therefore, a theoretical understanding of the decays of XYZ exotics might be an another avenue for understanding their structure. In this work, our objective is to study the inclusive decays of heavy quark hybrids to traditional quarkonium i.e, \( H^m \rightarrow \bar{Q} \Sigma + X \), where \( H^m \) is a low-lying hybrid, \( Q^\Sigma \) is a low-lying quarkonium state and \( X \) denotes other final state particles.

Within the QCD framework, one can use lattice simulations and effective field theories (EFTs) to describe the traditional quarkonium and quarkonium hybrids and compute its spectra. Since, the heavy quarks in quarkonium and heavy-quark hybrids are nonrelativistic, the appropriate framework to use is the nonrelativistic effective theory NRQCD [5,6]. More specifically, if we are only interested in the dynamics of the two heavy quarks, then the appropriate framework to use is the potential NRQCD effective theory known as pNRQCD [7, 8]. In case of quarkonium hybrids, there are well-separated energy scales: \( m_Q \) (mass of heavy quark) >> \( m_Q v \) (relative momentum scale) >> \( \Lambda_{QCD} \) (energy scale for gluonic excitations) >> \( m_Q v^2 \) (dynamics of two heavy quark). The above momentum hierarchy suggests of an energy gap between the gluonic excitations and the excitations of the heavy quark-antiquark pair that has also been confirmed by the lattice data [9, 10]. This justifies the use of effective theory based on Born-Oppenheimer approximation (BOEFT) to describe the hybrids [11–15]. On the other hand, the lattice inputs are essential for determining the static potentials that are used for solving the Schrödinger equation for computing the spectra. Traditionally, the lattice studies of the heavy quark hybrids have mainly focused in the charmonium sector and recently in bottomonium sector [16]. In the charm sector, the recent lattice studies have predicted the existence of a lowest hybrid spin-multiplet \( J^{PC} = [(0, 1), 1^{+}, 1^{--}] \) at about 4.3 GeV [17–20]. In the bottom sector, the lattice study in Ref. [16], predicted hybrid states with quantum numbers \( J^{PC} = [(0, 1), 2^{++}, 1^{--}] \) approximately 1500 MeV above the ground-state \( \eta_b \) meson.

In this work, we will use the BOEFT for the hybrids and pNRQCD for the low-lying quarkonium states. For computing the decay rates, we perform a matching calculation between BOEFT and pNRQCD to obtain the imaginary terms in the BOEFT potential. In Sec. 2, we compute the quarkonium and the hybrid spectrum, in Sec. 3, we perform the matching calculation and compute the decay rates and we conclude in Sec. 4.

# 2 Spectrum

## 2.1 Quarkonium

The conventional quarkonium states (\( Q\bar{Q} \)) are color singlet bound states of a heavy quark and antiquark in the ground state static potential \( V_{\Sigma}^{\bar{Q}Q}(r) \). The shape of the static potential \( V_{\Sigma}^{\bar{Q}Q}(r) \) is well described by the Cornell potential. The Schrödinger equation describing the quarkonium
spectrum is given by
\[
\left(-\nabla^2 + V_{\Sigma^g}(r)\right) \Phi^Q_{(n)}(r) = E^Q_{n} \Phi^Q_{(n)}(r),
\]
where \(m_Q\) is the heavy quark mass, \(E^Q_{n}\) is the quarkonium energy, \(\Phi^Q_{(n)}(r)\) denotes the quarkonium wave-function, the index \((n) = (n, j, l, s)\) denotes the usual set of quarkonium quantum numbers. We use the following form of the static potential
\[
V_{\Sigma^g}(r) = -\frac{\kappa_g}{r} + \sigma_g r + E^{QQ}_g,
\]
where \(\kappa_g = 0.489\) and the string tension parameter \(\sigma_g = 0.187\text{GeV}^2\) are determined from the fit to the lattice data. The constant \(E^{\bar{Q}Q}_g\) is different for both charmonium and bottomonium and is determined by comparison to the experimental spin-averaged mass from PDG 2020 \[4\]
\[
E^{\bar{c}c}_g = -0.254\text{GeV}, \quad E^{bb}_g = -0.195\text{GeV},
\]
where have used the RS-scheme charm and bottom mass: \(m_c = 1.477\text{GeV} and m_b = 4.863\text{GeV}\) to compute the quarkonium spectrum. The quarkonium mass is given by \(M_{QQ} = 2m_Q + E^Q_{n}\) for \(Q = (c, b)\), where \(E^Q_{n}\) is the eigenvalue in Eq. (1).

2.2 Hybrids

Hybrids \((Q\bar{Q}g)\) are exotic hadrons that are color singlet bound states of a color octet \(Q\bar{Q}\) source coupled to gluonic excitations. Therefore, hybrid states are more complicated compared to traditional quarkonium due to presence of active gluons. The energy scale for the gluonic excitations is the nonperturbative energy scale \(\Lambda_{QCD}\). In the BOEFT description, the nonperturbative gluon dynamics generate a background static potential in which the heavy quark-antiquark pair in the hybrids binds together. In the static limit \((m_Q \rightarrow \infty)\), the hybrid spectrum is composed of the static energies, which are characterized by the representation of the \(D_{ooch}\) cylindrical symmetry group just like in diatomic molecules. The hybrid static energies are nonperturbative quantities that are generally computed on the lattice. In the short-distance limit \(r \rightarrow 0\), where \(r\) is the relative coordinate of \(Q\bar{Q}\), the hybrid static energies are degenerate and quantum numbers are characterized by representations of spherical symmetry group \(O(3) \times C\) instead of \(D_{ooch} [7,13,14]\). We focus here on the low-lying hybrids coming from \(\Sigma_u^-\) and \(\Pi_u\) static potentials and we closely follow the notations in Ref. [13].

The BOEFT Lagrangian that describes the hybrid states is given by
\[
L_{BOEFT} = \int_{\mathcal{R}} \int_{\mathcal{R}} \sum_{\kappa, \lambda} \sum_{\lambda', \lambda''} \left[ \Psi_{\kappa, \lambda}(r, R, t) \left( i \partial_t - V_{\kappa, \lambda'}(r) + p^{i}_{\kappa, \lambda} \frac{\nabla^2}{m_Q} p^{i}_{\kappa, \lambda'} \right) \Psi_{\kappa, \lambda'}(r, R, t) 
+ \Psi_{\kappa, \lambda}(r) \Delta V(r) \delta_{\lambda, \lambda'} \Psi_{\kappa, \lambda'}(r, R, t) \right] + \cdots,
\]
where \(\mathcal{R} \equiv \int d^3r, r\) and \(R\) are the relative and center-of-mass coordinates of the heavy-quark-antiquark pair, the quantum number \(\kappa\) is \(\kappa \equiv K^{PC}\), with \(K\) defined as the angular momentum of the gluonic degrees of freedom, \(\Psi_{\kappa, \lambda}\) denotes the hybrid field (or the wave-function) and the ellipses represent higher order terms in the multipole expansion. \(p^{i}_{\kappa, \lambda}\) (i denotes vector index) is the projection operator that projects the gluonic degrees of freedom to an eigenstate of \(K \cdot \hat{r}\) with eigenvalue \(\lambda\). These projection operators correctly reproduce the hybrid quantum numbers in \(D_{ooch}\) representation. For low-lying hybrids coming from \(\Sigma_u^-\) and \(\Pi_u\) static potentials, the gluon quantum number \(\kappa = 1^{+-}\). Therefore, for our purpose, the projectors \(P^{i}_{1, \lambda}\) are given
by $P_{10}^i = \hat{P}_0^i(\theta, \phi) = \hat{P}_i^i$ for projecting onto the $\Sigma_u^-$ state and $P_{\pm 1}^i = \hat{P}_{\pm 1}^i(\theta, \phi) = (\hat{\theta}^i \pm i \hat{\phi}^i)/\sqrt{2}$ for projecting onto the two components of the $\Pi_u^-$ state, where $\hat{\theta}$ and $\hat{\phi}$ are the usual spherical unit vectors. In Eq. (4), the BOEFT potential $V_{\kappa\lambda\lambda'}(r)$ can be expanded in $1/m_q$ as

$$V_{\kappa\lambda\lambda'}(r) = E_{\kappa\lambda}(r)\delta_{\lambda\lambda'} + \frac{V_{\kappa\lambda\lambda'}^{(1)}(r)}{m_q} + \ldots,$$

(5)

where $E_{\kappa\lambda}(r)$ denotes the static potential, and $V_{\kappa\lambda\lambda'}^{(1)}(r)$ can be written as sum of spin-dependent and spin-independent pieces [15]. The effective potential $\Delta V$ in Eq. (4) (that is treated as a perturbation) is responsible for producing transitions to standard quarkonium states and the form will be determined by performing a matching calculation in section 3. From now on, we ignore the subscript $\kappa$ (as $\kappa = 1^+$ for our purpose) and we write the hybrid wave-function as

$$\Psi_{\lambda}^{(m)}(r) \equiv \Psi_{\lambda}^{(m,j,l,s)}(r) = \psi_{m}^{(j)}(r)\phi_{(l,s)}^{(\lambda)}(\theta, \phi),$$

(6)

where $m$ is the principal quantum number, the quantum number $j$ is the eigenvalue of the total angular momentum operator: $J = L + S$, where $S$ is the $Q\bar{Q}$ spin, $L = L_{Q\bar{Q}} + K$, where $K$ is the gluon angular momentum, and $L_{Q\bar{Q}}$ is the angular momentum operator of the two heavy quarks. We use the notation $m(L_{Q\bar{Q}})_L$ to denote the hybrid state.

At leading order, the equations of motion for the fields $\Psi_{\lambda}^{(m)}(r)$ that follow from Eq. (4) are the set of coupled Schrödinger equations which are given by

$$\sum_{\lambda = \pm 1} \hat{r}_{\lambda}^2(\theta, \phi) \left(-\frac{\nabla^2}{m_q} + E_{\kappa\lambda}(r)\right) \hat{r}_{\lambda}(\theta, \phi) \Psi_{\lambda}^{(m)}(r) = \mathcal{E}_m \Psi_{\lambda}^{(m)}(r),$$

(7)

where $\mathcal{E}_m$ is the eigenvalue. The hybrid mass is given by $M_{Q\bar{Q}} = 2m_q + \mathcal{E}_m$ for $Q = (c, b)$. Since, there are projection operators on both side of $\nabla^2_r$ in Eq. (7), the contributions from $\Sigma_u^-$ and $\Pi_u^-$ potentials mix together that results in pairs of solutions with same angular momentum quantum number but opposite parity [13]. The static potentials that we use for computing the hybrid spectrum is split into a short-distance part and long-distance part [13]:

$$E_{\kappa\lambda}^{(0)}(r) = \begin{cases} V_{o}^{RS}(v_f) + \Lambda_{RS}^{Q\bar{Q}}(v_f) + b_{\lambda}r^2, & r < 0.25\text{fm} \\ \mathcal{V}(r), & r > 0.25\text{fm} \end{cases},$$

(8)

where for the short-distance part ($r < 0.25\text{fm}$) we have used the RS-scheme octet potential $V_{o}^{RS}(r)$ up to order $\alpha_s^3$ in perturbation theory and the RS-scheme gluelump mass $\Lambda_{RS}^{Q\bar{Q}} = 0.87(15)$ GeV at the renormalon subtraction scale $v_f = 1\text{ GeV}$ [10,21]. The RS-scheme octet potential is given by [10,21]

$$V_{o}^{RS}(r, v_f, \mu) = V_{o}(r, \mu) - \delta V_{o}^{RS}(v_f),$$

(9)

with

$$V_{o}(r, \mu) = \left(\frac{C_A}{2} - C_F\right)\frac{\alpha_s(\mu)}{\mu},$$

(10)

$$\delta V_{o}^{RS}(v_f) = \sum_{n=1}^{\infty} N_{v_f}v_f \left(\frac{\beta_0}{2\pi}\right)^n \alpha_s^{n+1}(v_f) \sum_{k=0}^{\infty} c_k \frac{\Gamma(n+1+b-k)}{\Gamma(1+b-k)},$$

(11)

where $\mu$ denotes the energy scale parameter, $N_{v_f} = 0.114001$, the parameters $b$ and $c_k$ are defined in Ref. [10]. The form of $\alpha_s$ up to order $\alpha_s^3$ in perturbation theory is given in Ref. [22].

The long-distance ($r > 0.25\text{fm}$) part of the potential $\mathcal{V}(r)$ is given by

$$\mathcal{V}(r) = \frac{a_1}{r} + \sqrt{a_2r^2 + a_3 + a_4}.$$
The above form of the long-distance potential $V(r)$ is chosen so as to reproduce the short and long distance behavior of the Cornell potential. The parameters $b_\lambda$ in Eq. (8) and $a_1$, $a_2$, $a_3$ and $a_4$ in Eq. (12) are different for both $\Sigma$ and $\Pi$ static potentials. The parameters are determined by performing a fit to the lattice data in Refs. [9,10] and demanding that the short-range and the long-range pieces in Eq. (8) are continuous up to first derivatives (see Ref. [13] for details). The result for the spectrum is given in Table III of Ref. [13].

3 Inclusive Decay Rate

We want to compute the inclusive decay rate of low-lying quarkonium hybrids decaying to traditional quarkonium i.e., $H_m \rightarrow Q_n + X$, where $H_m$ is a low-lying hybrid, $Q_n$ is a low-lying quarkonium state and $X$ denotes other final state particles. We denote the energy (mass) difference by $\Delta E = \varepsilon_m - E^Q_n \gtrsim 1$ GeV, which for low-lying hybrid and quarkonium states satisfy the following hierarchy of energy scales: $m_Q \nu >> \Delta E >> \Lambda_{QCD} >> m_Q \nu^2$. This implies that the relevant theory at the energy scale $\Delta E$ is the weakly coupled pNRQCD effective theory which is obtained from NRQCD by integrating out gluons with momentum and energy of order $\sim m_Q \nu$ and quarks with energy of order $\sim m_Q \nu$. In order to describe the dynamics of two heavy quarks in the hybrids that happens at energy scale $m_Q \nu^2$, we will use the Born-Oppenheimer effective theory (BOEFT). Hence, starting with pNRQCD effective theory, we integrate out gluons with 4–momentum of order $\sim \Delta E$ and $\sim \Lambda_{QCD}$ in loops and match it to BOEFT that describes system at energy scale $m_Q \nu^2$. This matching condition leads to imaginary terms in the BOEFT potential which is related to the hybrid decay rate by the optical theorem.

The pNRQCD Lagrangian up to next-to-leading-order (NLO) in multipole expansion or in $1/m_Q$ is given by

$$L_{pNRQCD} = \int d^4 R \int d^3 r \left[ \left( i \partial_0 - h_s \right) S + \left( i D_0 - h_o \right) O \right] + g \text{Tr} \left[ S^i r \cdot E O + O^i r \cdot E S \right]$$

$$+ \frac{g c}{m_Q} \text{Tr} \left[ S^i (S_1 - S_2) \cdot B O + O^i (S_1 - S_2) \cdot B S \right] + \cdots,$$

(13)

where $\int_k \equiv \int d^3 R$, $S$ and $O$ denotes the singlet and the octet fields and ellipsis represents higher order terms as well as terms including light quarks and gluons. The singlet and octet Hamiltonian densities are given by

$$h_s = -\nabla^2 / m_Q + V_s(r), \quad h_o = -\nabla^2 / m_Q + V_o(r),$$

(14)

where $V_s(r)$ and $V_o(r)$ are the perturbative singlet and octet potentials. For our purpose, we use the following form for $V_s(r)$ and $V_o(r)$

$$V_s(r) = -\frac{\kappa_s}{r} + E_s^Q, \quad V_o(r) = V_o^{RS}(r),$$

(15)

where $V_o^{RS}(r)$ is given by Eq. (9), $\kappa_s = 0.489$, and the constant $E_s^Q$ for $Q = (c, b)$ is chosen so as to reproduce the spin-averaged $1s$ charmonium and $1s$ bottomonium mass.

The $r \cdot E$ vertex in Eq. (13) is responsible for the spin-conserving decays of hybrid whereas the $S \cdot B/m_Q$ vertex is responsible for the spin-flipping decays of hybrid (spin-0 hybrid decaying to spin-1 quarkonium and vice versa). Therefore, the spin-flipping decays are suppressed by powers of the heavy quark mass due to heavy quark spin symmetry.

Beginning with pNRQCD, we integrate out gluons with 4–momentum $\sim \Delta E$ and $\sim \Lambda_{QCD}$ in two steps, and obtain the BOEFT theory that describes system at energy scale $m_Q \nu^2$. We perform this by implementing the matching condition wherein we compute the two-point Green's
function in both the theories and equate them. For spin-preserving decays of hybrid to quarkonium, the two-point function in pNRQCD is expanded up to $O(r^2)$ in the multipole expansion using the NLO pNRQCD Lagrangian in Eq. (13) which is equated to the corresponding two-point function in BOEFT computed using the Lagrangian in Eq. (4). For the spin-flipping decay of hybrids, the two-point function is expanded up to $O\left(1/m^2_Q\right)$ using the pNRQCD Lagrangian in Eq. (13). After implementing this matching condition, we obtain the the following form of the effective potential $\Delta V$ in Eq. (4) for the spin-preserving decays

$$\Delta V = -\frac{g^2}{3} \frac{T_F}{N_c} \int_0^\infty dt \ e^{iA} e^{-ih_{t}/2} e^{-i\theta_{t}} e^{-i\theta_{h}/2} \int \frac{d^3k}{(2\pi)^3} |k| e^{-ik|t|}.$$  (16)

In case of spin-flipping decays, the $s\cdot E$ vertex is replaced by the $S\cdot B/m$ vertex. Thus, using the optical theorem, the (spin-conserving) decay rate of the hybrids for the process $H_n \to Q_n + X$ is given by $\Gamma_{H_n\to Q_n} = -2\langle H_m | \text{Im} \Delta V | H_m \rangle$ (see details of the calculation in Ref. [23])

$$\Gamma(m \to n) = \sum_{n'} \left| \int d^3r \Phi^{\dagger}_{(n')} (r) \Phi^Q_{(n)} (r) \right|^2 \Gamma_{mn'},$$  (17)

where in the above expression we have included the overlap between the quarkonium $\Phi^Q_{(n)}$ and the Coulomb singlet $\Phi^s_{(n')}$ wave-functions (we use compact notation $(n)$ and $(n')$ to denote the set of quantum numbers for quarkonium and singlet state) and $\Gamma_{mn'}$ is given by

$$\Gamma_{mn'} = \text{Re} \frac{8\pi a_s T_F}{3N_c} \sum_{r} \int_0^\infty \left[ \int d^3l \prod_{i=1}^3 \int d^3l' \prod_{i=1}^3 \right] \left[ \Phi^{\dagger}_{(n')} (r') \Phi^Q_{(n)} (r) \right] \left[ \Phi^{\dagger}_{(n')} (r') \Phi^s_{(n)} (r) \right] \left( \Lambda_{\text{glue}} + E_{s}^o / 2 + E_{s}^o / 2 - E_{s}^o \right),$$  (18)

where $\int_r \equiv \int d^3r$, $a_s$ is evaluated at the scale $\Delta E = E_{m} - E_{n}^Q$, $\psi^{\dagger}_{(m)}$ is the hybrid wave-function given in Eq. (6) (i is the vector index, $(m)$ denotes the set of quantum numbers), $\Phi_i^s$ is the octet wave-function, $E_i^o$ is the octet energy, $E_i^s$ is the singlet energy, and $\Lambda_{\text{glue}} = 0.87(15)$ GeV in RS-scheme. In order to obtain the second line from the first line in Eq. (18), we expand the singlet $(h_s)$ and octet $(h_o)$ Hamiltonians in terms of their eigenfunctions $\Phi^s_{(n')}$ and $\Phi^o_{(n')}$, which satisfy

$$h_s (r, p) \Phi^{s}_{(n')} (r) = E_{n}^s \Phi^{s}_{(n')} (r), \quad h_o (r, p) \Phi^{o}_{(n')} (r) = E_{n}^o \Phi^{o}_{(n')} (r).$$  (19)

For computing the decay rates using Eq. (17), we need the octet and the singlet wave-functions. We use the singlet and octet potentials in Eq. (15) to compute $\Phi^s_{(n')}$ and $\Phi^o_{(n')}$. Suppose, we assume: the singlet and quarkonium energies satisfy $E_{n}^s \approx E_{n}^o$, overlap between hybrid and octet wave-functions $\int d^3r \psi^{\dagger}_{(m)} (r) \Phi^s_{(n')} (r)$ is nonzero only for hybrid energy $E_{m}$: $E_{m} = E_{s}^o + \Lambda_{\text{glue}}$ (where ignoring the $b_2 r^2$ term in Eq. (8)) such that the cubic factor within the integrand in Eq. (18) is replaced by a constant $\Delta E$, and the overlap function of quarkonium and singlet wave-function satisfy $\int d^3r \Phi^{\dagger}_{(n')} (r) \Phi^Q_{(n)} (r) \approx 1$ for $(n') = (n)$, then Eq. (17) is simplified to

$$\Gamma^{\text{sim}} (m \to n) \approx \frac{4\pi a_s (\Delta E) T_F}{3N_c} \left[ \int_r \psi^{\dagger}_{(m)} (r) \right] \left[ \int_r \Psi^{Q}_{(n)} (r) \right] \Delta E^3.$$  (20)
The overlap between hybrid and octet wave-function is from changing the scale to $2\Delta E/2$ in $\alpha_s$, while the lower error bar is from changing the scale to $2\Delta E$ in $\alpha_s$.

The simplified decay rate in Eq. (20) is identical to Eq. (17) in Ref. [12] and Eq. (62) in Ref. [24]. However, in Ref. [12], the authors only consider the diagonal elements (where they contract the index $i$ and $j$ in Eq. (20)) instead of the full tensor structure. This led to a selection rule that hybrids with $L = L_{Q\bar{Q}}$ does not decay. This is incorrect as such decays are allowed by considering the tensor structure of the matrix element in Eq. (20). The results for the spin-conserving and spin-flipping decay rates are shown in tables 1 and 2. The spin-flipping decay rates in table 2 are suppressed by the two-powers of the heavy quark mass $m$. In both tables 1 and 2, we see that for most of the cases, the values of the hybrid decay rate from the general expression involving overlap functions of octet and singlet state in Eq. (17) differs from that obtained from the simplified expression in Eq. (20) even after considering the error bars. We find that this difference is mainly due to contributions from the cubic factor within the integrand and the Coulomb singlet wave-functions in Eq. (18). Therefore, this raises the questions on the validity of the approximations that were used to obtain the simplified expression in Eq. (20). Specifically, we find that the approximation about the singlet and quarkonium energy $E_n^2 \approx E_n^s$ is only valid for $1s$ charmonium and bottomonium states and the overlap between hybrid and octet wave-function $\int d^3r \bar{\psi}^{1s}_{(m)}(r)\Phi_i^n(r)$ is nonzero over wide
Table 2: Preliminary results for the spin-flipping inclusive decay rate for hybrids decays to traditional quarkonium states: $H_m \rightarrow Q_n +X$ due to due to $S \cdot B/m_Q$ vertex in Eq. (13). The decay rate in fourth and fifth column is computed using Eq. (20) and in last two column using Eq. (17). The notation $(1 \rightarrow 0)$ denotes spin-1 hybrid decaying into spin-0 quarkonium while $(0 \rightarrow 1)$ denotes spin-0 hybrid decaying into spin-1 quarkonium.

| $m\{q_0\} \rightarrow N L$ | $\Delta E$ (GeV) | $\alpha_s(\Delta E)$ | $\Gamma^{\text{birth}}$(MeV) $(1 \rightarrow 0)$ | $\Gamma^{\text{end}}$(MeV) $(0 \rightarrow 1)$ | $\Gamma$(MeV) $(1 \rightarrow 0)$ | $\Gamma$(MeV) $(0 \rightarrow 1)$ |
|--------------------------|-----------------|---------------------|--------------------------|--------------------------|-----------------|-----------------|
| Charmonium hybrid decay  |                 |                     |                          |                          |                 |                 |
| $1p_0 \rightarrow 1p$    | 1.096           | 0.347               | 45.54 $^{+12.90}_{-11.37}$ | 136.62 $^{+71.69}_{-34.10}$ | 0.09 $^{+0.04}_{-0.02}$ | 0.22 $^{+0.11}_{-0.05}$ |
| $2p_0 \rightarrow 1p$    | 1.560           | 0.295               | 1.66 $^{+0.35}_{-0.33}$ | 4.98 $^{+1.06}_{-1.04}$ | 0.06 $^{+0.02}_{-0.01}$ | 0.18 $^{+0.07}_{-0.04}$ |
| $2p_0 \rightarrow 2p$    | 1.087           | 0.348               | 44.17 $^{+23.33}_{-11.07}$ | 132.51 $^{+69.98}_{-31.20}$ | 0.14 $^{+0.07}_{-0.04}$ | 0.43 $^{+0.22}_{-0.11}$ |
| $3p_0 \rightarrow 1p$    | 1.764           | 0.270               | 0.73 $^{+0.26}_{-0.14}$ | 2.18 $^{+0.77}_{-0.40}$ | 0.09 $^{+0.02}_{-0.01}$ | 0.21 $^{+0.07}_{-0.04}$ |
| $2p_1 \rightarrow 1p$    | 1.173           | 0.335               | 5.09 $^{+2.54}_{-1.23}$ | 15.26 $^{+7.81}_{-3.69}$ | 0.07 $^{+0.04}_{-0.02}$ | 0.21 $^{+0.11}_{-0.05}$ |
| $3p_1 \rightarrow 1p$    | 1.542           | 0.296               | 2.05 $^{+0.86}_{-0.44}$ | 6.16 $^{+2.57}_{-1.32}$ | 0.07 $^{+0.03}_{-0.02}$ | 0.22 $^{+0.09}_{-0.05}$ |
| $3p_1 \rightarrow 2p$    | 1.068           | 0.351               | 3.71 $^{+1.99}_{-0.94}$ | 11.13 $^{+2.81}_{-1.96}$ | 0.18 $^{+0.09}_{-0.04}$ | 0.53 $^{+0.29}_{-0.13}$ |
| Bottonium hybrid decay   |                 |                     |                          |                          |                 |                 |
| $1s(d) \rightarrow 1s$  | 1.087           | 0.348               | 34.53 $^{+18.23}_{-6.65}$ | 103.60 $^{+54.70}_{-25.95}$ | 11.37 $^{+7.09}_{-3.49}$ | 34.11 $^{+18.00}_{-8.94}$ |
| $2s(d) \rightarrow 1s$  | 1.439           | 0.305               | 15.45 $^{+8.04}_{-3.42}$ | 46.35 $^{+20.16}_{-10.26}$ | 0.18 $^{+0.09}_{-0.04}$ | 0.53 $^{+0.29}_{-0.13}$ |
| $3s(d) \rightarrow 1s$  | 1.744           | 0.282               | 0.20 $^{+0.06}_{-0.04}$ | 0.59 $^{+0.03}_{-0.02}$ | 0.51 $^{+0.20}_{-0.11}$ | 1.53 $^{+0.80}_{-0.32}$ |

range of octet energies, if we don’t assume $\mathcal{E}_m \approx E_1^L + \Lambda_{\text{glue}}$ (see Ref. [23]).

4 Conclusions

In this work, we study the inclusive decays of heavy quark hybrids to traditional quarkonium by using the framework of Born-Oppenheimer effective field theory. We have derived an expression of the decay rate given in Eq. (17) that depends on the overlap functions of hybrid, octet, Coulomb singlet and quarkonium wave-functions. The values of the spin-conserving and the spin-flipping decay rates of hybrids are shown in tables 1 and 2. We also find that using certain assumptions, our expression for the decay rate in Eq. (17) reduces to a simplified expression given by Eq. (20) that was earlier derived in Refs. [12, 24]. However, the difference
in the values of the decay rate from Eqs. (17) and (20) shown in tables 1 and 2 raises questions on the validity of those approximations.

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