Z_b/Z'_b \rightarrow \Upsilon \pi \text{ and } h_b \pi \text{ decays in intermediate meson loops model}

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With the recent measurement of Z_b(10610) and Z_b(10650) → B\bar{B}^* + c.c. and B^*\bar{B}^*, we investigate the transitions from the Z_b(10610) and Z_b(10650) to bottomonium states with emission of a pion via intermediate B B^* meson loops. The experimental data can be reproduced in this approach with a commonly accepted range of values for the form factor cutoff parameter α. The Y(3S)π decay channels appear to experience obvious threshold effects which can be understood by the property of the loop integrals. By investigating the α-dependence of partial decay widths and ratios between different decay channels, we show that the intermediate B B^* meson loops are crucial for driving the transitions of Z_b/Z'_b → Y(nS)π with n = 1, 2, 3, and h_b(mP)π with m = 1 and 2.

I. INTRODUCTION

Recently, two charged bottomonium-like structures Z^\pm_b(10610) and Z^\pm_b(10650) (abbreviated to Z^\pm_b and Z^\pm_b in the following) were observed by the Belle Collaboration in the \pi^\pm \Upsilon(nS) (n = 1, 2, 3) and \pi^\pm h_b(mP) (m = 1, 2) invariant mass spectra of \Upsilon(5S) → \Upsilon(nS)π^+π^- and h_b(mP)π^+π^- decays [1,2]. The reported masses and widths of the two resonances are M_{Z^+_b} = 10607.2 ± 2.0 MeV, Γ_{Z^+_b} = 18.4 ± 2.4 MeV and M_{Z^-_b} = 10652.2 ± 1.5 MeV, Γ_{Z^-_b} = 11.5 ± 2.2 MeV [1,2]. Analyses of the charged pion angular distributions favor the quantum numbers of the Z-states J^P = 1^+ (1^+). Evidence for the charge neutral partner Z^0_b is found in a Dalitz plot analysis of \Upsilon(5S) → \Upsilon(2S)π^0π^0 with 4.9σ significance by Belle Collaboration [3]. Its measured mass M_{Z^0_b} = 10609.8 ± 6 MeV is also consistent with that measured in the charged mode. Since Z^0_b's are isorotplet states, they need at least four quarks as minimal constituents, which makes them ideal candidates for exotic hadrons beyond the conventional q\bar{q} mesons. Note that the decay rates of \Upsilon(5S) → Z^0_bπ → \Upsilon(nS)ππ are comparable to those of \Upsilon(5S) → Z_bπ → h_b(mP)ππ. This implies unusual dynamic mechanisms undergoing the decay process since the transition to h_b(mP) would require the flip of heavy quark spin and should be suppressed in the heavy quark mass limit.

Before the observation of Z^\pm_b and Z^\pm_b, the authors predicted the existence of loosely bound S-wave B\bar{B}^* molecular states [4,5]. In Ref. [6,7], the authors predicts the possible existence of B^*(b)\bar{B}^*(b) molecular candidates within one-boson-exchange model. Since the Z^\pm_b and Z^\pm_b are charged and close to the B\bar{B}^* and B^*\bar{B}^* thresholds, many studies show that they could be S-wave B\bar{B}^* and B^*\bar{B}^* molecular states [8,14]. In Ref. [15], the masses of S-wave heavy tetraquarks b\bar{b}d\bar{d} and b\bar{d}h\bar{u} with J^P = 1^+ are extracted by the chromomagnetic interaction Hamiltonian, which turn out to be compatible with the corresponding masses of Z^+_b and Z^+_b. The QCD sum rule calculations provide a tetraquark interpretation [16]. Meanwhile, the tetraquark picture is applied to the understanding of the decays of Z^\pm_b/Z^\pm_b \rightarrow π^± \Upsilon(nS) and π^± h_b(mP) [17].

Besides the spectrum study, the production and decay of Z^\pm_b and Z^\pm_b are also investigated extensively. Considering Z^\pm_b and Z^\pm_b to be B\bar{B}^* and B^*\bar{B}^* molecular states, Voloshin estimates their production in the radiative decay of \Upsilon(5S) [13], and the pion-emission transitions from Z^\pm_b and Z^\pm_b to lighter bottomonia are investigated by Refs. [18,19]. In Ref. [14], the properties of Z^\pm_b and Z^\pm_b were studied in the framework of a nonrelativistic effective field theory assuming that Z^\pm_b and Z^\pm_b are the B\bar{B}^* and B^*\bar{B}^* molecular states.

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The intermediate meson loop transitions have been one of the important nonperturbative transition mechanisms in many processes, and their impact on the heavy quarkonium transitions, sometimes called coupled-channel effects, has been noticed for a long time [24, 23]. By applying the on-shell approximation, the bottomed meson loops were suggested to play an important role in the $\Upsilon(5S)$ transitions to the lower $\Upsilon$ states with the emission of two pions [23] or one $\eta$ [24]. This mechanism seems to explain many unusual properties that make the $\Upsilon(5S)$ different from $\Upsilon(4S)$. Similar approach was also applied to the study of $Z_b$ and $Z'_b$ by Liu et al. [11].

In this work, we will investigate the decays of $Z_b \to \Upsilon(nS)\pi$ and $Z_b \to h_b(mP)\pi$ via intermediate $B$-meson loops in an effective Lagrangian approach (ELA) with the favored quantum numbers $I^G(J^{PC}) = 1^+(1^{--})$ for the $Z_b/Z'_b$. We try to enhance the scenario by quantitative calculations that the bottomed meson loops are crucial for explaining the experimental results for $Z_b$ and $Z'_b \to B\bar{B}^* + c.c.$ and $B^*\bar{B}^*$, and $Z_b$ and $Z'_b \to \Upsilon(nS)\pi$ and $h_b(mP)\pi$.

The paper is organized as follows. In Sec. II we will introduce the formulæ for the ELA. In Sec. III the numerical results are presented. The Summary will be given in Sec. IV.

**II. TRANSITION AMPLITUDE**

In order to calculate the leading contributions from the bottomed meson loops, we need the leading order effective Lagrangians for the couplings. Based on the heavy quark symmetry and chiral symmetry [25, 26], the relevant effective Lagrangians used in this work are as follows,

$$L_{\Upsilon(nS)B^{(*)}}(B^{(*)}) = ig_{\Upsilon B\bar{B}}\bar{\Upsilon}_\nu\bar{\nu}(\partial^\mu B\bar{B} - B^\mu\bar{B}) - g_{\Upsilon B B}e^{\nu\alpha\beta}\partial_\nu\Upsilon_\alpha B^\beta + B_\beta\bar{B}_\beta$$

$$L_{h_b(mP)B^{(*)}}(B^{(*)}) = g_{h_b B\bar{B}}h_b^\mu(B\bar{B}^\mu + B_\mu\bar{B}) + ig_{h_b B B}e^{\nu\alpha\beta}\partial_\nu h_b B_\beta\bar{B}_\alpha,$$

$$L_B^{B^{(*)}\pi} = -ig_{B B\pi}(B_\alpha\partial_\alpha P_j B^\mu_j - B^\mu_j\partial_\alpha P_j B_\alpha) + \frac{1}{2}g_{B B\pi}e^{\nu\alpha\beta}B_\mu\partial_\nu P_{ij} \bar{P}_{ij} B^\beta,$$

$$L_{Z_b^{(*)}B^{(*)}} = ig_{Z_b^{(*)}B}e^{\nu\alpha\beta}Z_b^\mu B_\mu B_\beta + i g_{Z_b^{(*)}B}e^{\nu\mu\alpha\beta}Z_b^\nu B_\mu B_\beta,$$

where $B^{(*)} = \left( B^{(*)+}, B^{(*)0} \right)$ and $B^{(*)T} = \left( B^{(*)-}, B^{(*)0} \right)$ correspond to the bottom meson isodoublets.
With the experimental data for $BR(Z_b^+ \rightarrow B^+ \bar{B}^0 + \bar{B}^0 B^+)$ = (86.0 ± 3.6)% and $BR(Z_b^+ \rightarrow B^+ \bar{B}^0)$ = (73.4 ± 7.0)% from [27], we obtain $g_{Z_b B^+ B^*}$ = 13.39 GeV and $g_{Z_b' B^+ B^*}$ = 0.32. The relations

$$g_{Z_b B^+ B^*} = -g_{Z_b B^+ B^*} m_{Z_b} \sqrt{m_{B^+}}/m_{B^*}, \quad g_{Z_b' B^+ B^*} = -g_{Z_b' B^+ B^*} m_{Z_b'} \sqrt{m_{B^+}}/m_{B^*},$$

are applied to extract the couplings for $g_{Z_b B^+ B^*}$ and $g_{Z_b' B^+ B^*}$.

In Eq. (1), the following expressions are adopted in the numerical calculations,

$$g_{YYB} = 2g_2 \sqrt{m_Y m_B}, \quad g_{YYB'} = -\frac{g_{YYB}}{\sqrt{m_Y m_{B'}}}, \quad g_{YYB'} = \frac{g_{YYB}}{m_Y} \sqrt{m_{B}/m_{B'}},$$

where $g_2 = \sqrt{m_Y/(2m_B f_Y)}$: $f_Y$ and $m_Y$ denote the constant and mass of $Y(nS)$, respectively. The decay constant $f_Y$ can be extracted in $Y(nS) \rightarrow e^+ e^-$:

$$\Gamma(Y(nS) \rightarrow e^+ e^-) = \frac{4\pi g_Y^2}{27} \frac{f_Y^2}{m_Y} \frac{\text{eV}}{\text{MeV}},$$

where $\alpha_{EM} = 1/137$ is the fine-structure constant. By adopting the mass values in Table I and data for the leptonic decay widths of $Y(nS)$ states: $\Gamma(Y(1S) \rightarrow e^+ e^-) = 1.340 \pm 0.018 \text{ keV}$, $\Gamma(Y(2S) \rightarrow e^+ e^-) = 0.612 \pm 0.011 \text{ keV}$, $\Gamma(Y(3S) \rightarrow e^+ e^-) = 0.443 \pm 0.008 \text{ keV}$ [28], we obtain $f_{Y(1S)} = 715.2 \text{ MeV}$, $f_{Y(2S)} = 497.5 \text{ MeV}$, and $f_{Y(3S)} = 430.2 \text{ MeV}$.

In addition, the coupling constants in Eq. (2) are determined as

$$g_{h_0 B^+ B^*} = -2g_1 \sqrt{m_{h_0} m_{B^+}}/m_{B^*}, \quad g_{h_0 B^+ B^*} = 2g_1 \frac{m_{B^*}}{\sqrt{m_{h_0}}},$$

with $g_1 = -\sqrt{m_{h_0}/3}/f_{h_0}$, where $m_{h_0}$ and $f_{h_0}$ are the mass and decay constant of $h_0(1P)$, respectively [29], i.e. $f_{h_0} = 175 \pm 55 \text{ MeV}$ [30]. $f_{h_0(1P)} = f_{h_0(2P)} = f_{h_0(1P)} = f_{h_0(2S)} = f_{h_0(1S)} = 121.6 \text{ MeV}$.

The coupling constants relevant to the pion interactions in Eq. (3) are

$$g_{B^+ B^* B^0} = \frac{g_2}{f_Y} \sqrt{m_{B^+} m_{B^0}/m_{B^*}}, \quad g_{B^+ B^* B^0} = \frac{g_{B^+ B^0}}{\sqrt{m_{B^+} m_{B^0}/m_{B^*}}},$$

where $g = 0.44 \pm 0.03^{+0.01}_{-0.00}$ [31] and $f_Y = 132 \text{ MeV}$ are adopted in this work.

The loop transition amplitudes for the transitions in Figs. 1 and 2 can be expressed in a general form in the effective Lagrangian approach as follows:

$$M_{fi} = \int \frac{d^4q_2}{(2\pi)^4} \sum_{D^+} T_1 T_2 T_3 \prod_i \int f_i(m_i, q_i^2)$$

where $T_i (i = 1, 2, 3)$ are the vertex functions; $a_i = q_i^2 - m_i^2$ (i = 1, 2, 3) are the denominators of the intermediate meson propagators. We adopt the form factor $\prod_i f_i(m_i, q_i^2)$, which is a product of monopole form factors for each internal meson, i.e.

$$\prod_i f_i(m_i, q_i^2) = f_1(m_1, q_1^2) f_2(m_2, q_2^2) f_3(m_3, q_3^2),$$

with

$$f_i(m_i, q_i^2) = \frac{\Lambda_i^2 - m_i^2}{\Lambda_i^2 - q_i^2},$$

where $\Lambda_i = m_i + \alpha \Lambda_{QCD}$ and the QCD energy scale $\Lambda_{QCD} = 220 \text{ MeV}$. This form factor is supposed to parameterize the non-local effects of the vertex functions and remove the loop integral divergence.
The explicit transition amplitudes for $Z_b(p_i) \to B^{(*)}(q_1)B^{(*)}(q_3)[B^{(*)}(q_2)] \to Y(nS)(p_f)\pi(p_\pi)$ via those triangle loops are given as follows:

$$M_{BB'[^0]} = (i)^3 \int \frac{d^4q^2}{(2\pi)^4} [g_{Z_bB'B}][g_{Y(nS)BB'}\epsilon^g_{f'}(q_1-q_2)p][g_{B'B'}\pi p_\pi]$$

$$\times \begin{pmatrix} \frac{i}{q_1^2-m_1^2} & \frac{i}{q_2^2-m_2^2} & \frac{i}{q_3^2-m_3^2} \end{pmatrix} \prod_i F_i(m_i, q_i^2)$$

(13)

$$M_{BB'[^*]} = (i)^3 \int \frac{d^4q^2}{(2\pi)^4} [g_{Z_bB'B}][g_{Y(nS)B'B'}\epsilon^g_{f'}(q_1-q_2)p][g_{B'B'}\pi p_\pi]$$

$$\times \begin{pmatrix} \frac{i}{q_1^2-m_1^2} & \frac{i}{q_2^2-m_2^2} & \frac{i}{q_3^2-m_3^2} \end{pmatrix} \prod_i F_i(m_i, q_i^2)$$

(14)

$$M_{B'B[^0]} = (i)^3 \int \frac{d^4q^2}{(2\pi)^4} [g_{Z_bB'B}][g_{Y(nS)B'B'}\epsilon^g_{f'}(q_1-q_2)p][g_{B'B'}\pi p_\pi]$$

$$\times \begin{pmatrix} \frac{i}{q_1^2-m_1^2} & \frac{i}{q_2^2-m_2^2} & \frac{i}{q_3^2-m_3^2} \end{pmatrix} \prod_i F_i(m_i, q_i^2)$$

(15)

$$M_{B'B[^*]} = (i)^3 \int \frac{d^4q^2}{(2\pi)^4} [g_{Z_bB'B}][g_{Y(nS)B'B'}\epsilon^g_{f'}(q_1-q_2)p][g_{B'B'}\pi p_\pi]$$

$$\times \begin{pmatrix} \frac{i}{q_1^2-m_1^2} & \frac{i}{q_2^2-m_2^2} & \frac{i}{q_3^2-m_3^2} \end{pmatrix} \prod_i F_i(m_i, q_i^2)$$

(16)

$$M_{B'^B[^0]} = (i)^3 \int \frac{d^4q^2}{(2\pi)^4} [g_{Z_bB'B}][g_{Y(nS)B'B'}\epsilon^g_{f'}(q_1-q_2)p][g_{B'B'}\pi p_\pi]$$

$$\times \begin{pmatrix} \frac{i}{q_1^2-m_1^2} & \frac{i}{q_2^2-m_2^2} & \frac{i}{q_3^2-m_3^2} \end{pmatrix} \prod_i F_i(m_i, q_i^2)$$

(17)

Also, the explicit transition amplitudes for $Z_b(p_i) \to B^{(*)}(q_1)B^{(*)}(q_3)[B^{(*)}(q_2)] \to h_b(mP)(p_f)\pi(p_\pi)$ via those triangle loops are given as follows:

$$M_{BB'[^0]} = (i)^3 \int \frac{d^4q^2}{(2\pi)^4} [g_{Z_bB'B}][g_{h_b(mP)B'B'}\epsilon^g_{f'}][g_{BB'}\epsilon_{f'}]$$

$$\times \begin{pmatrix} \frac{i}{q_1^2-m_1^2} & \frac{i}{q_2^2-m_2^2} & \frac{i}{q_3^2-m_3^2} \end{pmatrix} \prod_i F_i(m_i, q_i^2)$$

(18)

$$M_{BB'[^*]} = (i)^3 \int \frac{d^4q^2}{(2\pi)^4} [g_{Z_bB'B}][g_{h_b(mP)B'B'}\epsilon^g_{f'}][g_{BB'}\epsilon_{f'}]$$

$$\times \begin{pmatrix} \frac{i}{q_1^2-m_1^2} & \frac{i}{q_2^2-m_2^2} & \frac{i}{q_3^2-m_3^2} \end{pmatrix} \prod_i F_i(m_i, q_i^2)$$

(19)

$$M_{B'B[^0]} = (i)^3 \int \frac{d^4q^2}{(2\pi)^4} [g_{Z_bB'B}][g_{h_b(mP)B'B'}\epsilon^g_{f'}][g_{BB'}\epsilon_{f'}]$$

$$\times \begin{pmatrix} \frac{i}{q_1^2-m_1^2} & \frac{i}{q_2^2-m_2^2} & \frac{i}{q_3^2-m_3^2} \end{pmatrix} \prod_i F_i(m_i, q_i^2)$$

(20)

$$M_{B'B[^*]} = (i)^3 \int \frac{d^4q^2}{(2\pi)^4} [g_{Z_bB'B}][g_{h_b(mP)B'B'}\epsilon^g_{f'}][g_{BB'}\epsilon_{f'}]$$

$$\times \begin{pmatrix} \frac{i}{q_1^2-m_1^2} & \frac{i}{q_2^2-m_2^2} & \frac{i}{q_3^2-m_3^2} \end{pmatrix} \prod_i F_i(m_i, q_i^2)$$

(21)
FIG. 3: (a) The $\alpha$-dependence of the branching ratios of $Z_b^+ \rightarrow \Upsilon(1S)\pi^+$ (solid line), $\Upsilon(2S)\pi^+$ (dashed line) and $\Upsilon(3S)\pi^+$ (dotted line). (b) The $\alpha$-dependence of the branching ratios of $Z_b^+ \rightarrow \Upsilon(1S)\pi^+$ (solid line), $\Upsilon(2S)\pi^+$ (dashed line) and $\Upsilon(3S)\pi^+$ (dotted line). The experimental values are taken from [27] as a reference.

FIG. 4: (a) The $\alpha$-dependence of the branching ratios of $Z_b^+ \rightarrow h_b(1P)\pi^+$ (solid line) and $h_b(2P)\pi^+$ (dashed line). (b) The $\alpha$-dependence of the branching ratios of $Z_b^+ \rightarrow h_b(1P)\pi^+$ (solid line) and $h_b(2P)\pi^+$ (dashed line). The experimental values are taken from [27] as a reference.

where $p_i$, $p_f$, $p_\pi$ are the four-vector momenta of the initial $Z_b$, final state bottomonium and pion, respectively, and $q_1$, $q_2$, and $q_3$ are the four-vector momenta of the intermediate bottomed mesons as defined in Figs. 1 and 2.

III. RESULTS

| States   | $\Upsilon(1S)$ | $\Upsilon(2S)$ | $\Upsilon(3S)$ | $h_b(1P)$ | $h_b(2P)$ | $B$ | $B^*$ | $\pi$ |
|----------|----------------|----------------|----------------|-----------|-----------|-----|-------|-------|
| Mass (MeV) [28] | 9460 | 10023 | 10355 | 9898 | 10259 | 5279 | 5325 | 140  |

Proceeding to the numerical results, we list the meson masses involved in the hidden-bottom decays of $Z_b/Z_b'$ in Table I. Several points concerning the determination of the form factor cutoff parameter which would be the only free parameter in a decay channel, should be clarified. First, we determine the cutoff parameter $\alpha$ for each channel separately by the experimental data. As shown in Table I it is possible to find an appropriate range of $\alpha$ values for each decay channels that can account for the data via the intermediate bottomed meson loops. Meanwhile, one notices that the $\alpha$ values for $Z_b^+ / Z_b'^+ \rightarrow \Upsilon(3S)\pi^+$ are much smaller than other channels which
TABLE II: List of branching fractions for the $Z_b^+$ and $Z_b'^+$ decays. The last column values are obtained at the average of the central $\alpha$ values exclude the $\Upsilon(3S)\pi^+$ channels. The experimental values are taken from [27] as a reference.

![Image of the table](image.png)

TABLE III: The branching ratios of decay rates for $Z_b^+ \rightarrow \Upsilon(3S)\pi^+$ and $Z_b'^+ \rightarrow \Upsilon(3S)\pi^+$ with $M_{B'} = M_B = 5279$ MeV (Scheme-I) and $M_{B'} = M_B = 5325$ MeV (Scheme-II).

![Image of the table](image.png)

indicates some unusual feature with this channel. In TableII we also list the $\alpha$ values for each decay channels that can reproduce the experimental data. An alternative test is that we make an average of the $\alpha$ values for the $Z_b$ and $Z_b'$ decays separately without including the $\Upsilon(3S)\pi$ channel, and then check whether it is possible to describe the experimental data with single values of $\alpha$ for $Z_b$ and $Z_b'$, respectively. Interestingly, as shown by the sixth column of TableII with $\alpha = 1.81$ and 1.38 for the $Z_b$ and $Z_b'$ decays, respectively, the data can be reasonably accounted for except for the $\Upsilon(3S)\pi$ channel.

![Image of the graph](image.png)

FIG. 5: (a) The $\alpha$-dependence of the ratios $R_{Z_b}^{\Upsilon(3S)\pi}$ (solid line), $R_{Z_b}^{\Upsilon(3S)\pi}$ (dashed line), $R_{Z_b}^{\Upsilon(3S)\pi}$ (dotted line), $R_{Z_b}^{\Upsilon(3S)\pi}$ (dash-dotted line) defined in Eq. (23). (b) The $\alpha$-dependence of the ratios $r_{Z_b}$ (solid line), $r_{Z_b}$ (dashed line) defined in Eq. (23).
We also check the $\alpha$-dependence of the decay branching ratios in order to give a quantitative estimate of the cutoff uncertainties in the loop integrals. The numerical results are summarized in Figs. 3 and 4 for the $Z_b$ and $Z_b'$ decays into $\Upsilon(nS)\pi$ and $h_6(mP)\pi$, respectively.

In Fig. 3(a), we plot the $\alpha$ dependence of the branching ratios of $Z^+_b \to \Upsilon(1S)\pi^+$ (solid line), $\Upsilon(2S)\pi^+$ (dashed line), and $\Upsilon(3S)\pi^+$ (dotted line), respectively. A predominant feature is that the $\alpha$ dependence of the branching ratios are quite stable, which indicates a reasonable cutoff of the ultraviolet (UV) contributions by the empirical form factor. As shown in this figure, at the same $\alpha$, the intermediate $B$-meson loop effects turn out to be more important in $Z^+_b \to \Upsilon(3S)\pi$ than in $Z^+_b \to \Upsilon(1S, 2S)\pi^+$. As a result, a smaller value of $\alpha$ is favored in $Z^+_b \to \Upsilon(3S)\pi^+$. This is understandable since the mass of $\Upsilon(3S)$ is closer to the thresholds of $BB^*$ or $B'\bar{B}'$ than the other two states [28]. Thus, it gives rise to important threshold effects in $Z^+_b \to \Upsilon(3S)\pi^+$.

One notices that the $\alpha$-dependence of the branching ratios for $Z^+_b / Z^+_b' \to \Upsilon(3S)\pi^+$ is stabler than those for $\Upsilon(1S, 2S)\pi$. This indicates that the enhanced branching ratios are not from the off-shell part of the loop integrals. As we know that the dispersive contributions become rather model-dependent near threshold, the enhanced (but rather stable in terms of $\alpha$) branching ratios for $Z^+_b / Z^+_b' \to \Upsilon(3S)\pi^+$ suggests that more stringent dynamic constraints are presumably needed to describe the near-threshold phenomena where the local quark-hadron duality has been apparently violated. What makes this process different from e.g. $\psi' \to h_1\pi^0$ in Ref. [33] is that there is no cancellations between the charged and neutral meson loops. As a consequence, the subleading terms in Refs. [33, 34] become actually leading contributions. In this sense, a new power counting in the nonrelativistic effective field theory should be exploited for the $Z_b / Z_b'$ decays [35].

Figure 3(b) presents the branching ratios of $Z^+_b / Z^+_b' \to \Upsilon(nS)\pi^+$, and the notation are the same as Fig. 3(a).

The $\alpha$ dependence of the branching ratios of $Z^+_b / Z^+_b' \to h_6(1P)\pi^+$ (solid line) and $h_6(2P)\pi^+$ (dashed line) is presented in Fig. 4. The experimental data are denoted by points for corresponding decay channels. The data for $Z^+_b \to h_6(1P)\pi^+$ and $h_6(2P)\pi^+$ can be reproduced with $\alpha = 1.76^{+0.24}_{-0.30}$ and $2.90^{+0.60}_{-0.70}$, respectively. For $Z^+_b \to h_6(1P)\pi^+$ and $h_6(2P)\pi^+$, the values of $\alpha = 1.36^{+0.20}_{-0.24}$ and $1.62^{+0.38}_{-0.43}$ can be determined by the experimental data. As shown in Table I, the decay channels for both $Z_b$ and $Z_b' \to h_6(1P, 2P)\pi$ can be reasonably accounted for by the averaged values of $\alpha = 1.81$ and 1.38, respectively. Moreover, as shown in Fig. 4, the $\alpha$ dependence turns out to be stable for both $Z_b$ and $Z_b'$ decays. The stabler behaviors for $Z^+_b$ and $Z^+_b' \to h_6(2P)\pi^+$ than $h_6(1P)\pi^+$ indicates the closeness of the $B'\bar{B}'$ threshold to the $h_6(2P)\pi^+$ threshold and the dominance of the meson loop contributions due to the open threshold effects.

It would be interesting to further clarify the uncertainties arising from the introduction of form factors by studying the $\alpha$ dependence of the ratios between different partial decay widths. For the decays of $Z^+_b / Z^+_b' \to \Upsilon(nS)\pi^+$, we define the following ratios to the partial decay widths of $Z^+_b / Z^+_b' \to \Upsilon(2S)\pi^+$:

$$ R^{12}_{Z_b} = \frac{\Gamma(Z^+_b \to \Upsilon(1S)\pi^+)}{\Gamma(Z^+_b \to \Upsilon(2S)\pi^+)} $$

$$ R^{32}_{Z_b} = \frac{\Gamma(Z^+_b \to \Upsilon(3S)\pi^+)}{\Gamma(Z^+_b \to \Upsilon(2S)\pi^+)} $$

$$ R^{12}_{Z_b'} = \frac{\Gamma(Z^+_b' \to \Upsilon(1S)\pi^+)}{\Gamma(Z^+_b' \to \Upsilon(2S)\pi^+)} $$

$$ R^{32}_{Z_b'} = \frac{\Gamma(Z^+_b' \to \Upsilon(3S)\pi^+)}{\Gamma(Z^+_b' \to \Upsilon(2S)\pi^+)} $$

(22)

which are plotted in Fig. 5(a). The stabilities of the ratios in terms of $\alpha$ indicate a reasonably controlled cutoff for each channels by the form factor.

For the decays of $Z^+_b / Z^+_b' \to h_6(1P, 2P)\pi^+$, the following ratios are defined:

$$ r_{Z_b} = \frac{\Gamma(Z^+_b \to h_6(2P)\pi^+)}{\Gamma(Z^+_b \to h_6(1P)\pi^+)} $$

$$ r_{Z_b'} = \frac{\Gamma(Z^+_b' \to h_6(2P)\pi^+)}{\Gamma(Z^+_b' \to h_6(1P)\pi^+)} $$

(23)

The $\alpha$ dependence is then plotted in Fig. 5(b), which also appears to be highly stable. Since the first coupling vertices are the same for those decay channels when taking the ratio, the stability of the ratios suggests that the transitions of $Z_b / Z_b' \to \Upsilon(nS)\pi$ and $h_6(mP)\pi$ are largely driven by the open threshold effects via the intermediate $B$ meson loops. In order to obtain this, the following analysis is carried out. First, one notices that we have adopted the couplings for the $h_6$ and $\Upsilon$ to $BB^\ast$ or $B'\bar{B}'$ in the heavy quark approximation. Since the physical masses for $B$ and $B'$ are adopted in the loop integrals, the form factor will introduce unphysical pole contributions of which the interferences with the nearby physical poles would lead to model-dependent uncertainties. By assuming
$M_{B^*} = M_B = 5279$ MeV (Scheme-I) and $M_{B^*} = M_B = 5325$ MeV (Scheme-II), namely, by partially restoring the local quark-hadron duality, we calculate the partial widths of $Z_b^+ + b \rightarrow \Upsilon(3S)\pi$. We expect that the partial restoration of the local quark-hadron duality will significantly lower the partial widths since there will be only one physical pole in the loop and the unphysical one can be easily isolated away from the physical one. As a result, the inferences caused by the closeness of the unphysical pole will be reduced. As listed in Table III in the heavy quark limit, i.e. $M_{B^*} = M_B$, the partial widths of $Z_b^+ + b \rightarrow \Upsilon(3S)\pi$ can be reproduced at much larger $\alpha$. This is a rather direct demonstration of the sensitivity of the meson loop behaviors when close to open threshold and when the dispersive part becomes dominant.

In brief, we find it is possible that with the same values of $\alpha$ for different decay channels, experimental data for the $Z_b$ and $Z_b'$ hadronic decays can be accounted for in terms of intermediate $B$ meson loops except for the $\Upsilon(3S)\pi$ channel where the close-to-threshold effect plays an important role. Recognizing this is helpful for us to understand the experimental results, and establish the the intermediate $B$ meson loops as the dominant transition mechanisms for the $Z_b$ and $Z_b'$ decays.

IV. SUMMARY

In this work, we investigate hidden-bottom decays of the newly discovered resonances $Z_b^+$ and $Z_b'^+$ via intermediate $B$-meson loops. In this calculation, the quantum numbers of the neutral partners of these two resonances are fixed to be $I^G(J^{PC}) = 1^{+}(1^{+}-)$, which is currently favored by the experimental analysis. In the ELA, the present experimental data can be reproduced with a commonly accepted range of values for the cutoff parameter $\alpha$ except for the $\Upsilon(3S)\pi$ channel where the close-to-threshold effect plays an important role in the process of $Z_b + b \rightarrow B^* \bar{B}^*(B) \rightarrow \Upsilon(3S)\pi$.

Our results show that the $\alpha$ dependence of the branching ratios are quite stable, which indicate the dominant mechanism driven by the intermediate meson loops with a fairly well control of the UV contributions. We also pointed out that the results become sensitive to the meson loop contributions when the final state mass threshold are close to the intermediate meson thresholds in our calculation. Namely, the effects from the unphysical pole introduced by the form factors would interfere with the nearby physical poles from the internal propagators and lead to model-dependent uncertainties. Such a phenomenon has been discussed in Ref. [34]. Further experimental and theoretical studies of the consequences from such an intermediate meson loop effects would be important for providing more quantitative information on the structure of $Z_b^+$ and $Z_b'^+$.

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