Exploring the $\Upsilon(4S,5S,6S) \to h_b(1P)\eta$ hidden-bottom hadronic transitions

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Recently, the Belle Collaboration has reported the measurement of the spin-flipping transition $\Upsilon(4S) \to h_b(1P)\eta$ with an unexpectedly large branching ratio: $\mathcal{B}(\Upsilon(4S) \to h_b(1P)\eta) = (2.18 \pm 0.11 \pm 0.18) \times 10^{-3}$. Such a large branching fraction contradicts with the anticipated suppression for the spin flip. In this work, we examine the effects induced by intermediate bottomed meson loops and point out that these effects are significantly important. Using the effective Lagrangian approach (ELA), we find the experimental data on $\Upsilon(4S) \to h_b(1P)\eta$ can be accommodated with the reasonable inputs. We then explore the decays $\Upsilon(5S,6S) \to h_b(1P)\eta$ and find that these two channels also have sizable branching fractions. We also calculate these processes in the framework of nonrelativistic effective field theory (NREFT). For the decays $\Upsilon(4S) \to h_b(1P)\eta$, the NREFT results are at the same order of magnitude but smaller than the ELA results by a factor of 2 to 5. For the decays $\Upsilon(5S,6S) \to h_b(1P)\eta$, the NREFT results are smaller than the ELA results by approximately 1 order of magnitude. We suggest a future experiment Belle-II to search for the $\Upsilon(5S,6S) \to h_b(1P)\eta$ decays, which will be helpful for understanding the transition mechanism.

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I. INTRODUCTION

In recent years, bottomonium transitions with an η meson or two pions in the final state have been extensively studied on the experimental side [1–7]. In 2008, the *BABAR* Collaboration first observed an enhancement for the transition $\Upsilon(4S) \to \Upsilon(1S)\eta$ compared to the dipion transition [1]. In 2011, two charged bottomoniumlike structures, $Z_b^{\pm}(10610)$ and $Z_b^{\pm}(10650)$, were observed by the Belle Collaboration in the $\pi^{\pm}\Upsilon(nS)$ and $\pi^{\pm}h_b$ invariant mass spectra of $\Upsilon(5S) \to \Upsilon(nS)\pi^+\pi^-$ and $h_b(mP)\pi^+\pi^-$ decays [2,3]. In 2015, the Belle Collaboration measured for the first time the branching fraction $\mathcal{B}(\Upsilon(4S) \to h_b(1P)\eta) = (2.18 \pm 0.11 \pm 0.18) \times 10^{-3}$ [7]. This value is anomalously large since one would expect a power suppression for the transitions with the spin flip [8,9].

A low-lying heavy quarkonium system is expected to be compact and nonrelativistic, so the QCD multipole expansion (QCDME) [8–10] can be applied to explore the hadronic transitions. For the excited states that lie above open flavor thresholds, QCDME might be problematic due to the coupled channel effects. Several possible new mechanisms have been proposed in order to explain the anomalous decay widths of $\Upsilon(4S) \to h_b(1P)\eta$. For instance, a nonrelativistic effective field theory (NREFT) is used in Ref. [11], where the branching ratio can reach the order of 10^{-3}. It has been noticed for a long time that the intermediate meson loop (IML) is one prominent nonperturbative mechanism in hadronic transitions [12–14]. In recent years, this mechanism has been successfully applied to study the production and decays of ordinary and exotic states [15–45], and a global agreement with experimental data is found. This approach has also been extensively used to study the $\Upsilon(4S,5S,6S)$ hidden bottomonium decays [46–52]. In this work, we will investigate the process $\Upsilon(4S,5S,6S) \to h_b(1P)\eta$ via the IML model. As we will show in the following, the experimental data on $\Upsilon(4S) \to h_b(1P)\eta$ can be accommodated in this approach. We then predict the branching ratios of the decays $\Upsilon(5S,6S) \to h_b(1P)\eta$ and find that they are measurable in the future.

The rest of this paper is organized as follows. We will first introduce the effective Lagrangian for our calculation in Sec. II and calculate the IML contributions to decay widths. Then, we will present our numerical results in Sec. III. A brief summary will be given in Sec. IV.

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II. RADIATIVE DECAYS

Generally speaking, all the possible intermediate meson loops should be included in the calculation. In reality, we only pick up the leading order contributions as a

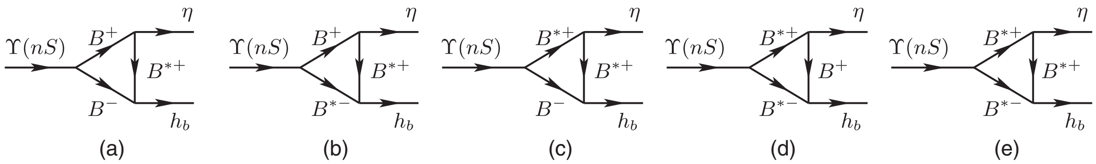

FIG. 1. The hadron-level diagrams for $\Upsilon(4S,5S,6S)\to h_b(1P)\eta$ via charged intermediate bottomed meson loops. Similar diagrams for neutral and strange intermediate bottomed meson loops.

reasonable approximation due to the breakdown of the local quark-hadron duality [12,53]. In this work, we consider the IML illustrated in Fig. 1 as the leading order contributions of $\Upsilon(4S,5S,6S)\to h_b(1P)\eta$. To calculate these diagrams, we need the effective Lagrangians to derive the couplings. Based on the heavy quark symmetry and chiral symmetry [54,55], the Lagrangian for the S- and P-wave bottomonia at leading order is given as

$$\mathcal{L}_1 = ig_1 \mathrm{Tr}[P^{\mu}_{b\bar{b}}\bar{H}_{2i}\gamma_\mu \bar{H}_{1i}] + \mathrm{H.c.}, \tag{1}$$

$$\mathcal{L}_2 = g_2 \mathrm{Tr}[R_{b\bar{b}}\bar{H}_{2i}\overleftrightarrow{\partial}_\mu \gamma^\mu \bar{H}_1] + \mathrm{H.c.}. \tag{2}$$

The S-wave bottomonium doublet and P-wave bottomonium multiplet states are expressed as

$$R_{b\bar{b}} = \frac{1+\not{v}}{2}(\Upsilon^\mu\gamma_\mu - \eta_b\gamma_5)\frac{1-\not{v}}{2}, \tag{3}$$

$$P^{\mu}_{b\bar{b}} = \frac{1+\not{v}}{2}\left(\chi^{\mu\alpha}_{b2}\gamma_\alpha + \frac{1}{\sqrt{2}}\varepsilon^{\mu\nu\alpha\beta}v_\alpha\gamma_\beta\chi_{b1\nu} + \frac{1}{\sqrt{3}}(\gamma^\mu - v^\mu)\chi_{b0} + h^\mu_b\gamma_5\right)\frac{1-\not{v}}{2}, \tag{4}$$

where Υ and η_b are the S-wave bottomonium fields. The h_b and χ_{bJ} (J = 0, 1, 2) are the P-wave bottomonium fields. The v^μ is the 4 velocity of these bottomonium states.

The bottomed and antibottomed meson triplet read

$$H_{1i} = \frac{1+\not{v}}{2}[\mathcal{B}^{*\mu}_i\gamma_\mu - \mathcal{B}_i\gamma_5], \tag{5}$$

$$H_{2i} = [\bar{\mathcal{B}}^{*\mu}_i\gamma_\mu - \bar{\mathcal{B}}_i\gamma_5]\frac{1-\not{v}}{2}, \tag{6}$$

$$\bar{H}_{1i,2i} = \gamma^0 H^\dagger_{1i,2i}\gamma^0, \tag{7}$$

where $\mathcal{B}$ and $\mathcal{B}^*$ denote the pseudoscalar and vector bottomed meson fields, respectively, i.e., $\mathcal{B}^{(*)} = (B^{+(*)}, B^{0(*)}, B^{0(*)}_s)$. v^μ is the 4 velocity of the bottomed mesons. $\varepsilon_{\mu\nu\alpha\beta}$ is the antisymmetric Levi-Civita tensor, and $\varepsilon_{0123} = +1$.

Consequently, the relevant effective Lagrangian for S-wave $\Upsilon(nS)$ and P-wave $h_b(1P)$ read

$$\begin{aligned}\mathcal{L}_{\Upsilon(nS)\mathcal{B}^{(*)}\mathcal{B}^{(*)}} = {} & ig_{\Upsilon\mathcal{B}\mathcal{B}}\Upsilon_\mu(\partial^\mu\mathcal{B}\bar{\mathcal{B}} - \mathcal{B}\partial^\mu\bar{\mathcal{B}}) \\ & - g_{\Upsilon\mathcal{B}^*\mathcal{B}}\varepsilon_{\mu\nu\alpha\beta}\partial^\mu\Upsilon^\nu(\partial^\alpha\mathcal{B}^{*\beta}\bar{\mathcal{B}} + \mathcal{B}\partial^\alpha\bar{\mathcal{B}}^{*\beta}) \\ & - ig_{\Upsilon\mathcal{B}^*\mathcal{B}^*}\{\Upsilon^\mu(\partial_\mu\mathcal{B}^{*\nu}\bar{\mathcal{B}}^*_\nu - \mathcal{B}^{*\nu}\partial_\mu\bar{\mathcal{B}}^*_\nu) \\ & + (\partial_\mu\Upsilon_\nu\mathcal{B}^{*\nu} - \Upsilon_\nu\partial_\mu\mathcal{B}^{*\nu})\bar{\mathcal{B}}^{*\mu} \\ & + \mathcal{B}^{*\mu}(\Upsilon^\nu\partial_\mu\bar{\mathcal{B}}^*_\nu - \partial_\mu\Upsilon^\nu\bar{\mathcal{B}}^*_\nu)\},\end{aligned} \tag{8}$$

$$\begin{aligned}\mathcal{L}_{h_b\mathcal{B}^{(*)}\mathcal{B}^{(*)}} = {} & g_{h_b\mathcal{B}^*\mathcal{B}}h^\mu_b(\mathcal{B}\bar{\mathcal{B}}^*_\mu + \mathcal{B}^*_\mu\bar{\mathcal{B}}) \\ & + ig_{h_b\mathcal{B}^*\mathcal{B}^*}\varepsilon_{\mu\nu\alpha\beta}\partial^\mu h^\nu_b\mathcal{B}^{*\alpha}\bar{\mathcal{B}}^{*\beta},\end{aligned} \tag{9}$$

where the coupling constants will be determined later.

The effective Lagrangian for a light pseudoscalar meson coupled to a bottomed mesons pair can be constructed using the heavy quark symmetry and chiral symmetry [54–56]

$$\begin{aligned}\mathcal{L}_{\mathcal{B}^{(*)}\mathcal{B}^{(*)}\mathcal{P}} = {} & -ig_{\mathcal{B}^*\mathcal{B}\mathcal{P}}(\mathcal{B}^i\partial^\mu\mathcal{P}_{ij}\mathcal{B}^{*j\dagger}_\mu - \mathcal{B}^{*i}_\mu\partial^\mu\mathcal{P}_{ij}\mathcal{B}^{j\dagger}) \\ & + \frac{1}{2}g_{\mathcal{B}^*\mathcal{B}^*\mathcal{P}}\varepsilon_{\mu\nu\alpha\beta}\mathcal{B}^{*\mu}_i\partial^\nu\mathcal{P}_{ij}\overleftrightarrow{\partial}_\alpha\mathcal{B}^{*\beta\dagger}_j,\end{aligned} \tag{10}$$

where $\mathcal{P}$ is a 3×3 matrix for the pseudoscalar octet. The physical states η is the linear combinations of $n\bar{n} = (u\bar{u}+d\bar{d})/\sqrt{2}$ and $s\bar{s}$ with the mixing scheme,

$$|\eta\rangle = \cos\alpha_P|n\bar{n}\rangle - \sin\alpha_P|s\bar{s}\rangle. \tag{11}$$

The mixing angle is given as $\alpha_P \simeq \theta_P + \arctan\sqrt{2}$, where the empirical value for the θ_P should be in the range $-24.6°{-}-11.5°$ [57]. In this work, we will take $\theta_P = -19.3°$ [58].

With the above Lagrangians, we can derive the transition amplitudes for $\Upsilon(nS)(p_1)\to[B^{(*)}(q_1)\bar{B}^{(*)}(q_3)]B^{(*)}(q_2)\to h_b(1P)(p_2)\eta(p_3)$ shown in Fig. 1,

$$\begin{aligned}
\mathcal{M}_{BB[B^*]} &= \int \frac{d^4q_2}{(2\pi)^4}[-2g_{\Upsilon BB}\varepsilon_{1\mu}q_2^\mu][-g_{B^*BP}p_{3\nu}][g_{h_bB^*B}\varepsilon_{2\alpha}] \\
&\quad \times \frac{i}{q_1^2-m_1^2}\frac{i(-g^{\nu\alpha}+q_2^\nu q_2^\alpha/m_2^2)}{q_2^2-m_2^2}\frac{i}{q_3^2-m_3^2}\mathcal{F}(m_2,q_2^2), \\
\mathcal{M}_{BB*[B^*]} &= \int \frac{d^4q_2}{(2\pi)^4}[g_{\Upsilon B^*B}\varepsilon_{\mu\nu\alpha\beta}p_1^\mu\varepsilon_1^\nu q_3^\alpha][-g_{B^*BP}p_{3\delta}][g_{h_bB^*B^*}\varepsilon_{\theta\phi\kappa\lambda}p_2^\theta\varepsilon_2^\phi] \\
&\quad \times \frac{i}{q_1^2-m_1^2}\frac{i(-g^{\delta\kappa}+q_2^\delta q_2^\kappa/m_2^2)}{q_2^2-m_2^2}\frac{i(-g^{\beta\lambda}+q_3^\beta q_3^\lambda/m_3^2)}{q_3^2-m_3^2}\mathcal{F}(m_2,q_2^2), \\
\mathcal{M}_{B^*B[B^*]} &= \int \frac{d^4q_2}{(2\pi)^4}[-g_{\Upsilon B^*B}\varepsilon_{\mu\nu\alpha\beta}p_1^\mu\varepsilon_1^\nu q_1^\alpha][g_{B^*B^*P}\varepsilon_{\theta\phi\kappa\lambda}p_3^\phi 2q_2^\kappa][g_{h_bB^*B}\varepsilon_{2\delta}] \\
&\quad \times \frac{i(-g^{\beta\theta}+q_1^\beta q_1^\theta/m_1^2)}{q_1^2-m_1^2}\frac{i(-g^{\lambda\delta}+q_2^\lambda q_2^\delta/m_2^2)}{q_2^2-m_2^2}\frac{i}{q_3^2-m_3^2}\mathcal{F}(m_2,q_2^2), \\
\mathcal{M}_{B^*B^*[B]} &= \int \frac{d^4q_2}{(2\pi)^4}[-g_{\Upsilon B^*B^*}\varepsilon_1^\mu(2q_{1\alpha}g_{\mu\nu}-q_{1\mu}g_{\alpha\nu}+q_{3\nu}g_{\mu\alpha})][g_{B^*BP}p_{3\beta}][g_{h_bB^*B}\varepsilon_{2\delta}] \\
&\quad \times \frac{i(-g^{\nu\beta}+q_1^\nu q_1^\beta/m_1^2)}{q_1^2-m_1^2}\frac{i}{q_2^2-m_2^2}\frac{i(-g^{\alpha\delta}+q_3^\alpha q_3^\delta/m_3^2)}{q_3^2-m_3^2}\mathcal{F}(m_2,q_2^2), \\
\mathcal{M}_{B^*B^*[B^*]} &= \int \frac{d^4q_2}{(2\pi)^4}[-g_{\Upsilon B^*B^*}\varepsilon_1^\mu(2q_{1\alpha}g_{\mu\nu}-q_{1\mu}g_{\alpha\nu}+q_{3\nu}g_{\mu\alpha})][g_{B^*B^*P}\varepsilon_{\theta\phi\kappa\lambda}p_3^\phi q_2^\kappa][g_{h_bB^*B^*}\varepsilon_{\beta\rho\sigma\delta}p_2^\beta\varepsilon_2^\rho] \\
&\quad \times \frac{i(-g^{\nu\theta}+q_1^\nu q_1^\theta/m_1^2)}{q_1^2-m_1^2}\frac{i(-g^{\alpha\delta}+q_2^\alpha q_2^\delta/m_2^2)}{q_2^2-m_2^2}\frac{i(-g^{\lambda\rho}+q_3^\lambda q_3^\rho/m_3^2)}{q_3^2-m_3^2}\mathcal{F}(m_2,q_2^2), \qquad (12)
\end{aligned}$$

where p_1, p_2, and p_3 are the four momenta of the initial state $\Upsilon(nS)$, final state $h_b(1P)$, and η, respectively. ε_1 and ε_2 are the polarization vector of $\Upsilon(nS)$ and $h_b(1P)$, respectively. q_1, q_3, and q_2 are the four momenta of the bottomed meson connecting $\Upsilon(nS)$ and η, the bottomed meson connecting $\Upsilon(nS)$ and $h_b(1P)$, and the exchanged bottomed meson, respectively.

In the triangle diagrams of Fig. 1, the exchanged bottomed mesons are off shell. To compensate the off shell effects and regularize the ultraviolet divergence [59–61], we introduce the monopole form factor,

$$\mathcal{F}(m_2,q_2^2)=\frac{\Lambda^2-m_2^2}{\Lambda^2-q_2^2}, \qquad (13)$$

where q_2 and m_2 are the momentum and mass of the exchanged bottomed meson, respectively. The parameter $\Lambda \equiv m_2+\alpha\Lambda_{\rm QCD}$, and the QCD energy scale $\Lambda_{\rm QCD}=220$ MeV. The dimensionless parameter α, which is usually of order 1, depends on the specific process.

III. NUMERICAL RESULTS

With the experimental data on the decay width of $\Upsilon(4S)\to B\bar{B}$ [57], the coupling constant $g_{\Upsilon(4S)BB}$ is determined as $g_{\Upsilon(4S)BB}=24.2$ which is comparable to the estimation in the vector meson dominance model. Since the mass of $\Upsilon(4S)$ is only above the $B\bar{B}$ threshold, the coupling constants $g_{\Upsilon(4S)B^*B}$ and $g_{\Upsilon(4S)B^*B^*}$ are determined as follows:

$$g_{\Upsilon(4S)B^*B}=\frac{g_{\Upsilon(4S)BB}}{\sqrt{m_{B^*}m_B}},$$
$$g_{\Upsilon(4S)B^*B^*}=g_{\Upsilon(4S)B^*B}\sqrt{\frac{m_{B^*}}{m_B}}m_{B^*}. \qquad (14)$$

For the coupling constants between $\Upsilon(5S)$ and $B^{(*)}\bar{B}^{(*)}$, we use the experimental data on the decay width of $\Upsilon(5S)\to B^{(*)}\bar{B}^{(*)}$ [57]. The measured branching ratios and the corresponding coupling constants are given in Table I. One can see that the values determined from the $\Upsilon(5S)$ data in Table I are very small. This is partly due to the fact that as a high-excited $b\bar{b}$ state, the wave function of $\Upsilon(5S)$ has a complicated node structure, and the coupling constants will be small if the p values of $B^{(*)}\bar{B}^{(*)}$ channels (1060–1270 MeV) are those corresponding to the zeros in the amplitude [48]. Since there is no experimental information on $\Upsilon(6S)\to B^{(*)}\bar{B}^{(*)}$ [57], we choose the same values as the $\Upsilon(5S)$ ones.

The coupling constants between $h_b(1P)$ and $B^{(*)}\bar{B}^*$ in Eq. (9) are determined as

$$g_{h_bBB^*}=-2g_1\sqrt{m_{h_b}m_Bm_{B^*}}, \qquad g_{h_bB^*B^*}=2g_1\frac{m_{B^*}}{\sqrt{m_{h_b}}}, \qquad (15)$$

TABLE I. The coupling constants of $\Upsilon(5S)$ interacting with $B^{(*)}\bar{B}^{(*)}$. Here, we list the corresponding branching ratios of $\Upsilon(5S) \to B^{(*)}\bar{B}^{(*)}$.

Final state	$\mathcal{B}(\%)$	Coupling	Final state	$\mathcal{B}(\%)$	Coupling	Final state	$\mathcal{B}(\%)$	Coupling
$B\bar{B}$	5.5	1.76	$B\bar{B}^* + \text{c.c.}$	13.7	$0.14\ \text{GeV}^{-1}$	$B^*\bar{B}^*$	38.1	2.22
$B_s\bar{B}_s$	0.5	0.96	$B_s\bar{B}_s^* + \text{c.c.}$	1.35	$0.10\ \text{GeV}^{-1}$	$B_s^*\bar{B}_s^*$	17.6	5.07

where $g_1 = -\sqrt{m_{\chi_{b0}}/3}/f_{\chi_{b0}}$. $m_{\chi_{b0}}$ and $f_{\chi_{b0}}$ are the mass and decay constant of $\chi_{b0}(1P)$, respectively [62], i.e., $f_{\chi_{b0}} = 175 \pm 55$ MeV [63].

In the chiral and heavy quark limits, the couplings between bottomed meson pair and light pseudoscalar mesons have the following relationships [55]:

$$g_{\mathcal{B}^*\mathcal{B}^*P} = \frac{g_{\mathcal{B}^*\mathcal{B}P}}{\sqrt{m_{\mathcal{B}} m_{\mathcal{B}^*}}} = \frac{2}{f_\pi} g, \tag{16}$$

where $f_\pi = 132$ MeV is the pion decay constant and $g = 0.59$ [64].

For the tree-level contributions to $\Upsilon(nS) \to h_b(1P)\eta$, the amplitude scales as the quark mass difference

$$\mathcal{M}^{\text{tree}} \sim \delta \tag{17}$$

with $\delta = m_s - (m_u + m_d)/2$.

For the bottom meson loop contributions in Fig. 1, the decay amplitude scales as follows:

$$\mathcal{M}^{loop} \sim \mathcal{N} \frac{q^2}{\bar{v}^3 M_B^2} \triangle, \tag{18}$$

where $\mathcal{N} = 1/(2\sqrt{3}\pi v_b^4)$, q is the final η momentum, $\bar{v}$ is understood as the average velocity of the intermediate bottomed mesons. The meson mass difference $\triangle$ denotes the violation of the $SU(3)$ symmetry, which has a similar size as δ. v_b denotes the bottom quark velocity inside the bottomonia, and we take $v_b = \sqrt{0.1}$ here.

For $\Upsilon(4S) \to h_b(1P)\eta$ decay, the momentum of the emitted η is $q \simeq 388$ MeV, and the velocity v is about $\sqrt{[2m_B - (m_{\Upsilon(4S)} + m_{h_b})/2]/m_B} \simeq 0.28$. As a result, the factor $\mathcal{N} q^2/(\bar{v}^3 M_B^2)$ is about 2.17, which gives an enhancement compared with the tree-level contributions. For $\Upsilon(5S) \to h_b(1P)\eta$, the velocity $\bar{v} \simeq 0.23$ and $q = 750$ MeV, so the factor $\mathcal{N} q^2/(\bar{v}^3 M_B^2)$ is about 15. For $\Upsilon(6S) \to h_b(1P)\eta$, the velocity $\bar{v} \simeq 0.19$ and $q = 930$ MeV, so the factor $\mathcal{N} q^2/(\bar{v}^3 M_B^2)$ is about 37. According to our power counting analysis, the transitions $\Upsilon(4S, 5S, 6S) \to h_b(1P)\eta$ are dominated by the meson loops.

In Fig. 2(a), we plot the branching ratios for $\Upsilon(4S) \to h_b(1P)\eta$ in terms of the cutoff parameter α with the monopole form factor. We also zoom into details of the figure with a narrow range $\alpha = 0.1$–0.2 in order to show the best fit of the α parameter. As shown in Fig. 2(a), the branching ratios are not drastically sensitive to the cutoff parameter α. Our calculated branching ratios can reproduce the experimental data [57] at about $\alpha = 0.12$. In Figs. 2(b) and 2(c), we plot the predicted branching ratios for $\Upsilon(5S) \to h_b(1P)\eta$ and $\Upsilon(6S) \to h_b(1P)\eta$ in terms of the cutoff parameter α with the monopole form factor. The behavior is similar to that of $\Upsilon(4S) \to h_b(1P)\eta$ in Fig. 2(a). The predicted branching ratios of $\Upsilon(5S) \to h_b(1P)\eta$ are about 10^{-3}–10^{-2} with the commonly accepted α range. For $\Upsilon(6S) \to h_b(1P)\eta$, the results are much smaller, which are about 10^{-4}–10^{-2}. At the same α, the predicted branching ratios of $\Upsilon(5S) \to h_b(1P)\eta$ are about 1 order of magnitude smaller than that of $\Upsilon(4S) \to h_b(1P)\eta$. For $\Upsilon(6S) \to h_b(1P)\eta$, the predicted branching ratios are about 2 orders smaller than that of $\Upsilon(4S) \to h_b(1P)\eta$. We suggest a future experiment Belle-II to carry out the search for the spin-flipping transitions $\Upsilon(5S, 6S) \to h_b(1P)\eta$, which will help us understand the decay mechanism. Here, we should notice several uncertainties that may influence our numerical results, such as the coupling constants and off shell effects arising from the exchanged particles of the loops, and the cutoff parameter can also be different in decay channels.

In order to illustrate the impact of the η-η' mixing angle, in Fig. 3, we present the branching ratios in terms of the η-η' mixing angle with $\alpha = 0.15$ (solid line) and 0.25 (dashed line), respectively. As shown in this figure, when the η-η' mixing angle α_P increases, the branching ratios of $\Upsilon(4S) \to h_b(1P)\eta$ decrease, while the branching ratios of $\Upsilon(5S, 6S) \to h_b(1P)\eta$ increase. This behavior suggests how the η-η' mixing angle influences our calculated results to some extent.

As a comparison, in Fig. 3, we also give the results using the NREFT approach denoted as dotted lines. The NREFT approach provides a systematic tool to control the uncertainties [11,34,65]. From the figure, one can see that for the decays $\Upsilon(4S) \to h_b(1P)\eta$, the NREFT results are at the same order of magnitude but smaller than the effective Lagrangian approach (ELA) results by a factor of 2 to 5. These differences may give some sense of the theoretical uncertainties for the calculated rates and indicate the viability of our model to some extent. However, for transitions where the mass difference between the initial and final state becomes large, the NREFT may be not applicable. From Fig. 3, one can see that for the decays $\Upsilon(5S, 6S) \to h_b(1P)\eta$, the NREFT results are smaller than

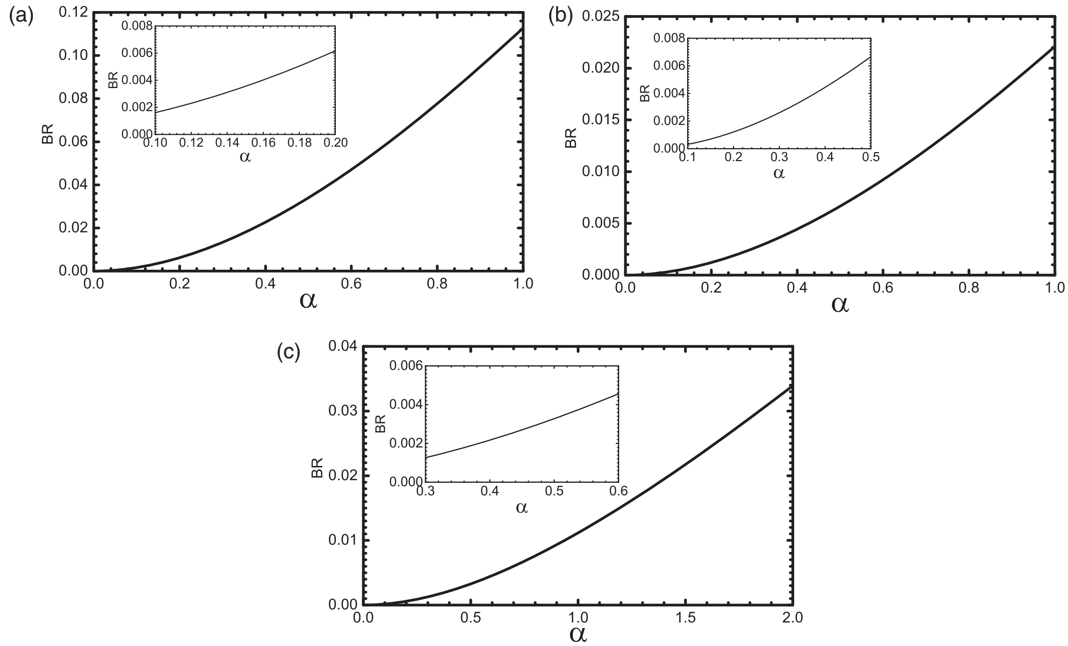

FIG. 2. (a) The α dependence of the branching ratios of $\Upsilon(4S) \to h_b(1P)\eta$. (b) The α dependence of the branching ratios of $\Upsilon(5S) \to h_b(1P)\eta$. (c). The α dependence of the branching ratios of $\Upsilon(6S) \to h_b(1P)\eta$.

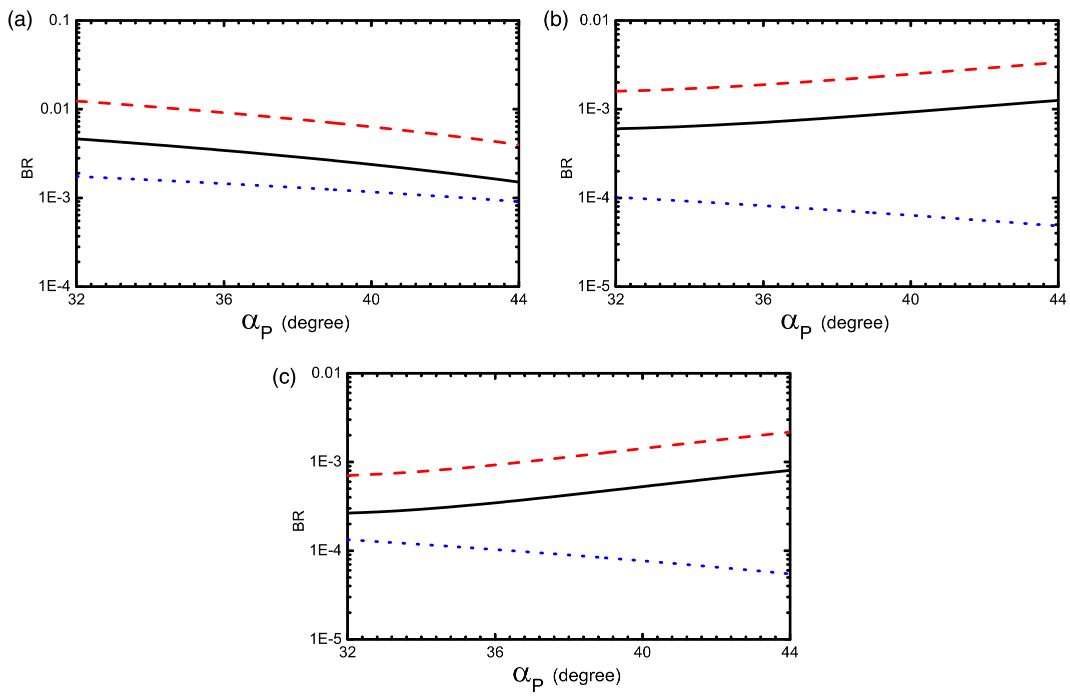

FIG. 3. (a) The dependence of branching ratios of $\Upsilon(4S) \to h_b(1P)\eta$ on the η-η' mixing angle with the cutoff parameter $\alpha = 0.15$ (solid line) and 0.25 (dashed line), respectively. The calculated branching ratios in the NREFT approach are presented with a dotted line. (b) The branching ratios of $\Upsilon(5S) \to h_b(1P)\eta$ in terms of the η-η' mixing angle with $\alpha = 0.15$ (solid line) and 0.25 (dashed line), respectively. The calculated branching ratios in NREFT approach are presented with a dotted line. (c). The branching ratios of $\Upsilon(6S) \to h_b(1P)\eta$ in terms of the η-η' mixing angle with $\alpha = 0.15$ (solid line) and 0.25 (dashed line), respectively. The calculated branching ratios in the NREFT approach are presented with a dotted line.

the ELA results by approximately 1 order of magnitude. We suggest a future experiment Belle-II to carry out the search for this anomalous $\Upsilon(5S,6S) \to h_b(1P)\eta$ transitions, which will help us test this point.

IV. SUMMARY

Recent experiments on the $\Upsilon(4S) \to h_b(1P)\eta$ transition show anomalously large decay rates. This seems to contradict the naive expectation that hadronic transitions with spin flipping terms should be suppressed with respect those that do not have these terms. In this work, we have studied the spin-flipping transitions of $\Upsilon(4S,5S,6S) \to h_b(1P)\eta$ via intermediate bottomed meson loops in an effective Lagrangian approach. Our results have shown that the intermediate bottomed meson loops can play an important role in these processes, especially when the initial states are close to the two particle thresholds. For $\Upsilon(4S) \to h_b(1P)\eta$, the experimental data can be reproduced in this approach with a commonly accepted range of values for the form factor cutoff parameter α. We also predict the branching ratios of $\Upsilon(5S) \to h_b(1P)\eta$, which are about orders of 10^{-3}–10^{-2}. For $\Upsilon(6S) \to h_b(1P)\eta$, the results are much smaller, which are about 10^{-4}–10^{-2}. As a cross-check, we also calculated the branching ratios of the decays in the framework of NREFT. We suggest a future experiment Belle-II to carry out the search for the spin-flipping transitions $\Upsilon(5S,6S) \to h_b(1P)\eta$, which will help us understand the decay mechanism.

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