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Investigating possible decay modes of $Y(4260)$ under $D_1(2420)\bar{D} + c.c.$ molecular state ansatz

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By assuming that $Y(4260)$ is a $D_1\bar{D}$ molecular state, we investigate some hidden-charm and charmed pair decay channels of $Y(4260)$ via intermediate $D_1\bar{D}$ meson loops with an effective Lagrangian approach. Through investigating the α dependence of branching ratios and ratios between different decay channels, we show that the intermediate $D_1\bar{D}$ meson loops are crucial for driving these transitions of $Y(4260)$ studied here. The coupled channel effects turn out to be more important in $Y(4260) \to D^*\bar{D}^*$, which can be tested in future experiments.

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I. INTRODUCTION

During the past years, the experimental observation of a large number of so-called XYZ states has initiated tremendous efforts to unravel their nature beyond the conventional quark model (for recent reviews, see, e.g., Refs. [1–5]). $Y(4260)$ was reported by the *BABAR* Collaboration in the $\pi^+\pi^- J/\psi$ invariant spectrum in $e^+e^- \to \gamma_{ISR}\pi^+\pi^- J/\psi$ [6], which has been confirmed both by the CLEO and Belle collaborations [7,8]. Its mass and total width are well determined as $m = 4263^{+8}_{-9}$ MeV and $\Gamma_Y = 95 \pm 14$ MeV, respectively [9]. The new datum from BESIII confirms the signal in $Y(4260) \to J/\psi\pi^+\pi^-$ with much higher statistics [10]. The mass of $Y(4260)$ does not agree to what is predicted by the potential quark model. Furthermore, the most mysterious fact is that as a charmonium state with $J^{PC} = 1^{--}$, it is only "seen" as a bump in the two pion transitions to J/ψ, but not in any open charm decay channels like $D\bar{D}$, $D^*\bar{D} + c.c.$, and $D^*\bar{D}^*$, or other measured channels. The line shapes of the cross section for e^+e^- annihilations into $D^{(*)}$ meson pairs appear to have a dip at its peak mass 4.26 GeV instead of a bump.

Since the observation of the $Y(4260)$, many theoretical investigations have been carried out (for a review see Ref. [11]). It has variously been identified as a conventional $\psi(4S)$ based on a relativistic quark model [12], a tetraquark $c\bar{c}s\bar{s}$ state [13], a charmonium hybrid [14–16], hadronic molecule of $D_1\bar{D}$ [17–19],[1] $\chi_{c1}\omega$ [23], $\chi_{c1}\rho$ [24], $J/\psi f_0$ [25], a cusp [26,27] or a nonresonance explanation [28,29], etc. The dynamical calculation of tetraquark states indicated that $Y(4260)$ cannot be interpreted as a P-wave 1^{--} state of the charm-strange diquark-antidiquark, because the corresponding mass is found to be 200 MeV heavier [30]. In Ref. [31], the authors also studied the possibility of $Y(4260)$ as a P-wave 1^{--} state of the charm-strange diquark-antidiquark state in the framework of QCD sum rules and arrived at the same conclusion as Ref. [30]. Some lattice calculations give the mass of the vector hybrid within this mass region [32], which is very close to the new charmoniumlike state $Y(4360)$ [33]. With the $D_1\bar{D}$ molecular ansatz, a consistent description of some of the experimental observations can be obtained, such as its nonobservation in open charm decays, or the observation of $Z_c(3900)$ as mentioned in Ref. [19]; the threshold behavior in its main decay channels is investigated in Ref. [34] and the production of $X(3872)$ is studied in the radiative decays of $Y(4260)$ [22]. Under such a molecular state assumption, a consistent description of many experimental observations could be obtained. However, as studied in [35], the production of an S-wave $D_1\bar{D}$ pair in e^+e^- annihilation is forbidden in the limit of exact heavy quark spin symmetry, which substantially weakens the arguments for considering the $Y(4260)$ charmoniumlike resonance as a $D_1\bar{D}$ molecular state.

The intermediate meson loop (IML) transition is one of the possible nonperturbative dynamical mechanisms, especially when we investigate the pertinent issues in the energy region of charmonium [36–61]. During the past decade, many interesting observations were announced by Belle, *BABAR*, CLEO, BESIII, and so on. And in theoretical study, it is widely recognized that the IML may be closely related to many nonperturbative phenomena observed in experiments [44–64], e.g., apparent Okubo-Zweig-Iizuka–rule violations, sizable non-$D\bar{D}$ decay branching ratios for $\psi(3770)$ [44–49], the helicity selection rule violations in charmonium decays [56–58], and the hidden charmonium decays of the newly discovered Z_c [62].

In this work, we will investigate the hidden-charm decays of $Y(4260)$ and $Y(4260) \to D^{(*)}\bar{D}^{(*)}$ via the $D_1\bar{D}$ loop with an effective Lagrangian approach (ELA) under the $D_1\bar{D}$ molecular assumption. The paper is organized as follows. In Sec. II, we will introduce the ELA briefly and give some

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[1]Notice that there are two D_1 states of similar masses, and the one in question should be the narrower one, i.e., the $D_1(2420)$ ($\Gamma = 27$ MeV); the $D_1(2430)$ ($\Gamma \simeq 384$ MeV) is too broad to form a molecular state [20–22].

relevant formulas. In Sec. III, the numerical results are presented. The summary will be given in Sec. IV.

II. THE MODEL

To calculate the leading contributions from the charmed meson loops, we need the leading order effective Lagrangian for the couplings. Based on the heavy quark symmetry and chiral symmetry [65,66], the relevant effective Lagrangian used in this work read

$$\begin{aligned}\mathcal{L}_{\psi D^{(*)}D^{(*)}} = {} & ig_{\psi DD}\psi_\mu(\partial^\mu D\bar{D} - D\partial^\mu \bar{D}) \\ & - g_{\psi D^* D}\varepsilon^{\mu\nu\alpha\beta}\partial_\mu\psi_\nu(\partial_\alpha D^*_\beta \bar{D} + D\partial_\alpha \bar{D}^*_\beta) \\ & - ig_{\psi D^* D^*}\{\psi^\mu(\partial_\mu D^{*\nu}\bar{D}^*_\nu - D^{*\nu}\partial_\mu \bar{D}^*_\nu) \\ & + (\partial_\mu \psi_\nu D^{*\nu} - \psi_\nu \partial_\mu D^{*\nu})\bar{D}^{*\mu} \\ & + D^{*\mu}(\psi^\nu\partial_\mu \bar{D}^*_\nu - \partial_\mu\psi^\nu \bar{D}^*_\nu)\}, \end{aligned} \tag{1}$$

$$\begin{aligned}\mathcal{L}_{h_c D^{(*)}D^{(*)}} = {} & g_{h_c D^* D} h_c^\mu (D\bar{D}^*_\mu + D^*_\mu \bar{D}) \\ & + ig_{h_c D^* D^*}\varepsilon^{\mu\nu\alpha\beta}\partial_\mu h_{c\nu} D^*_\alpha \bar{D}^*_\beta, \end{aligned} \tag{2}$$

where $D^{(*)} = (D^{(*)+}, D^{(*)0}, D_s^{(*)+})$ and $\bar{D}^{(*)T} = (D^{(*)-}, \bar{D}^{(*)0}, D_s^{(*)-})$ correspond to the charmed meson isodoublets. The following couplings are adopted in the numerical calculations:

$$g_{\psi DD} = 2g_2\sqrt{m_\psi}m_D, \qquad g_{\psi D^* D} = \frac{g_{\psi DD}}{\sqrt{m_D m_{D^*}}}, \qquad g_{\psi D^* D^*} = g_{\psi D^* D}\sqrt{\frac{m_{D^*}}{m_D}}m_{D^*}. \tag{3}$$

In principle, the parameter g_2 should be computed with nonperturbative methods. It shows that vector meson dominance would provide an estimate of these quantities [65]. The coupling g_2 can be related to the J/ψ leptonic constant f_ψ, which is defined by the matrix element $\langle 0|\bar{c}\gamma_\mu c|J/\psi(p,\epsilon)\rangle = f_\psi m_\psi \epsilon^\mu$, and $g_2 = \sqrt{m_\psi}/2m_D f_\psi$, where $f_\psi = 405 \pm 14$ MeV, and we have applied the relation $g_{\psi DD} = m_\psi/f_\psi$. The ratio of the coupling constants $g_{\psi' DD}$ to $g_{\psi DD}$ is fixed as that in Ref. [57], i.e.,

$$\frac{g_{\psi' DD}}{g_{\psi DD}} = 0.9. \tag{4}$$

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

$$g_{h_c DD^*} = -2g_1\sqrt{m_{h_c} m_D m_{D^*}}, \quad g_{h_c D^* D^*} = 2g_1\frac{m_{D^*}}{\sqrt{m_{h_c}}}, \tag{5}$$

with $g_1 = -\sqrt{m_{\chi_{c0}}/3}/f_{\chi_{c0}}$, where $m_{\chi_{c0}}$ and $f_{\chi_{c0}}$ are the mass and decay constant of $\chi_{c0}(1P)$, respectively [67]. We take $f_{\chi_{c0}} = 510 \pm 40$ MeV [68].

The light vector mesons nonet can be introduced by using the hidden gauge symmetry approach, and the effective Lagrangian containing these particles are as follows [69,70]:

$$\begin{aligned}\mathcal{L}_{\mathrm{D^*DV}} = {} & ig_{\mathrm{D^*DV}}\epsilon_{\alpha\beta\mu\nu}(\mathrm{D}_b\overleftrightarrow{\partial}_\alpha \mathrm{D}_a^{*\beta\dagger} - \mathrm{D}_b^{*\beta\dagger}\overleftrightarrow{\partial}_\alpha \mathrm{D}_a^j)(\partial^\mu V^\nu)_{ba} \\ & + ig_{\bar{\mathrm{D}}^*\bar{\mathrm{D}}V}\epsilon_{\alpha\beta\mu\nu}(\bar{\mathrm{D}}_b\overleftrightarrow{\partial}_\alpha \bar{\mathrm{D}}_a^{*\beta\dagger} - \bar{\mathrm{D}}_b^{*\beta\dagger}\overleftrightarrow{\partial}_\alpha \bar{\mathrm{D}}_a^j) \\ & \times(\partial^\mu V^\nu)_{ab} + \mathrm{H.c.}, \\ \mathcal{L}_{\mathrm{DDV}} = {} & ig_{\mathrm{DDV}}(\mathrm{D}_b\overleftrightarrow{\partial}_\mu \mathrm{D}_a^\dagger)V^\mu_{ba} + ig_{\bar{\mathrm{D}}\bar{\mathrm{D}}V}(\bar{\mathrm{D}}_b\overleftrightarrow{\partial}_\mu \bar{\mathrm{D}}_a^\dagger)V^\mu_{ab}, \\ \mathcal{L}_{\mathrm{DD_1}V} = {} & g_{\mathrm{DD_1}V}\mathrm{D}^\mu_{1b}V_{\mu ba}\mathrm{D}_a^\dagger + g'_{\mathrm{DD_1}V}(\mathrm{D}^\mu_{1b}\overleftrightarrow{\partial}^\nu \mathrm{D}_a^\dagger) \\ & \times(\partial_\mu V_\nu - \partial_\nu V_\mu)_{ba} + g_{\bar{\mathrm{D}}\bar{\mathrm{D}}_1V}\bar{\mathrm{D}}_a^\dagger V_{\mu ab}\bar{\mathrm{D}}^\mu_{1b} \\ & + g'_{\bar{\mathrm{D}}\bar{\mathrm{D}}_1V}(\bar{\mathrm{D}}^\mu_{1b}\overleftrightarrow{\partial}^\nu \bar{\mathrm{D}}_a^\dagger)(\partial_\mu V_\nu - \partial_\nu V_\mu)_{ab} + \mathrm{H.c.}, \\ \mathcal{L}_{\mathrm{D^*D^*V}} = {} & ig_{\mathrm{D^*D^*V}}(\mathrm{D}^*_{b\nu}\overleftrightarrow{\partial}_\mu \mathrm{D}_a^{*\nu\dagger})V^\mu_{ba} \\ & + ig'_{\mathrm{D^*D^*V}}(\mathrm{D}_b^{*\mu}\mathrm{D}_a^{*\nu\dagger} - \mathrm{D}_a^{*\mu\dagger}\mathrm{D}_b^{*\nu})(\partial_\mu V_\nu - \partial_\nu V_\mu)_{ba} \\ & + ig_{\bar{\mathrm{D}}^*\bar{\mathrm{D}}^*V}(\bar{\mathrm{D}}^*_{b\nu}\overleftrightarrow{\partial}_\mu \bar{\mathrm{D}}_a^{*\nu\dagger})V^\mu_{ab} \\ & + ig'_{\bar{\mathrm{D}}^*\bar{\mathrm{D}}^*V}(\bar{\mathrm{D}}_b^{*\mu}\bar{\mathrm{D}}_a^{*\nu\dagger} - \bar{\mathrm{D}}_a^{*\mu\dagger}\bar{\mathrm{D}}_b^{*\nu})(\partial_\mu V_\nu - \partial_\nu V_\mu)_{ab}. \end{aligned} \tag{6}$$

And the coupling constants read

$$\begin{aligned} g_{\mathrm{DD}V} &= -g_{\bar{\mathrm{D}}\bar{\mathrm{D}}V} = \frac{1}{\sqrt{2}}\beta g_V, \\ g_{\mathrm{DD_1}V} &= -g_{\bar{\mathrm{D}}\bar{\mathrm{D}}_1V} = -\frac{2}{\sqrt{3}}\zeta_1 g_V\sqrt{\mathrm{m_D m_{D_1}}}, \\ g'_{\mathrm{DD_1}V} &= -g'_{\bar{\mathrm{D}}\bar{\mathrm{D}}_1V} = \frac{1}{\sqrt{3}}\mu_1 g_V, \\ g_{\mathrm{D^*D^*}V} &= -g_{\bar{\mathrm{D}}^*\bar{\mathrm{D}}^*V} = -\frac{1}{\sqrt{2}}\beta g_V, \\ g'_{\mathrm{D^*D^*}V} &= -g'_{\bar{\mathrm{D}}^*\bar{\mathrm{D}}^*V} = -\sqrt{2}\lambda g_V \mathrm{m_{D^*}}, \end{aligned} \tag{7}$$

where $f_\pi = 132$ MeV is the pion decay constant, and the parameter g_V is given by $g_V = m_\rho/f_\pi$ [66]. We take $\lambda = 0.56\ \mathrm{GeV}^{-1}$, $g = 0.59$, and $\beta = 0.9$ in our calculation [71].

The effective Lagrangian for the light pseudoscalar mesons are constructed by imposing invariance under both heavy quark spin-flavor transformation and chiral transformation [66,72–74]. The pertinent interaction terms for this work read

$$\begin{aligned}\mathcal{L}_{D_1 D^* P} = {} & g_{D_1 D^* P}[3D^\mu_{1a}(\partial_\mu\partial_\nu\phi)_{ab}D_b^{*\dagger\nu} \\ & - D^\mu_{1a}(\partial^\nu\partial_\nu\phi)_{ab}D^{*\dagger}_{b\mu}] \\ & + g_{\bar{D}_1\bar{D}^* P}[3\bar{D}_a^{*\dagger\mu}(\partial_\mu\partial_\nu\phi)_{ab}\bar{D}^\nu_{1b} \\ & - \bar{D}_a^{*\dagger\mu}(\partial^\nu\partial_\nu\phi)_{ab}\bar{D}_{1b\nu}] + \mathrm{H.c.}, \end{aligned} \tag{8}$$

$$\begin{aligned}\mathcal{L}_{\mathrm{DD^*P}} = {} & g_{\mathrm{DD^*P}}\mathrm{D}_b(\partial_\mu\phi)_{ba}\mathrm{D}_a^{*\mu\dagger} + g_{\mathrm{DD^*P}}\mathrm{D}_b^{*\mu}(\partial_\mu\phi)_{ba}\mathrm{D}_a^\dagger \\ & + g_{\bar{\mathrm{D}}\bar{\mathrm{D}}^*\mathrm{P}}\bar{\mathrm{D}}_a^{*\mu\dagger}(\partial_\mu\phi)_{ab}\bar{\mathrm{D}}_b + g_{\bar{\mathrm{D}}\bar{\mathrm{D}}^*\mathrm{P}}\bar{\mathrm{D}}_a^\dagger(\partial_\mu\phi)_{ab}\bar{\mathrm{D}}_b^{*\mu}, \end{aligned} \tag{9}$$

with $D^{(*)} = (D^{(*)+}, D^{(*)0}, D_s^{(*)+})$ and $\bar{D}^{(*)} = (D^{(*)-}, \bar{D}^{(*)0}, D_s^{(*)-})$. ϕ is the 3×3 Hermitian matrix for the octet of Goldstone bosons. In the chiral and heavy quark limit, the above coupling constants are

$$g_{DD^*P} = -g_{\bar{D}\bar{D}^*P} = -\frac{2g}{f_\pi}\sqrt{m_D m_{D^*}}, \tag{10}$$

$$g_{D^*D_1P} = g_{\bar{D}^*\bar{D}_1P} = -\frac{\sqrt{6}}{3}\frac{h'}{\Lambda_\chi f_\pi}\sqrt{m_{D^*}m_{D_1}}, \tag{11}$$

with the chiral symmetry breaking scale $\Lambda_\chi \simeq 1$ GeV and the coupling $h' = 0.65$ [75].

By assuming $Y(4260)$ is a $D_1\bar{D}$ molecular state, the effective Lagrangian is constructed as

$$\mathcal{L}_{Y(4260)D_1D} = i\frac{y}{\sqrt{2}}(\bar{D}_a^\dagger Y^\mu D_{1a}^{\mu\dagger} - \bar{D}_{1a}^{\mu\dagger} Y^\mu D_a^\dagger) + \text{H.c.}, \tag{12}$$

which is an S-wave coupling. Since the mass $Y(4260)$ is slightly below an S-wave $D_1\bar{D}$ threshold, the effective coupling $g_{Y(4260)D_1D}$ is related to the probability of finding the D_1D component in the physical wave function of the bound state, c^2, and the binding energy, $\delta E = m_D + m_{D_1} - m_Y$ [22,76,77],

$$g_{\rm NR}^2 \equiv 16\pi(m_D + m_{D_1})^2 c^2\sqrt{\frac{2\delta E}{\mu}}[1 + \mathcal{O}(\sqrt{2\mu\epsilon r})], \tag{13}$$

where $\mu = m_D m_{D_1}/(m_D + m_{D_1})$ and r is the reduced mass and the range of the forces. The coupling constants in Eq. (12) are given by the first term in the above equation. The coupling constant gets maximized for a pure bound state, which corresponds to $c^2 = 1$ by definition. In the following, we present the numerical results with $c^2 = 1$.

With the mass $m_Y = 4263^{+8}_{-9}$ MeV, and the averaged masses of the D and D_1 mesons [9], we obtain the mass differences between the $Y(4260)$ and their corresponding thresholds,

$$m_D + m_{D_1} - m_Y = 27^{+9}_{-8}\ \text{MeV}, \tag{14}$$

and with $c^2 = 1$, we obtain

$$|y| = 14.62^{+1.11}_{-1.25} \pm 6.20\ \text{GeV}, \tag{15}$$

where the first errors are from the uncertainties of the binding energies, and the second ones are due to the approximate nature of Eq. (13).

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:

$$M_{fi} = \int \frac{d^4 q_2}{(2\pi)^4}\sum_{D^*\text{pol}}\frac{T_1T_2T_3}{a_1a_2a_3}\mathcal{F}(m_2, q_2^2), \tag{16}$$

where T_i and $a_i = q_i^2 - m_i^2$ ($i = 1, 2, 3$) are the vertex functions and the denominators of the intermediate meson propagators. For example, in Fig. 2(a), T_i ($i = 1, 2, 3$) are

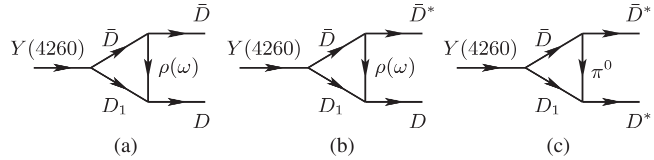

FIG. 1. The hadron-level diagrams for $Y(4260) \to D^{(*)}\bar{D}^{(*)}$ with $D_1\bar{D}$ as the intermediate states.

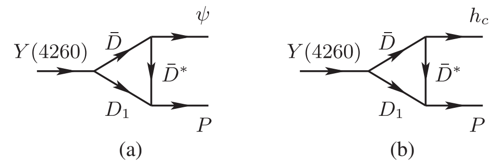

FIG. 2. The hadron-level diagrams for hidden-charm decays of $Y(4260)$ with $D_1\bar{D}$ as intermediate states. P denotes the pseudoscalar meson π^0 or η.

the vertex functions for the initial $Y(4260)$, final charmonium, and final light pseudoscalar mesons, respectively. a_i ($i = 1, 2, 3$) are the denominators for the intermediate $\bar{D}$, D^*, and D_1 mesons, respectively. We introduce a dipole form factor,

$$\mathcal{F}(m_2, q_2^2) \equiv \left(\frac{\Lambda^2 - m_2^2}{\Lambda^2 - q_2^2}\right)^2, \tag{17}$$

where $\Lambda \equiv m_2 + \alpha\Lambda_{\rm QCD}$ and the QCD energy scale $\Lambda_{\rm QCD} = 220$ MeV. This form factor is supposed to kill the divergence, compensate the off-shell effects arising from the intermediate exchanged particle and the nonlocal effects of the vertex functions [36,78,79].

III. NUMERICAL RESULTS

Since $Y(4260)$ has a large width 95 ± 14 MeV, one has to take into account the mass distribution of the $Y(4260)$ when calculating its decay widths. Its two-body decay width can then be calculated as follows [80]:

$$\Gamma(Y(4260))_{2\text{-body}} = \frac{1}{W}\int_{(m_Y-2\Gamma_Y)^2}^{(m_Y+2\Gamma_Y)^2} ds\frac{(2\pi)^4}{2\sqrt{s}} \times \int d\Phi_2|\mathcal{M}|^2\frac{1}{\pi}\text{Im}\left(\frac{-1}{s - m_Y^2 + im_Y\Gamma_Y}\right), \tag{18}$$

where $\int d\Phi_2$ is the two-body phase space [9]. $\mathcal{M}$ are the loop transition amplitudes for the processes in Figs. 1 and 2. The factor $1/W$ with

$$W = \frac{1}{\pi}\int_{(m_Y-2\Gamma_Y)^2}^{(m_Y+2\Gamma_Y)^2}\text{Im}\left(\frac{-1}{s - m_Y^2 + im_Y\Gamma_Y}\right)ds \tag{19}$$

is considered in order to normalize the spectral function of the $Y(4260)$ state.

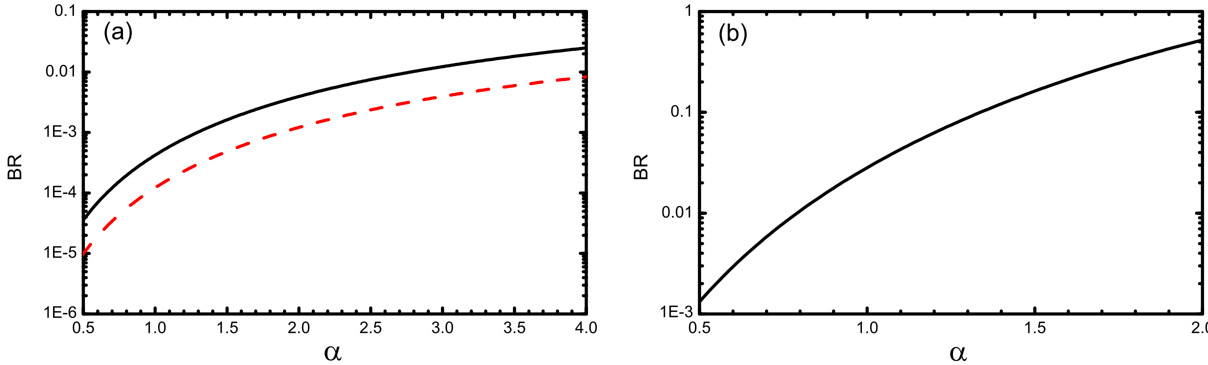

FIG. 3 (color online). (a) The α dependence of the branching ratios of $Y(4260) \to D\bar{D}$ (solid line) and $D^*\bar{D} + c.c.$ (dashed line). (b) The α dependence of the branching ratios of $Y(4260) \to D^*\bar{D}^*$.

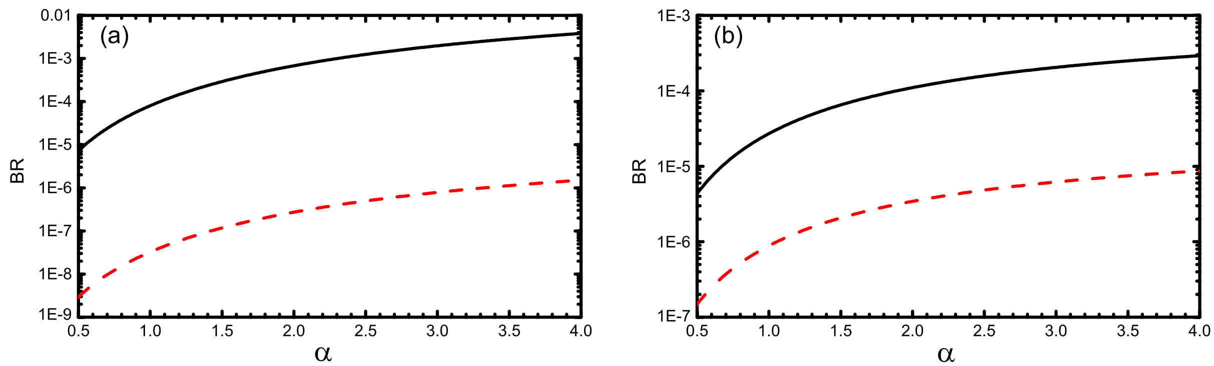

FIG. 4 (color online). (a) The α dependence of the branching ratios of $Y(4260) \to J/\psi\eta$ (solid line) and $J/\psi\pi^0$ (dashed line). (b) The α dependence of the branching ratios of $Y(4260) \to \psi'\eta$ (solid line) and $\psi'\pi^0$ (dashed line).

The numerical results are presented in Figs. 3–5. In Table I, we list the predicted branching ratios of $Y(4260)$ at different α values, and the errors are from the uncertainties of the coupling constants in Eq. (15). We have checked that including the width for D_1 only causes a minor change of about 1%–3%.

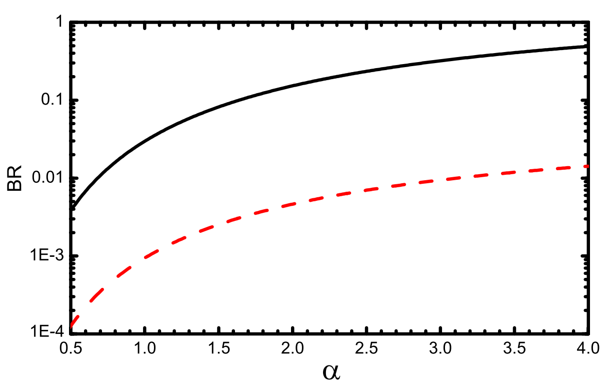

FIG. 5 (color online). The α dependence of the branching ratios of $Y(4260) \to h_c\eta$ (solid line) and $h_c\pi^0$ (dashed line).

In Fig. 3(a), we plot the α dependence of the branching ratios of $Y(4260) \to D\bar{D}$ (solid line) and $Y(4260) \to D^*\bar{D} + c.c.$ (dashed line), respectively. The branching ratios of $Y(4260) \to D^*\bar{D}^*$ in terms of α are shown in Fig. 3(b). In this figure, no cusp structure appears. This is because the mass of $Y(4260)$ lies below the intermediate $D_1\bar{D}$ threshold. The α dependence of the branching ratios is not drastically sensitive to some extent, which indicates a reasonable cutoff of the ultraviolet contributions by the empirical form factors. As shown in this figure, at the same α, the intermediate $D_1\bar{D}$ meson loops turns out to be more important in $Y(4260) \to D^*\bar{D}^*$ than that in $Y(4260) \to D\bar{D}$ and $D^*\bar{D} + c.c.$ This behavior can also be seen from Table I. As a result, a smaller value of α is favored in $Y(4260) \to D^*\bar{D}^*$. This phenomenon can easily be explained from Fig. 1. For the decay $Y(4260) \to D^*\bar{D}^*$, the off-shell effects of intermediate mesons $D_1D(\pi)$ are not significant, which makes this decay favor a relatively smaller α value. For the decay $Y(4260) \to D\bar{D}$ and $D^*\bar{D} + c.c.$, since the exchanged mesons of the intermediate meson loops are ρ and ω, which makes their off-shell effects relatively significant, this decay favors a relatively larger α value.

TABLE I. The predicted branching ratios of $Y(4260)$ decays with different α values. The uncertainties are dominated by the use of Eq. (13).

Final states	$\alpha = 0.5$	$\alpha = 1.0$	$\alpha = 1.5$	$\alpha = 2.0$
$D\bar{D}$	$(3.54^{+3.71}_{-2.34}) \times 10^{-5}$	$(4.21^{+4.41}_{-2.78}) \times 10^{-4}$	$(1.62^{+1.70}_{-1.07}) \times 10^{-3}$	$(3.94^{+4.13}_{-2.60}) \times 10^{-3}$
$D^*\bar{D}$ + *c.c.*	$(9.86^{+10.33}_{-6.51}) \times 10^{-6}$	$(1.22^{+1.28}_{-0.80}) \times 10^{-4}$	$(4.82^{+5.05}_{-3.18}) \times 10^{-4}$	$(1.20^{+1.28}_{-0.79}) \times 10^{-3}$
$D^*\bar{D}^*$	$(1.41^{+1.48}_{-0.93}) \times 10^{-3}$	$(2.78^{+2.91}_{-1.83}) \times 10^{-2}$	$(16.24^{+17.01}_{-10.72})\%$	$(52.21^{+54.69}_{-34.48})\%$
$J/\psi\eta$	$(7.43^{+7.78}_{-4.91}) \times 10^{-6}$	$(8.19^{+8.58}_{-5.41}) \times 10^{-5}$	$(2.95^{+3.09}_{-1.95}) \times 10^{-4}$	$(6.80^{+7.12}_{-4.49}) \times 10^{-4}$
$J/\psi\pi^0$	$(3.04^{+3.18}_{-2.01}) \times 10^{-9}$	$(3.32^{+3.48}_{-2.19}) \times 10^{-8}$	$(1.19^{+1.24}_{-0.78}) \times 10^{-7}$	$(2.72^{+2.85}_{-1.79}) \times 10^{-7}$
$\psi'\eta$	$(4.34^{+4.54}_{-2.84}) \times 10^{-6}$	$(2.71^{+2.84}_{-1.79}) \times 10^{-5}$	$(6.50^{+6.81}_{-4.29}) \times 10^{-5}$	$(1.10^{+1.15}_{-0.73}) \times 10^{-4}$
$\psi'\pi^0$	$(1.76^{+1.84}_{-1.16}) \times 10^{-7}$	$(9.71^{+10.17}_{-6.41}) \times 10^{-7}$	$(2.14^{+2.24}_{-1.41}) \times 10^{-6}$	$(3.43^{+3.59}_{-2.26}) \times 10^{-6}$
$h_c\eta$	$(3.87^{+4.05}_{-2.55}) \times 10^{-3}$	$(2.99^{+3.13}_{-1.97}) \times 10^{-2}$	$(8.20^{+8.59}_{-5.41}) \times 10^{-2}$	$(15.26^{+15.98}_{-10.08})\%$
$h_c\pi^0$	$(1.27^{+1.33}_{-0.84}) \times 10^{-4}$	$(9.50^{+9.95}_{-6.27}) \times 10^{-4}$	$(2.54^{+2.66}_{-1.67}) \times 10^{-3}$	$(4.62^{+4.83}_{-3.05}) \times 10^{-3}$

In a fit to the total hadronic cross sections measured by BES [81], authors set an upper limit on $\Gamma_{e^+e^-}$ for $Y(4260)$ to be less than 580 eV at 90% confidence level (C.L.) [82]. This implies that its branching fraction to $J/\psi\pi^+\pi^-$ is greater than 0.6% at 90% C.L. [82]. Recently, BESIII has reported a study of $e^+e^- \to h_c\pi^+\pi^-$, and observes a state with a mass of $4021.8 \pm 1.0 \pm 2.5$ MeV and a width of $5.7 \pm 3.4 \pm 1.1$ MeV in the $h_c\pi^\pm$ mass distribution, called the $Z_c(4020)$ [83]. The Belle Collaboration did a comprehensive search for $Y(4260)$ decays to all possible final states containing open charmed meson pairs and found no sign of a $Y(4260)$ signal in any of them [84–89]. The *BABAR* Collaboration measured some upper limits of the ratios $\mathcal{B}(Y(4260) \to D\bar{D})/\mathcal{B}(Y(4260) \to J/\psi\pi^+\pi^-) < 7.6$ at 95% C.L. [90], $\mathcal{B}(Y(4260) \to D^*\bar{D})/\mathcal{B}(Y(4260) \to J/\psi\pi^+\pi^-) < 34$, and $\mathcal{B}(Y(4260) \to D^*\bar{D}^*)/\mathcal{B}(Y(4260) \to J/\psi\pi^+\pi^-) < 40$ at 90% C.L. [91], respectively. Within the parameter range considered in this work, the results displayed in Table I could be compatible with these available experimental limits. However, since there are still several uncertainties coming from the undetermined coupling constants, and the cutoff energy dependence of the amplitude is not quite stable, the numerical results would be lacking in high accuracy. Especially, since the kinematics, off-shell effects arising from the exchanged particles and the divergence of the loops in these open charmed channels studied here are different, the cutoff parameter can also be different in different decay channels. We expect more precise experimental measurements on these open charmed pairs to test this point in the near future.

In Ref. [59], a nonrelativistic effective field theory method was introduced to study the meson loop effects in $\psi' \to J/\psi\pi^0$ transitions. And a power counting scheme was proposed to estimate the contribution of the loop effects, which is helpful to judge how important the coupled-channel effects are. This power counting scheme was analyzed in detail in Ref. [61]. Recently, the authors studied that the S-wave threshold plays a more important role than the P-wave, especially for the S-wave molecule with large coupling to its components, such as $Y(4260)$ coupling to $D_1\bar{D}$ in Ref. [22]. Before giving the explicit numerical results, we will follow the similar power counting scheme to qualitatively estimate the contributions of the coupled-channel effects discussed in this work. Corresponding to the diagrams in Figs. 2(a) and 2(b), the amplitudes for $Y(4260) \to J/\psi\pi^0$ ($J/\psi\eta$, $\psi'\pi^0$, $\psi'\eta$) and $Y(4260) \to h_c\pi^0$ ($h_c\eta$) scale as

$$\frac{v^5}{(v^2)^3} q^3 \frac{\Delta}{v^2} \sim \frac{q^3\Delta}{v^3} \tag{20}$$

and

$$\frac{v^5}{(v^2)^3} q^2 \frac{\Delta}{v^2} \sim \frac{q^2\Delta}{v^3}, \tag{21}$$

respectively. There are two scaling parameters v and q appearing in the above two formulas. As illustrated in Ref. [92], v is understood as the average velocity of the intermediate charmed meson. q denotes the momentum of the outgoing pseudoscalar meson. And Δ denotes the charmed meson mass difference, which is introduced to account for the isospin or SU(3) symmetry violation. For the π^0 and η production processes, the factors Δ are about $M_{D^+} + M_{D^-} - 2M_{D^0}$ and $M_{D^+} + M_{D^0} - 2M_{D_s}$, respectively. According to Eqs. (20) and (21), it can be concluded that the contributions of the coupled channel effects would be significant here since the amplitudes scale as $\mathcal{O}(1/v^3)$. And the branching ratio of $Y(4260) \to h_c\pi^0$ is expected to be larger than that of $Y(4260) \to J/\psi\pi^0$, because the corresponding amplitudes scale as $\mathcal{O}(q^2)$ and $\mathcal{O}(q^3)$, respectively. However, the momentum q in $Y(4260) \to J/\psi\pi^0$ is larger than that in $Y(4260) \to h_c\pi^0$, which may compensate for this discrepancy to some extent.

For the open charmed decays in Fig. 1, the exchanged intermediate mesons are light vector mesons or light pseudoscalar mesons, which will introduce different scales. Since we cannot separate different scales, we just give possible numerical results in the form factor scheme.

For the hidden-charm transitions $Y(4260) \to J/\psi \eta(\pi^0)$, we plot the α dependence of the branching ratios of $Y(4260) \to J/\psi \eta(\pi^0)$ in Fig. 4(a) as shown by the solid and dashed lines, respectively. The π^0-η mixing has been taken into account. (Using Dashen's theorem [93], one may express the mixing angle in terms of the masses of the Goldstone bosons at leading order in chiral perturbation theory, and the value is about 0.01.) Some points can be learned from this figure: (1) A predominant feature is that the branching ratios are not drastically sensitive to the cutoff parameter, which indicates a reasonable cutoff of the ultraviolet contributions by the empirical form factors to some extent. (2) The leading contributions to $Y(4260) \to J/\psi \pi^0$ are given by the differences between the neutral and charged charmed meson loops and also from the π^0-η mixing through the loops contributing to the eta transition. (3) At the same α, the branching ratios for the $Y(4260) \to J/\psi \eta$ transition are 2–3 orders of magnitude larger than that of $Y(4260) \to J/\psi \pi^0$. It is because there are no cancellations between the charged and neutral meson loops.

The branching ratios of $Y(4260) \to \psi' \eta$ (solid line) and $Y(4260) \to \psi' \pi^0$ (dashed line) in terms of α are shown in Fig. 4(b). The behavior is similar to that of Fig. 4(a). Since the mass of ψ' is closer to the thresholds of $\bar{D}D^*$ than J/ψ, it should give rise to more important threshold effects in $Y(4260) \to \psi' \eta(\pi^0)$ than in $Y(4260) \to J/\psi \eta(\pi^0)$. At the same α value, the obtained branching ratio of $Y(4260) \to \psi' \pi^0$ is larger than that of $Y(4260) \to J/\psi \pi^0$. The three-momentum of final η is only about 167 MeV in $Y(4260) \to \psi' \eta$, which leads to a smaller branching ratio in $Y(4260) \to J/\psi \eta$ than in $Y(4260) \to J/\psi \eta$ at the same α value.

In Fig. 5, we plot the α dependence of the branching ratios of $Y(4260) \to h_c \pi^0$ (solid line) and $Y(4260) \to h_c \eta$ (dashed line), respectively. The branching ratios for $Y(4260) \to h_c \pi^0(\eta)$ are larger than that of $Y(4260) \to J/\psi \pi^0(\eta)$ and $\psi' \pi^0(\eta)$, which is consistent with the power counting analysis in Eqs. (20) and (21).

To study the exclusive threshold effects via the intermediate meson loops, we define the following ratios:

$$R_1 \equiv \frac{|\mathcal{M}_{Y(4260)\to\psi'\pi}|^2}{|\mathcal{M}_{Y(4260)\to J/\psi\pi}|^2}, \quad R_2 \equiv \frac{|\mathcal{M}_{Y(4260)\to\psi'\eta}|^2}{|\mathcal{M}_{Y(4260)\to J/\psi\eta}|^2}, \tag{22}$$

and

$$r_1 \equiv \frac{|\mathcal{M}_{Y(4260)\to J/\psi\pi}|^2}{|\mathcal{M}_{Y(4260)\to J/\psi\eta}|^2}, \qquad r_2 \equiv \frac{|\mathcal{M}_{Y(4260)\to\psi'\pi}|^2}{|\mathcal{M}_{Y(4260)\to\psi'\eta}|^2},$$
$$r_3 \equiv \frac{|\mathcal{M}_{Y(4260)\to h_c\pi}|^2}{|\mathcal{M}_{Y(4260)\to h_c\eta}|^2}. \tag{23}$$

These ratios are plotted in Figs. 6(a) and 6(b), respectively. The stabilities of the ratios in terms of α indicate a reasonably controlled cutoff for each channel by the form factor. Since the coupling vertices are the same for those decay channels when taking the ratio, the stability of the ratios suggests that the transitions of $Y(4260) \to J/\psi \pi^0(\eta)$ and $\psi' \pi^0(\eta)$ are largely driven by the open threshold effects via the intermediate $D_1\bar{D}$ meson loops to some extent. The future experimental measurements of these decays can help us investigate this issue deeply.

IV. SUMMARY

In this work, we have investigated the hidden-charm decays of $Y(4260)$ and the decays $Y(4260) \to D\bar{D}$, $D\bar{D}^*$, and $D^*\bar{D}^*$ in ELA. In this calculation, $Y(4260)$ is assumed to be the $D_1\bar{D}$ molecular state. Our results show that the α dependence of the branching ratios is not drastically sensitive, which indicates the dominant mechanism driven by the intermediate meson loops with a fairly good control of the ultraviolet contributions.

For the hidden charmonium decays, we also carried out the power counting analysis, and our results for these decays in ELA are qualitatively consistent with the power counting analysis. For the open charmed decays $Y(4260) \to D\bar{D}$, $D\bar{D}^*$, and $D^*\bar{D}^*$, the exchanged intermediate mesons are light vector mesons or light pseudoscalar mesons, which will introduce different

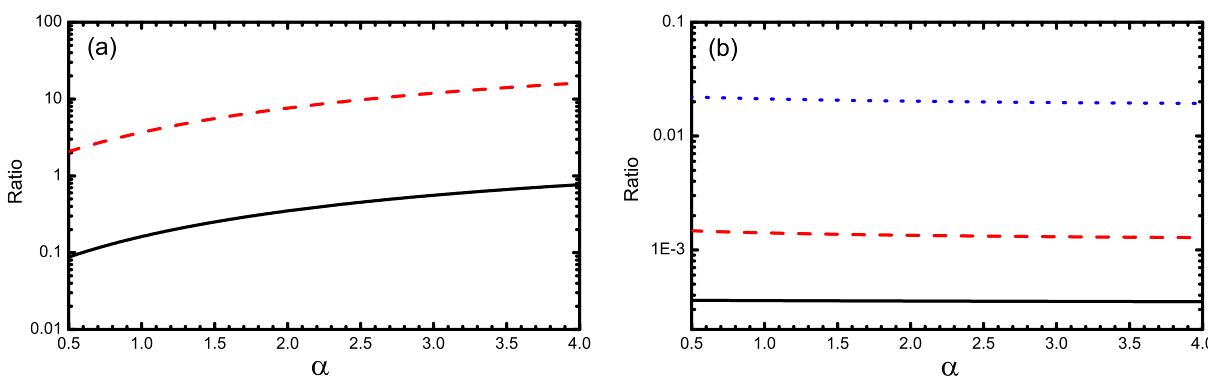

FIG. 6 (color online). (a) The α dependence of the ratios of R_1 (solid line) and R_2 (dashed line) defined in Eq. (22). (b) The α dependence of the ratios of r_1 (solid line), r_2 (dashed line), and r_3 (dotted line) defined in Eq. (23).

scales, so we cannot separate different scales and only give possible numerical results in the form factor scheme. For the decay $Y(4260) \to D^*\bar{D}^*$, the exchanged meson π is almost on-shell, so the coupled channel effects are more important than other channels studied here. We expect the experiments to search for the hidden-charm and charmed meson pair decays of $Y(4260)$, which will help us investigate the nature and decay mechanisms of $Y(4260)$ deeply.

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