New structure around 3250 MeV in the baryonic $B$ decay and the $D_0^*(2400)N$ molecular hadron

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Abstract In this work, we first propose the isovector $nD_0^*(2400)^0$ molecular state to explain the enhancement structure around 3250 MeV ($X_c(3250)^0$) in the $\Sigma_c^{++}\pi^-\pi^-$ invariant mass spectrum newly observed by the BaBar Collaboration. Under this molecular state configuration, both the analysis of the mass spectrum and the study of its dominant decay channel can well shed light on its resonance parameters measured by BaBar. Our investigation also shows that the isovector $nD_0^*(2400)^0$ molecular state can decay into $\Sigma_c^{++}\pi^-\pi^-$, which is consistent with experimental observation. These studies provide direct support to the isovector $nD_0^*(2400)^0$ molecular state assignment to $X_c(3250)^0$.

The BaBar Collaboration recently reported a new enhancement structure in the $\Sigma_c^{++}\pi^-\pi^-$ invariant mass spectrum of the $B^-\rightarrow \Sigma_c^{++}\bar{\nu}\pi^-\pi^-$ decay. The mass and width are given by $M = 3245 \pm 20 \text{ MeV}$ and $\Gamma = 108 \pm 6 \text{ MeV}$, respectively [1]. We will call the new structure $X_c(3250)^0$ in this work. In terms of the decay channel observed, we conclude that $X_c(3250)^0$ is an isosritplet with charm number $C = +1$. It is an intriguing research topic to understand this new enhancement structure observed in the baryonic $B$ decay. In this work, we propose a novel approach to explain the BaBar’s observation of $X_c(3250)^0$, which can be naturally explained as a molecular hadron composed of a charmed meson $D_0^*(2400)$ and a nucleon $N$. In the following, we illustrate why the explanation of the $D_0^*(2400)N$ molecular hadron is reasonable for the observed $X_c(3250)^0$ in detail.

A molecular explanation requires that the mass of the observed state should be below and close to the sum of masses of its components. The $X_c(3250)$ newly observed by BaBar just meets this necessary condition because the mass is near the threshold of $D_0^*(2400)$ and $N$. For the $D_0^*(2400)N$ molecular system, the flavor wave function is written as

$$ I = 1 : \begin{cases} |X_c(3250)^0\rangle = |D_0^*(2400)^0n\rangle \\ |X_c(3250)^+\rangle = \frac{|D_0^*(2400)^0n\rangle - |D_0^*(2400)^0p\rangle}{\sqrt{2}} \\ |X_c(3250)^{++}\rangle = |D_0^*(2400)^+p\rangle \end{cases}, \quad (1) $$

$$ I = 0 : |Y_c(3250)^+\rangle = \frac{|D_0^*(2400)^+n\rangle + |D_0^*(2400)^0p\rangle}{\sqrt{2}}, \quad (2) $$

whose expressions correspond to isrotplet and isosinglet, respectively, where $|X_c(3250)^0\rangle$ is the flavor wave function of $X_c(3250)^0$, etc. Under the assignment of the $D_0^*(2400)N$ molecular state, we further deduce the quantum number of $X_c(3250)$ as $I(J^P) = 1(\frac{1}{2}^+)$. With the neutral $X_c(3250)^0$ as an example, we first carry out the calculation of the binding energy for the $D_0^*(2400)N$ molecular system. According to the mass values of the $D_0^*(2400)^0$ and neutron $n$ listed in Particle Data Group (PDG) [2], we obtain the binding energy of $X_c(3250)^0$ as $\sim -13 \text{ MeV}$, with $M_{D_0^*(2400)^0} = 2318 \text{ MeV}$ and $M_n = 940 \text{ MeV}$. Then we have $M_{D_0^*(2400)^0 + n} = 3258 \text{ MeV} > M_{X_c(3250)^0}$. In the following, we will examine whether it is reasonable to explain $X_c(3250)^0$ as the $nD_0^*(2400)^0$ molecular hadron. To deduce the effective potential between $D_0^*(2400)^0$ and $n$, we adopt the effective Lagrangian of the light mesons interacting with the charmed meson $D_0^*(2400)$ or nucleon, i.e.,

$$ \mathcal{L}_{mD_0^*D_0^*} = -i \frac{\beta^g}{\sqrt{2}} D_0^* D_0^*(2i v \cdot \nabla_{\bar{u}}) + 2 g^g \sigma D_0^* D_0^* a \cdot \nabla_{\bar{u}}. $$

(3)
$\mathcal{L}_{mNN} = -\sqrt{2}g_{\sigma NN}N_b\left(\gamma^\mu + \frac{k\sigma^{\mu\nu}}{2m_N} \partial^\nu\right)\gamma_\mu \gamma_\nu N_a$

\[ + g_{\sigma NN}N_b N \tag{4} \]

with $\nu = (1, 0), \sigma^{\mu\nu} = \frac{i}{2}(\gamma^\mu \gamma^\nu - \gamma^\nu \gamma^\mu)$ and the vector matrix $\mathbb{V}$

\[ \mathbb{V} = \left( \begin{array}{ccc} \frac{1}{\sqrt{2}}\rho^0 + \frac{1}{\sqrt{2}}\omega \\
 \frac{1}{\sqrt{2}}\rho^0 + \frac{1}{\sqrt{2}}\omega \\
 \end{array} \right) \tag{5} \]

The mass of the exchanged sigma is taken as 660 MeV.

The effective potential of $X_c(3250)^0$ in the coordinate space is given by

\[ \mathcal{V}_{Total}(r) = \frac{1}{2} \mathcal{V}_\rho(r) + \frac{1}{2} \mathcal{V}_{\sigma}(r) \tag{6} \]

with $\mathcal{V}_\rho(r) = -2\beta'^2 g_{\rho NN}g_{\rho N} \rho(A, m_\rho, r)$ and $\mathcal{V}_{\sigma}(r) = -2g_{\rho NN}g_{\rho N} Y(A, m_\sigma, r)$, where the $Y(A, m, r)$ function is defined as

\[ Y(A, m, r) = \frac{1}{4\pi r} \left[ e^{-mr} - e^{-Ar} - \frac{\beta^2 - m^2}{2A} r e^{-Ar} \right]. \]

By solving the Schrödinger equation with the effective potential obtained in Eq. (5), the dependence on the cutoff $A$ of the bound state solution for $X_c(3250)^0$ is shown in Fig. 1. When $A = 1.23$ GeV, the theoretical result of the binding energy for $X_c(3250)^0$ is consistent with the experimental data [1]. Usually, in the one-boson-exchange model the general criterion of forming a molecular state is that we obtain a negative binding energy of this system and the corresponding cutoff $A$ should be around 1 GeV. Thus, by our calculation, we conclude that $D_0^*(2400)$ and nucleon can form a loosely bound state with a small binding energy since the adopted $A$ value is close to 1 GeV.

As shown in Fig. 1, the potential of the $D_0(2400)^0$ h molecular state is mainly from $\rho$ exchange, where the $\pi$ exchange is absent. It is well known that the vector meson exchanges provide short-range interaction. Usually, the necessary condition of forming a bound state requires that the mass of the exchanged meson is larger than the widths of the components of this bound state. We notice that the $\rho$ meson mass ($m_\rho = 770$ MeV) is larger than the width of $D_0(2400)^0$ ($\Gamma \sim 267$ MeV), which basically satisfies the necessary condition of forming a molecular state. Considering the above reasons, $D_0(2400)^0$ and $n$ can interact with each other before $D_0(2400)^0$ decays into other final states.

This is why we still expect to discuss the molecular configuration to $X_c(3250)^0$ although $D_0(2400)^0$ is a very broad meson.

Since the total width of a state is mainly determined by its dominant decay, studying the dominant decay mode of $X_c(3250)^0$ provides important information of its total decay width. What is more important is that this study can be used as a critical test of the $nD_0^*(2400)$ molecular state assignment to $X_c(3250)^0$.

As the $nD_0^*(2400)$ molecular state, $X_c(3250)^0$ decays into $D_0^{+,0}\pi^{-0}n$. Here, $X_c(3250)^0$ first falls apart into $n$ and $D_0^*(2400)^0$. Then, decays $X_c(3250)^0 \to D_0^{+,0}\pi^{-0}n$ occur via the intermediate $D_0^*(2400)^0$ as shown in Eq. (7). Since the branching ratio of $D_0^*(2400)^0 \to D \pi$ is almost 100 %, $X_c(3250)^0 \to D_0^{+,0}\pi^{-0}n$ is the dominant mode. For this process, the differential decay width is written as

\[ d\Gamma = \frac{1}{2J + 1} \frac{1}{2E} (2\pi)^4 \sum_\lambda |\mathcal{M}|^2 \times \delta^4 \left( \sum_{i=1}^3 q_i - P \right) \frac{d^3q_1}{(2\pi)^32e_1} \frac{d^3q_2}{(2\pi)^32e_2} \frac{d^3q_3}{(2\pi)^32e_3}, \tag{7} \]

with $J = 1/2$, where $q_i(e_i)$ ($i = 1, 2, 3$) denotes the momentum (energy) of final states. The decay amplitude $\mathcal{M}$ is ex-
pressed as
\[
\mathcal{M} = \frac{A(X_c(3250)^0 \rightarrow D_0^*(n)\bar{D}^*_0 \rightarrow D\pi)}{q^2 - M^2_{D_0^*} + iM_{D_0^*} \Gamma_{D_0^*}},
\]
(8)
where \(A_{X_c(3250)0} = A(X_c(3250)^0 \rightarrow D_0^*(n))\) and \(A_{D_0^*} = A(D_0^*(n) \rightarrow D\pi)\) describe the interactions \(X_c(3250)^0 \rightarrow nD_0^*(2400)^0\) and \(D_0^*(2400)^0 \rightarrow D\pi\), respectively.

By the covariant spectator theory (CST), we can describe the collapse of \(X_c(3250)^0\) into the on-shell \(n\) and the \(D_0^*(2400)\), where the vertex function \(|\Gamma\rangle\) satisfies the relation
\[
|\Gamma\rangle = VG(|\Gamma\rangle),
\]
(9)
which is obtained by the Gross equation. In Fig. 2, we present the diagrammatic representation of the Gross equation for the vertex function \(\Gamma\). We further obtain the Gross equation, which is dependent on three-momentum \(p\) as
\[
\Gamma(p) = \int \frac{d^3k}{(2\pi)^3} \mathcal{V}(p, k, W)G(k, W)\Gamma(k)
\]
(10)
with \(\Gamma(p) = \hat{\Gamma}(p)u_R(W)\). We use the convention \(\bar{u}u = 2m\) here and hereafter, where we perform the integration over \(k^0\). \(P = (W, 0)\) denotes the four-momentum of the \(nD_0^*(2400)^0\) system. \(p = (p_1 - p_2)/2\) and \(k = (k_1 - k_2)/2\) are the relative momenta as depicted in Fig. 2. Just because neutron is on the mass shell, one gets \(k = (k_0, k)\), \(p = (p_0, p)\), \(k_0 = E_1(k) - \frac{1}{2} W\), \(E_1(k) = \sqrt{m_1^2 + k^2}\). \(\mathcal{V}(p, k, W)\) is the interaction kernel with neutron on the mass shell. The two-body Green function \(G(k, W)\) in Eq. (10) is expressed as
\[
G(k, W) = \frac{1}{2E_1(k)}\bar{u}^{(r)}(k)\psi^{(r)}(k) \frac{1}{2E_2(k)}
\]
\[
\times \left( \frac{1}{2E_1(k) - E_1(k) + W} + \frac{1}{2E_1(k) + E_1(k) - W} \right).
\]
(11)

The normalization of the vertex \(\Gamma(p)\) requires [9, 10]
\[
1 = \int \frac{d^3p}{(2\pi)^3} \Gamma^\dagger(p) \frac{\partial}{\partial W^2} \left[ G(p, W) \Gamma(p) \right].
\]
(12)

The wave functions for the \(nD_0^*(2400)^0\) bound state are
\[
\psi^+(p) = \frac{1}{\sqrt{(2\pi)^32W}} \left[ \frac{1}{2E_1(p)2E_2(p)} \times \bar{u}^{(r)}(p)\tilde{\mathcal{F}}(p)u_R(W) \right]
\]
(13)
\[
\psi^-(p) = \frac{1}{\sqrt{(2\pi)^32W}} \left[ \frac{1}{2E_1(p)2E_2(p)} \times \bar{u}^{(r)}(p)\tilde{\mathcal{F}}(p)u_R(W) \right]
\]
(14)
where \(\mathcal{V}(p, k, W)\) is
\[
V_{rr'}(p, k, W) = -\sqrt{\frac{1}{2E_1(p)} \frac{1}{2E_2(p)} \frac{1}{2E_2(k)} \frac{1}{2E_2(k)}} \bar{u}^{(r)}(p)\psi^{(r)}(k),
\]
(15)

After taking the nonrelativistic approximation [11–13] and the Fourier transformation, we get the integral equations in the coordinate space,
\[
\left( \frac{\nabla^2}{\mu} + \epsilon \right) \psi^+(r) = -\left[ V(r) + V(r) F(r) V(r) \right] \psi^+(r),
\]
(16)
\[
\psi^-(r) = -\left[ F(r) V(r) \right] \psi^+(r),
\]
(17)
where \(F(r) = [m_2 - m_1 + W + V(r)]^{-1}\) and \(E_1(p) + E_2(p) - W \approx -\epsilon + \frac{p^2}{\mu}\) with the reduced mass \(\mu\) and the binding energy \(\epsilon = W - m_1 - m_2\).
For the loosely bound system discussed in this work, it is reasonable to assume $m_{1,2} \gg (\langle V \rangle).$ Then, we have

$$\left(-\left(\frac{\nabla^2}{\mu} + \epsilon\right)\psi^+(r) = -V(r)\psi^+(r), \quad (18)\right.$$ \hspace{1cm} \psi^-(r) = 0, \quad (19)$$

where Eq. (18) corresponds to the Schrödinger equation. We need to specify that $V(r)$ is the total effective potential of the $nD^*_0(2400)$ molecular state shown in Eq. (6). Thus, by solving the Schrödinger equation with the deduced effective potential for the $nD^*_0(2400)$ molecular state, we obtain $\psi^+(r).$ By the Fourier transformation, we get the wave function $\psi^+(p)$ in the momentum space

$$\psi^+(p) = \frac{1}{(2\pi)^{3/2}} \int d^3r e^{-ipr} \psi^+(r), \quad (20)$$

which satisfies the normalization condition $\int d^3p \times |\psi^+(p)|^2 = 1$ required by Eq. (12), where $p$ denotes the relative momentum.

With the calculated $\psi^+(p)$ and Eq. (13), we obtain the vertex $\Gamma(p),$ which directly corresponds to $A_{Xc(3250)}$ in Eq. (8). By the CST, the calculated wave function is related to the vertex of the $Xc(3250)^0$ collapse into $nD^*_0(2400)^0.$ Thus, we find the relation

$$\frac{A_{Xc(3250)}}{q^2 - M^2_{D^*_0} + iM_{D^*_0}\Gamma_{D^*_0}} = \sqrt{2W(2\pi)^3} \frac{\sqrt{2E_n(q)2E_{D^*_0(2400)}}(q)\psi^+_c(q)}{-W + E_n(q) - E_{D^*_0(2400)}(q)}. \quad (21)$$

In addition, the amplitude $A_{D^*_0}$ reads

$$A_{D^*_0(2400)} = i\sqrt{\frac{M_{D^*_0}}{M^2_{D^*_0} M_D}}, \quad (22)$$

where the isospin factor $I$ is taken as 1 and $1/\sqrt{2}$ for decays $D^*_0(2400) \rightarrow D^+\pi^-$ and $D^*_0(2400) \rightarrow D^0\pi^0,$ respectively. The coupling constant $g_{D^*_0}$ is determined by the decay width ($\Gamma_{D^*_0} = 267 \pm 40$ MeV [2]), i.e.,

$$\Gamma(D^0_{D^*_0(2400)} \rightarrow D^+\pi^-) = \frac{M_{D^+} |p|}{8\pi M_{D^*_0}} g_{D^*_0}^2 \quad (23)$$

with $p$ being the three-momentum of the daughter meson in the rest frame of the $D^0_{D^*_0(2400)}$ meson, where the $D\pi$ channel contributes most to the total width of $D^*_0(2400)$ [14, 15].

In Table 1, for several typical values of $\Lambda$ we give the decay width of $Xc(3250)^0 \rightarrow nD\pi,$ which is the dominant decay mode of $Xc(3250)^0$ but does not strongly depend on $\Lambda$. We compare the theoretical value of the decay width $Xc(3250)^0 \rightarrow nD\pi$ with the BaBar’s data [1], which shows that the total width of $Xc(3250)^0$ under the assignment of the $nD^*_0(2400)^0$ molecular state is comparable with the BaBar’s measurement [1]. The study of the dominant decay channels of $Xc(3250)^0$ also supports the $nD^*_0(2400)^0$ molecular state explanation for the observed $Xc(3250)^0.$

Apart from calculating the dominant decay width of $Xc(3250)^0,$ we also calculate the corresponding line shape of the pion spectrum of the $Xc(3250)^0 \rightarrow nD^+\pi^-$ process with the help of the CERNLIB program FOWL (see Fig. 3).

Since $Xc(3250)^0$ was observed in the $\Sigma^+_c\pi^+\pi^-\pi^-\pi^-$ invariant mass spectrum, in the following analysis we illustrate how this process can happen under the $nD^*_0(2400)^0$ molecular state assignment. Due to this molecular picture, $Xc(3250)^0$ first disassociate into off-shell $n$ and $D^*_0(2400)^0,$ which then transits into the final states $\Sigma^+_c\pi^+\pi^-\pi^-$ by exchanging the $D^+$ meson (see Fig. 4(a)–(b) for more details). Thus, we can naturally explain why $Xc(3250)^0$ is reported in the $\Sigma^+_c\pi^+\pi^-\pi^-\pi^-$ invariant mass spectrum.

We also qualitatively discuss other possible decay modes of $Xc(3250)^0,$ which are shown in Fig. 4 with the corresponding schematic diagrams. Here, the $Xc(3250)^0$ decays into $D^0n,$ $D^{*0}n,$ $\Lambda^+_c\pi^-$, $\Lambda^+_c\rho^-$, $\Sigma^+_c\pi^-$, $\Sigma^+_c\rho^-$ occur via these triangle diagrams listed in Fig. 4(c)–(h). We notice that these neutrons in the final states of the $Xc(3250)^0 \rightarrow D^0n,$ $D^{*0}n$ decays are neutral. Thus experimental search for these two decay modes of $Xc(3250)^0$ is difficult. Besides decaying into $D^{(*)0}n,$ $Xc(3250)^0$ can decay into a charmed baryon plus a light meson such that pion and $\rho.$ Searching for $Xc(3250)^0$ by these predicted decay channels will be an interesting research topic. Due to the subsequent decays $\Sigma^+_c \rightarrow \Lambda^+_c + \pi$ and $\rho \rightarrow 2\pi,$ finally the processes of $Xc(3250)^0$ into $\Lambda^+_c\rho^-$, $\Sigma^+_c\pi^-$, $\Sigma^+_c\rho^-$ are related to these final states of a $\Lambda^+_c$ baryon plus multipion. Comparing with these processes, $Xc(3250)^0 \rightarrow \Lambda^+_c\pi$ is a typical two-body

| $Xc(3250)^0$ with $I = 1$ | $Yc(3250)^+$ with $I = 0$ |
|--------------------------|--------------------------|
| $\Lambda$ | $E$ | $\Gamma$ | $A$ | $E$ |
| 1.17 | -3 | 121 $\pm$ 18 | 3.30 | -3 |
| 1.20 | -7 | 111 $\pm$ 17 | 3.60 | -6 |
| 1.23 | -13 | 105 $\pm$ 16 | 4.20 | -14 |
| 1.26 | -22 | 100 $\pm$ 15 | 4.80 | -23 |
| 1.30 | -35 | 92 $\pm$ 14 | 5.70 | -35 |

BaBar [1] | -13 | 108 $\pm$ 6 | - | - |
Fig. 3 The Dalitz plot analysis and the pion spectrum for the $X_c(3250)^0 \rightarrow n D^+ \pi^-$ decay (Color figure online)

Fig. 4 The $X_c(3250)^0 \rightarrow \Sigma_c^{++} \pi^-$ decay and other possible decays of $X_c(3250)^0$ (Color figure online)

decay. We suggest an experiment to carry out the search for $X_c(3250)^0$ by its $\Lambda^+ \pi^-$ invariant mass spectrum.

If $X_c(3250)^0$ is a $D_0(2400)^0 n$ molecular state, its dominant decay mode is $D \pi n$. In addition, we also listed its other decay modes (see Fig. 4), which occur via hadronic loop. Thus, these decay modes are suppressed compared with its $D \pi n$ decay mode. The final states of $X_c(3250)^0 \rightarrow D \pi n$ contain neutral neutron, which makes that it is difficult to find its $D \pi n$ decay mode in experiment. It is the reason why experimental firstly reported $X_c(3250)^0$ in $\Sigma^{++} c \pi^- \pi^-$. In this work we also study the possible existence of the isoscalar partner of $X_c(3250)^0$, which is named as $Y_c(3250)^+$ whose flavor wave function is given in Eq. (2). We find that the obtained $\Lambda$ value corresponding to its bound state solution is around 3.3 GeV, which is larger than 1 GeV. Thus, according to the criterion ($\Lambda \sim 1$ GeV), we conclude that it is impossible to form $Y_c(3250)^+$ molecular state.

In this work we also study the possible existence of the isoscalar partner of $X_c(3250)^0$, which is named as $Y_c(3250)^+$ whose flavor wave function is given in Eq. (2). We find that the obtained $\Lambda$ value corresponding to its bound state solution is around 3.3 GeV, which is larger than 1 GeV. Thus, according to the criterion ($\Lambda \sim 1$ GeV), we conclude that it is impossible to form $Y_c(3250)^+$ molecular state.

Now let us draw a brief conclusion. Being stimulated by the recent observation of an enhancement structure $X_c(3250)^0$ in the $\Sigma_c^{++} \pi^- \pi^-$ invariant mass spectrum of $B^- \rightarrow \Sigma_c^{++} \pi^- \pi^-$ [1], we find that $X_c(3250)^0$ can be well explained as the isovector $n D_0^*(2400)^0$ molecular hadron, which is supported both by the analysis of the mass spectrum and by the study of its dominant decay channel. Furthermore, the observed $X_c(3250)^0 \rightarrow \Sigma_c^{++} \pi^- \pi^-$ can be described reasonably under this picture. In addition, we also mention several other possible decay modes of $X_c(3250)^0$, which can be studied in future experiments. We expect the contributions from BaBar, Belle, LHCb, and forthcoming BelleII, SuperB, which are the ideal places to further investigate the observed $X_c(3250)^0$ in the $B$ decay.

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