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Possible new charmed baryon states from $\bar{B}^0 \to p\bar{p}D^0$ decays

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We examine the invariant-mass spectrum of D^0p in $\bar{B}^0 \to p\bar{p}D^0$ decay measured by *BABAR* and find that, through the two-step processes of $\bar{B}^0 \to \mathbf{B}_c^+(\to D^0p)\bar{p}$, where $\mathbf{B_c}$ denotes a charmed baryon state, some of the peaks can be identified with the established $\Sigma_c(2800)^+$, $\Lambda_c(2880)^+$, and $\Lambda_c(2940)^+$. Moreover, to account for the measured spectrum, it is necessary to introduce a new charmed baryon resonance with $(m,\Gamma) = (3212 \pm 20, 167 \pm 34)$ MeV.

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I. INTRODUCTION

The charmed baryon spectroscopy is expected to be an ideal place to study the dynamics of the light quarks in the environment of a heavy quark. Apart from a charm quark, the charmed baryon contains two light quarks as a diquark, which can be either an antitriplet or a sextet in the quark model. Because of various combinations of individual spins and orbital angular momenta between two light quarks as well as between the diquark and the charm quark, charmed baryons exhibit a rich spectrum of states [1–5]. However, so far, only a few of charmed baryons have been experimentally identified [6,7]. The B factories as well as the LHCb are able to play an important role in studying charmed baryons. For example, the charmed baryon states of $\Lambda_c(2665)^+$, $\Lambda_c(2880)^+$, $\Lambda_c(2940)^+$, $\Sigma_c(2800)$, $\Xi_c(2980)$, and $\Xi_c(3080)$ have been observed by CLEO [8], BELLE [9], and *BABAR* [10]. In particular, the decays of $\Lambda_c(2880)^+ \to \Sigma_c(2520)^{0(++)}\pi^{+(-)}$ have been analyzed to assign the quantum number of $\Lambda_c(2880)^+$ to be $J^P = \frac{5}{2}^+$ [11].

Recently, the *BABAR* collaboration has reported that both $m_{p\bar{p}}$ and m_{D^0p} invariant-mass distributions in $\bar{B}^0 \to p\bar{p}D^0$ exhibit two sets of peaking data points in the spectra [12]. In the $m_{p\bar{p}}$ spectrum, the low-mass peaks near the threshold can be recognized as the threshold effect [13]. On the one hand, the exclusion for the range of $m^2_{D^0p} < 9$ GeV^2 improves the knowledge of the threshold effect, a generic feature in three-body baryonic $\bar{B} \to \mathbf{B}\bar{\mathbf{B}}'M$ decays. On the other hand, the peaks in the m_{Dp} spectrum with $m^2_{p\bar{p}} > 5$ GeV^2 seem puzzling as the data points disagree with what is expected from the uniform phase-space model. In this paper, we find that, unlike the peak around the threshold area in the $m_{p\bar{p}}$ spectrum, those in the m_{D^0p} spectrum *cannot* be understood based on the pQCD effect in the creation of the dibaryon [14,15]. We will demonstrate that the charmed-baryon resonances decaying into D^0p fit well to these peaking data points in $\bar{B}^0 \to p\bar{p}D^0$. It turns out that one of the resonant charmed baryon states is experimentally not observed yet but theoretically studied in the literature [1–4].

II. FORMALISM

To explain the data points in the m_{D^0p} invariant mass distribution in $\bar{B}^0 \to p\bar{p}D^0$, we write down the factorizable amplitude induced at the tree process $b \to c\bar{u}d$ [15],

$$\mathcal{A}(\bar{B}^0 \to p\bar{p}D^0) = \frac{G_F}{\sqrt{2}} V_{cb}V_{ud}^* a_2 \langle D^0|(\bar{c}u)_{V-A}|0\rangle \times \langle p\bar{p}|(\bar{d}b)_{V-A}|\bar{B}^0\rangle, \tag{1}$$

where G_F is the Fermi constant, V_{cb} and V_{ud} represent the Cabibbo-Kobayashi-Maskawa matrix elements for the $b \to c\bar{u}d$ transition at the quark level, and $(\bar{q}_1q_2)_{V-A}$ stands for $\bar{q}_1\gamma_\mu(1-\gamma_5)q_2$. Here, a_2 is a parameter that takes into account the nonfactorizable effects. For the D meson production, the matrix element in Eq. (1) is given by $\langle D|\bar{c}\gamma^\mu\gamma_5 u|0\rangle = -if_D p^\mu$. The matrix elements for the $\bar{B}^0 \to p\bar{p}$ transition are parametrized as [15,16]

$$\begin{aligned}\langle p\bar{p}|\bar{d}\gamma_\mu b|\bar{B}^0\rangle &= i\bar{u}[g_1\gamma_\mu + g_2 i\sigma_{\mu\nu}p^\nu + g_3p_\mu + g_4q_\mu \\ &\quad + g_5(p_{\bar{p}} - p_p)_\mu]\gamma_5 v, \\ \langle p\bar{p}|\bar{d}\gamma_\mu\gamma_5 b|\bar{B}^0\rangle &= i\bar{u}[f_1\gamma_\mu + f_2 i\sigma_{\mu\nu}p^\nu + f_3p_\mu \\ &\quad + f_4q_\mu + f_5(p_{\bar{p}} - p_p)_\mu]v,\end{aligned} \tag{2}$$

where $p = p_B - p_p - p_{\bar{p}}$ and $q = p_p + p_{\bar{p}}$ with $p_i (i = B, p, \bar{p})$ representing the momenta of the particles and the momentum dependences of the form factors are given by

$$f_i = \frac{D_{f_i}}{t^3}, \qquad g_i = \frac{D_{g_i}}{t^3}, \tag{3}$$

with $t \equiv m^2_{p\bar{p}}$ and D_{f_i,g_i} being constants. Note that Eq. (3) is based on perturbative QCD (pQCD) counting rules [17,18] as three hard gluons are needed to produce $\mathbf{B}\bar{\mathbf{B}}'$. Since the $1/t^3$ function peaks at the low mass and flattens out at the large energy region, the threshold effect in the $m_{p\bar{p}}$ spectrum [15] can be easily accommodated. In the approach of pQCD counting rules, we have explained the experimental data observed in baryonic B decays, in particular, the branching fractions of the following decay modes: $B^- \to p\bar{p}K^{(*)-}(\pi^-)$, $\bar{B}^0 \to p\bar{p}K^{(*)0}$, $B^- \to \Lambda\bar{p}\rho^0(\gamma)$, $\bar{B}^0 \to \Lambda\bar{p}\pi^+$, $\bar{B}^0 \to n\bar{p}D^{*+}$, and $\bar{B}^0 \to p\bar{p}D^{(*)0}$. Moreover, the predicted $\mathcal{B}(\bar{B} \to \Lambda\bar{\Lambda}\bar{K}(\pi))$ [19], $\mathcal{B}(\bar{B}^0 \to \Lambda\bar{\Lambda}D^0)$, and $\mathcal{B}(B^- \to \Lambda\bar{p}D^{(*)0})$ [15] are consistent with the latest measurements [20]. These results show that our approach is a reliable one for tackling the three-body baryonic B decays. As we shall see, although the amplitude in Eq. (1) is not suitable to explain the peaking data points in the m_{D^0p} spectrum, it is trustworthy for describing the nonpeaking data points. Inspired by both experimental [10] and theoretical [21] studies of $\Lambda_c(2880)^+$ and $\Lambda_c(2940)^+ \to D^0p$, we shall examine the $\bar{B}^0 \to p\bar{p}D^0$ decay via the resonant $\bar{B}^0 \to B_c(\to D^0p)\bar{p}$ channels to reveal the m_{D^0p} invariant mass spectrum. As in Eq. (1), via the same effective Hamiltoniam for $b \to c\bar{u}d$ at the quark level, the resonant amplitude is derived as

$$\mathcal{A}_R(\bar{B}^0 \to (\mathbf{B_c} \to D^0p)\bar{p}) = \frac{G_F}{\sqrt{2}}V_{cb}V_{ud}^*a_2\langle D^0p|\mathbf{B_c}\rangle\mathcal{R} \times \langle\mathbf{B_c}\bar{p}|(\bar{c}u)_{V-A}(\bar{d}b)_{V-A}|\bar{B}^0\rangle, \tag{4}$$

where the factor $\mathcal{R}$ is the Breit–Wigner factor given by

$$\mathcal{R} = \frac{i}{q^2 - m^2 + im\Gamma}, \tag{5}$$

with $q = p_p + p_{D^0}$ and $m(\Gamma)$ the mass (width) of the resonant $\mathbf{B_c}$ state. We are now ready to evaluate the resonant amplitude $\mathcal{A}_R(\bar{B}^0 \to (\mathbf{B_c} \to D^0p)\bar{p})$. For example, if the quantum number of $\mathbf{B_c}$ is $J^P = 1/2^-$, the combination of the matrix elements for $\bar{B}^0 \to \mathbf{B_c}\bar{p}$ and $\mathbf{B_c} \to D^0p$ in Eq. (4) results in

$$\langle D^0p|\mathbf{B_c}\rangle\langle\mathbf{B_c}\bar{p}|(\bar{c}u)_{V-A}(\bar{d}b)_{V-A}|\bar{B}^0\rangle = g\bar{u}_p\sum u_{\mathbf{B_c}}\bar{u}_{\mathbf{B_c}}(a' + b'\gamma_5)v, \tag{6}$$

with $\Sigma u_{\mathbf{B_c}}\bar{u}_{\mathbf{B_c}} = (\not{q} + m)$ when the $\mathbf{B_c}$ spin is summed over, where g is the coupling constant for the strong decay $\mathbf{B_c} \to D^0p$. Without losing generality, this leads to a reduced form of the resonant amplitude $\mathcal{A}_R(\bar{B}^0 \to (\mathbf{B_c} \to D^0p)\bar{p})$ given by

$$\mathcal{A}_R(\bar{B}^0 \to p\bar{p}D^0) = \frac{G_F}{\sqrt{2}}V_{cb}V_{ud}^*\bar{u}_p\mathcal{R}(a + b\gamma_5)v_{\bar{p}}, \tag{7}$$

where $(a + b\gamma_5) = a_2g(\not{q} + m)(a' + b'\gamma_5)$. The parameters a and b in the above equation can be treated as constants for resonances with a narrow width. Besides, even if $\mathbf{B_c}$ carries a higher spin, such as $\Lambda_c(2880)^+$ with spin $5/2$, the spin structure in the propagator can still be summed over and factored into a and b. Since the coupling of $\mathbf{B_c} \to D^0p$ and the parity of $\mathbf{B_c}$ may not be well determined, there exist some other possibilities for the amplitude in Eq. (4), such as $(a - b\gamma_5)$ and $(b \pm a\gamma_5)$. By taking $|a| = |b|$, we are able to circumvent this complexity. Hence, we can obtain the amplitude squared $|\bar{A}|^2$, with the bar denoting the summation over p and $\bar{p}$ spins. Using the general partial rate formula for the three-body decay given by [6]

$$d\Gamma = \frac{1}{(2\pi)^3}\frac{|\bar{\mathcal{A}}|^2}{32M_{\bar{B}^0}^3}dm_{D^0p}^2dm_{p\bar{p}}^2, \tag{8}$$

we integrate over $m_{p\bar{p}}^2$ to obtain the partial rate as a function of m_{D^0p} in $\bar{B}^0 \to p\bar{p}D^0$.

III. NUMERICAL ANALYSIS

We perform a best χ^2 fit to the data. For the theoretical inputs, we take the Cabibbo-Kobayashi-Maskawa matrix elements $V_{cb} = A\lambda^2$ and $V_{ud} = 1 - \lambda^2/2$ with $A = 0.808$ and $\lambda = 0.2253$ from the Particle Data Group [6] along with the D meson decay constant $f_D = 0.23$ GeV [6,22]. For the parameter a_2 and the constants D_{f_i,g_i} in Eq. (3), we adopt their values with errors as in Ref. [15], which have been fitted to explain the total branching fraction of $\bar{B}^0 \to p\bar{p}D^0$.

We now discuss two different scenarios for the data analysis. In the first scenario, we take the amplitude in Eq. (1) only and fit the data points in Fig. 1. Since the m_{Dp} distribution in Fig. 1 is shown for $m_{p\bar{p}}^2 > 5\ \text{GeV}^2$, we integrate $m_{p\bar{p}}^2$ in Eq. (8) over the range of $5\ \text{GeV}^2 < m_{p\bar{p}}^2 < (m_{p\bar{p}}^2)_{\max}$ with $(m_{p\bar{p}})_{\max}^2$ as a function of m_{D^0p} defined in Ref. [6]. The dashed line in Fig. 1, drawn to link the nonpeaking data points, corresponds to the amplitude of Eq. (1) based on pQCD. The value of χ^2/degrees of freedom is fitted to be $133/27 = 4.9$, where $4.9 = 4.2 + 0.7$, with the last two numbers coming from the peaking data points of 2.82–3.28 GeV and nonpeaking

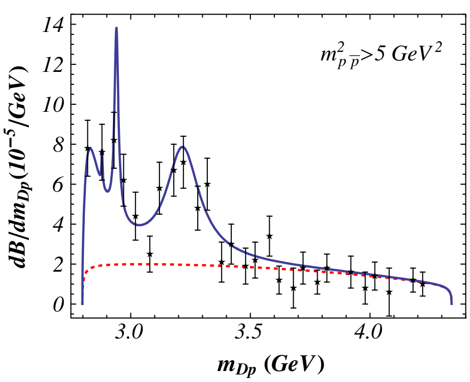

FIG. 1 (color online). Invariant-mass spectrum as a function of the invariant mass m_{D^0p} in $\bar{B}^0 \to p\bar{p}D^0$, where the solid line includes both resonant and nonresonant contributions, while the dashed line corresponds to the nonresonant contribution only.

TABLE I. The resonant charmed baryon ($\mathbf{B_c}$) states in the decay of $\bar{B}^0 \to p\bar{p}D^0$, where the values for $|a|$ and $|b|$ are fitting results. In column 3, input parameters of masses and widths are denoted by the notation*.

$\mathbf{B_c}$ state	$\lvert a\rvert = \lvert b\rvert$	(m, Γ) MeV	p value
$\Sigma_c(2800)^+$	4.61 ± 0.79	*$(2792 \pm 14, 62 \pm 60)$	3.3×10^{-9}
$\Lambda_c(2880)^+$	0.09 ± 0.02	*$(2881.53 \pm 0.35, 5.8 \pm 1.1)$	3.4×10^{-6}
$\Lambda_c(2940)^+$	0.65 ± 0.14	*$(2939.3 \pm 1.5, 17 \pm 8)$	2.1×10^{-6}
$\mathbf{B}_c(3212)^+$	3.37 ± 0.68	$(3212 \pm 20, 167 \pm 34)$	2.9×10^{-7}

data points above 3.28 GeV in the spectrum, respectively. This fitting result strongly suggests the necessity for the presence of resonances.

In the second scenario, we fit the data points in Fig. 1 by combining the nonresonant amplitude in Eq. (1) with the resonant amplitudes for various $\mathbf{B_c}$ states in the form of Eq. (4). It is reasonable to identify the established $\Sigma_c(2800)^+$, $\Lambda_c(2880)^+$, and $\Lambda_c(2940)^+$, with the first three peaks in the Dp spectrum of $\bar{B}^0 \to p\bar{p}D^0$. Particularly, the decays of $\Lambda_c(2880)^+ \to D^0 p$ and $\Lambda_c(2940)^+ \to D^0 p$ have been observed [10]. The theoretical inputs of the masses and the decay widths for the established $\mathbf{B_c}$ states are taken from the Particle Data Group [6], which are also shown in Table I. On the other hand, we need to introduce a new state to accommodate the peak observed around 3.2–3.3 GeV. The relative phases among different resonances are assumed to be positive in the fit. Moreover, the interference between different resonant states is small due to the constrained overlap. The results are presented in Table I and Fig. 1, in which the resonances are accommodated by the solid line. We note that, by adding the resonant states, χ^2/degrees of freedom is reduced to be $17/21 = 0.8$. The new resonant state is named $\mathbf{B}_c(3212)$ with the fitted value of $(m, \Gamma) = (3212 \pm 20, 167 \pm 34)$ MeV. In Table I, as $|a|$ and $|b|$ present the magnitude of a decay, each central value divided by the error indicates the confidence level for a resonant production. As a result, the fitting gives the confidence levels of 5.8σ, 4.5σ, 4.6σ, and 5.0σ for the resonant states of $\Sigma_c(2800)^+$, $\Lambda_c(2880)^+$, $\Lambda_c(2940)^+$, and $\mathbf{B}_c(3212)^+$, with the converted p values shown in Table I, respectively.

With the fit mass in the $m_{D^0 p}$ spectrum, the charmed baryon state $\mathbf{B_c}(3212)^+$ can be identified with one of the states, $\Sigma_c(3262)$ with $J^P = 3/2^+$ and $\Sigma_c(3268)$ with $J^P = 5/2^+$, or the combination of them. The latter state may arise from the orbital (two-dimensional) excitations of the c–(ud) bound state [3]. As for the resonant charmed baryon states found in the $\bar{B}^0 \to p\bar{p}D^0$ decay, one can perform experimental searches. In analog to the first observation of $\Lambda_c(2940)^+$ in the process of $e^+e^- \to \Lambda_c(2940)^+(\to D^0 p)X$ by *BABAR* [10], the $\mathbf{B}_c$ states can be searched for by scanning the $D^0 p$ spectra in $\mathbf{B_c} \to D^0 p$ from the e^+e^-, $p\bar{p}$, and pp colliders at the B factories, Tevatron, and LHC, respectively. We also expect that some of the resonant $\mathbf{B_c}$ states should be observed in the B decays through $\bar{B}^0 \to \mathbf{B_c}^+(\to \Sigma_c(2455)^{0,++}\pi^\pm)\bar{p}$, $\bar{B}^0 \to \mathbf{B_c}^+(\to \Sigma_c(2520)^{0,++}\pi^\pm)\bar{p}$, $B^- \to \mathbf{B_c}^+(\to \Sigma_c(2455)^{0,++}\pi^\pm)\bar{p}\pi^-$, and $B^- \to \mathbf{B_c}^+(\to \Sigma_c(2520)^{0,++}\pi^\pm)\bar{p}\pi^-$. It is important to note that, as the resonant $\bar{B}^0 \to p\bar{p}D^0$ decays imply the existence of the two-body $\bar{B}^0 \to \mathbf{B_c^+}\bar{p}$ decays, e.g., $\bar{B}^0 \to \Lambda_c(2940)^+\bar{p}$, the angular analysis in the $\bar{B}^0 \to \mathbf{B_c}\bar{p}$ decays will be very useful to extract the right quantum numbers for $\mathbf{B}_c$ states. Note that the spin of $\Lambda_c(2940)$ is still unknown. Finally, we remark that the charmed baryons Σ_c and Λ_c in our study would have the counterparts with the well-measured charm-strange baryons Ξ_c [5,23] based on the $SU(3)$ flavor symmetry.

IV. CONCLUSIONS

We have analyzed the peaks in the $m_{D^0 p}$ spectrum observed in the $\bar{B}^0 \to p\bar{p}D^0$ decay with $m^2_{p\bar{p}} > 5\ \text{GeV}^2$ and found that some of the peaking points can be recognized as the charmed baryon resonances in $\bar{B}^0 \to \mathbf{B_c}(\to D^0 p)\bar{p}$ channels. In addition to the established $\Sigma_c(2800)^+$, $\Lambda_c(2880)^+$, and $\Lambda_c(2940)^+$, a new resonance named $\mathbf{B}_c(3212)^+$ has been introduced, which is fitted to have $(m, \Gamma) = (3212 \pm 20, 167 \pm 34)$ MeV for $\bar{B}^0 \to p\bar{p}D^0$. Moreover, since the resonant $\bar{B}^0 \to \mathbf{B_c}(\to D^0 p)\bar{p}$ decays are the consequences of the two-body $\bar{B}^0 \to \mathbf{B_c}\bar{p}$ decays, the studies of the angular distributions in the two-body modes will help determine the spins for the new charmed baryon states, such as, $\Lambda_c(2940)$ and $\mathbf{B_c}(3212)$.

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Note added.—After we posted the paper to the e-print archive (arXiv:1205.0117 [hep-ph]), the *BABAR* collaboration presented the study of $B^- \to \Sigma_c^{++}\bar{p}\pi^-\pi^-$ [24], where they observed a new resonance $\mathbf{B_c}(3245)^0$ with $(m, \Gamma) = (3245 \pm 20, 108 \pm 6)$ MeV in the spectrum of $\Sigma_c^{++}\pi^-\pi^-$. With similar masses and decay widths, $\mathbf{B}_c(3245)^0$ and $\mathbf{B}_c(3212)^+$ can be regarded as the neutral and charged-one components of the excited baryon states of $(\Sigma_c^{++}, \Sigma_c^+, \Sigma_c^0)$ predicted in Ref. [3].

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