Possibility of search for new physics at the CERN LHCb

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It is interesting to search for new physics beyond the standard model at Large Hadron Collider beauty (LHCb). We suggest that weak decays of doubly charmed baryon such as $\Xi_{cc}(3520)^+$, Ξ_{cc}^{++} to charmless final states would be a possible signal for new physics. In this work we consider two models, i.e. the unparticle and Z' as examples to study such possibilities. We also discuss the cases for Ξ_{bb}^0, Ξ_{bb}^- which have not been observed yet, but one can expect to find them when LHCb begins running. Our numerical results show that these two models cannot result in sufficiently large decay widths, therefore if such modes are observed at LHCb, there must be a new physics other than the unparticle or Z' models.

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I. INTRODUCTION

LHC will begin its first run pretty soon, and besides searching for the long-expected Higgs boson, its main goal is to explore new physics beyond the standard model (SM). Many schemes have been proposed to reach the goal. Indeed, the LHCb detector, even though it is not responsible for the Higgs hunting, will provide an ideal place to study heavy flavor physics and search for evidence of new physics. One can make careful measurements on rare decays of B-mesons, b-baryons, B-mixing, and *CP* violation with a huge database available at LHCb; moreover, we are inspired by the possibilities of discovering new physics. It would be beneficial to conjecture more possible processes which would signal the existence of new physics.

In 2002, the first event for s doubly charmed baryon, $\Xi_{cc}^+(3520)$, was observed by the SELEX Collaboration in the channel of $\Xi_{cc}^+ \to \Lambda_c^+ K^- \pi^+$ [1–3]. Ξ_{cc}^+ has the mass $m = 3519 \pm 1$ MeV and width $\Gamma < 5$ MeV. By studying an alternative channel of pD^+K^- conducted later, the mass of the baryon-resonance was confirmed as $m = 3518 \pm 3$ MeV [4], which is consistent with that given in Ref. [1]. In the present theory, there definitely is no reason to exclude the existence of Ξ_{cc}^{++} which contains ccu valence quarks and as well $\Xi_{bb}^0(bbu)$ and $\Xi_{bb}^-(bbd)$, by the flavor-SU(3) symmetry.

In this work, we propose that direct decays of Ξ_{cc} with charmless final states or Ξ_{bb} with bottomless final states would be signals for new physics. By the quark-diagrams, one can easily notice that the main decay modes of Ξ_{cc} would be $D^+\Lambda(\Sigma^0)$, $\Lambda_c K^0$, D^+PK^-, and $\Lambda_c K^-\pi^+$. The latter two modes are just the channels where the SELEX collaboration observed the baryon Ξ_{cc}. Meanwhile the direct decays of Ξ_{cc} (or Ξ_{bb}) into charmless (bottomless) final states are suppressed in the standard model, so that they would be sensitive to new physics beyond the SM.

Since in Ξ_{cc} there are two identical charm quarks, neither can neither annihilate nor exchange W-boson to convert into other quarks. In the SM, direct transition of Ξ_{cc} into charmless final states may be realized via the double-penguin mechanism which is shown in Fig. 1(a), the crossed box-diagram [Fig. 1(b)], and a possible two-step process shown in Fig. 1(c). The mechanism includes two penguin loops or a crossed box-diagram is very sup-

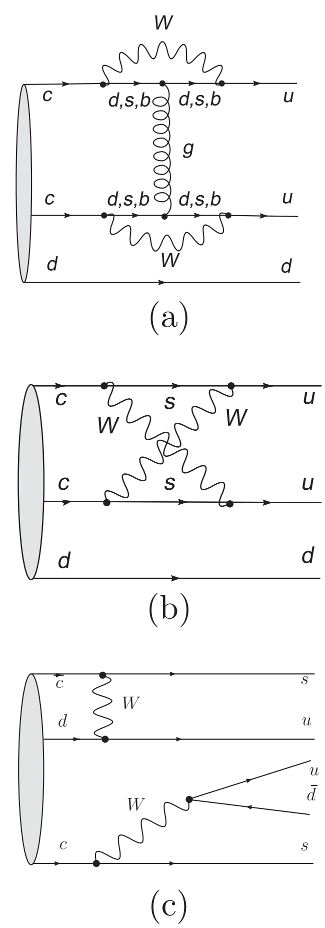


FIG. 1. (a) The double-penguin diagram which can induce the decay of $\Xi_{cc}(\Xi_{bb})$ into noncharm (nonbottom) final states. (b) The crossed-box diagram. (c) An emission where the effective interaction would be nonlocal and for charmless decays of Ξ_{cc}^{++}, it does not exist.

pressed, so that cannot result in any observable effects and we can ignore them completely. If a nonzero rate is observed at LHCb, it should be a signal of new physics. Definitely the diagram of Fig. 1(c) may cause a nonzero contribution and contaminate our situation for exploring new physics. If we consider the charmless decays of Ξ_{cc}^{++} or bottomless decays of Ξ_{bb}^{-}, that diagram [Fig. 1(c)] does not exist at all. Then, the first question is whether we can distinguish such direct decays of Ξ_{cc} into charmless final states (or Ξ_{bb} into bottomless final states) from the secondary decays, which results in charmless (or bottomless) products and are the regular modes in the framework of the SM. The answer is that the direct transitions are favorably two-body decays, namely, in the final states there are only two noncharmed hadrons by whose momenta one can reconstruct the invariant mass spectra of Ξ_{cc} (or Ξ_{bb}), whereas, in the regular modes with sequent decays, there are at least three hadrons in the final states.

The second question is whether there is any mechanism beyond the standard model available which can result in such direct decays. Below, we use two models to demonstrate how such direct decay modes are induced and estimate the widths accordingly. One of them is the unparticle scenario and another one is the $SU(3) \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model where a new gauge boson Z' exists and mediates an interaction to turn the charm quark into a u-quark. Thus by exchange of an unparticle or Z' between the two charm quarks in $\Xi_{cc}^{+(++)}$ (or between the two bottom quarks in $\Xi_{bb}^{-(0)}$), these direct transitions occur.

In this work, for simplicity, we only consider the inclusive decays of $\Xi_{cc}^{++}(\Xi_{bb}^{-})$ into charmless (bottomless) final states. The advantage of only considering the inclusive processes is obvious in that we do not need to worry about the hadronization of quarks into final states because such processes are fully governed by the nonperturbative QCD effects and bring up much uncertainty.

Below, we will investigate the processes caused by exchanging unparticles and Z' separately and then make a brief discussion on the possibility that new physics may result in observable phenomena at LHCb.

II. THE INCLUSIVE DECAY OF DOUBLY CHARMED BARYON

A. The unparticle scenario

Before entering the concrete calculation, we briefly review the concerned knowledge on the unparticle physics [5], which is needed in later derivation. The effective Lagrangian describing the interaction of the unparticle with the SM quarks is

$$\mathcal{L} = \frac{c_S^{qq'}}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \bar{q}\gamma_\mu(1-\gamma_5)q'\partial^\mu O_{\mathcal{U}} + \frac{c_V^{qq'}}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \bar{q}\gamma_\mu(1-\gamma_5)q' O_{\mathcal{U}}^\mu + \text{H.c.}, \tag{1}$$

where $O_{\mathcal{U}}$ and $O_{\mathcal{U}}^\mu$ are the scalar and vector unparticle fields, respectively. q and q' denote the SM quark fields. Generally, the dimensionless coefficients $c_{S,V}^{qq'}$ is related to the flavor of the quark field. This interaction induces a flavor-changing neutral current (FCNC) and contributes to the processes of concern.

For a scalar unparticle field, the propagator with momentum p and scale dimension $d_{\mathcal{U}}$ is [6]

$$\int d^4xd^{ip\cdot x}\langle 0|TO_{\mathcal{U}}(x)O_{\mathcal{U}}(0)|0\rangle = i\frac{A_{d_{\mathcal{U}}}}{2\sin(d_{\mathcal{U}}\pi)}\frac{1}{(-p^2-i\epsilon)^{2-d_{\mathcal{U}}}} \tag{2}$$

with

$$A_{d_{\mathcal{U}}} = \frac{16\pi^{5/2}}{(2\pi)^{2d_{\mathcal{U}}}}\frac{\Gamma(d_{\mathcal{U}}+1/2)}{\Gamma(d_{\mathcal{U}}-1)\Gamma(2d_{\mathcal{U}})}, \tag{3}$$

with $d_{\mathcal{U}}$ the scale dimension.

For the vector unparticle, the propagator reads

$$\int d^4xe^{ip\cdot x}\langle 0|TO_{\mathcal{U}}^\mu O_{\mathcal{U}}^\nu(0)|0\rangle = i\frac{A_{d_{\mathcal{U}}}}{2\sin(d_{\mathcal{U}}\pi)} \times \frac{-g^{\mu\nu}+p^\mu p^\nu/p^2}{(-p^2-i\epsilon)^{2-d_{\mathcal{U}}}}, \tag{4}$$

where the transverse condition $\partial_\mu O_{\mathcal{U}}^\mu = 0$ is required.

In the unparticle physics, the inclusive decay of doubly charmed baryons into light quarks $ccq \to uuq$ occurs at the tree level, and the transition is depicted in Fig. 2. Here the exchanged agent between the two charm quarks can be either scalar or vector unparticle.

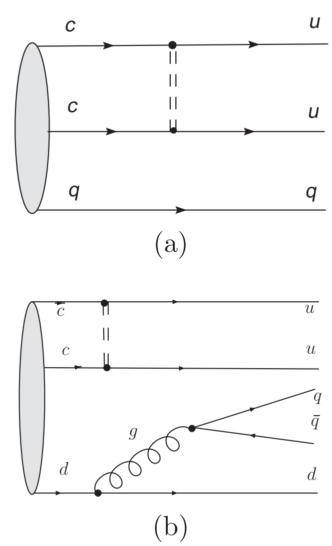


FIG. 2. The inclusive transition of doubly charmed baryon in unparticle physics, where the double-dashed line denotes the scalar or vector unparticle in the unparticle model or Z' in the left-right model.

Even though we only consider the inclusive processes where the quarks in the final states are treated as on-shell free particles and the wave functions of the light hadrons in the final states are not needed, the binding effect of the initial baryons (Ξ_{cc}^{++} or Ξ_{bb}^{-}) which are composed of three valence quarks must be taken into account. Namely, when we calculate the hadronic matrix elements, we need to invoke concrete phenomenological models to carry out the computations where the wave function of the initial baryon is needed. In this work, we adopt a simple nonrelativistic model, i.e. the harmonic oscillator model [7]. This model has been widely and successfully employed in similar research [8–14]. Thus one can trust that for heavy hadrons, such a simple nonrelativistic model can work well and the results are relatively reliable, even though certain errors are not avoidable. Thus in this work the matrix elements of the effective operators evaluated in terms of the harmonic oscillator wave function are believed to be a good approximation. According to the references listed above, the errors in the estimate, especially as we only need the wave function of the initial hadron, are expected to be less than 10%. By changing the input parameters and the model parameters, which are obtained by fitting other experiments, we scan the region of changes of the numerical results and find that the error range is indeed consistent with our expectation.

In the harmonic oscillator model, the wave function of the doubly charmed baryon $\Xi_{cc}(3520)^+$ is expressed as

$$|\Psi_{\Xi_{cc}^+}(P,s)\rangle = \mathcal{N}\sum_{\text{color,spin}}\chi_{\mathcal{SF}}\varphi_{\mathcal{C}}\int d^3p_\rho d^3p_\lambda \times \psi_{\Xi_{cc}^+}(\mathbf{p}_\rho,\mathbf{p}_\lambda)b_c^\dagger(p_1',s_1')b_c^\dagger(p_2',s_2') \times b_u^\dagger(p_3',s_3')|0\rangle,$$

which satisfies the normalization condition

$$\langle\Psi_{\Xi_{cc}^+}(\mathbf{P},s)|\Psi_{\Xi_{cc}^+}(\mathbf{P}',s')\rangle = (2\pi)^3\frac{M_{\Xi_{cc}^+}}{E_P}\delta^3(\mathbf{P}-\mathbf{P}')\delta(s-s'),$$

where $\mathcal{N}$ is the normalization constant. $\chi_{\mathcal{SF}}$ and $\varphi_{\mathcal{C}}$ denote the spin-flavor and color parts of the wave function of doubly charmed baryon $\Xi_{cc}(3520)^+$ respectively whose explicit expressions are

$$\mathcal{N} = \sqrt{\frac{E}{M}\frac{m_1'm_2'm_3'}{E_1'E_2'E_3'}}, \qquad \varphi_{\mathcal{C}} = \frac{1}{\sqrt{6}}\epsilon_{ijk},$$

$$\chi_{\mathcal{SF}} = \frac{1}{\sqrt{6}}[2|c\uparrow c\uparrow d\downarrow\rangle - |c\uparrow c\downarrow d\uparrow\rangle - |c\downarrow c\uparrow d\uparrow\rangle].$$

In the harmonic oscillator model, the spatial wave function $\psi_{\Xi_{cc}^+}(\mathbf{p}_\rho,\mathbf{p}_\lambda)$ reads as

$$\psi_{\Xi_{cc}^+}(\mathbf{p}_\rho,\mathbf{p}_\lambda) = 3^{3/4}\left(\frac{1}{\pi\alpha_\rho^2}\right)^{3/4}\left(\frac{1}{\pi\alpha_\lambda^2}\right)^{3/4} \times \exp\left[-\frac{\mathbf{p}_\rho^2}{2\alpha_\rho}-\frac{\mathbf{p}_\lambda^2}{2\alpha_\lambda}\right]$$

with the definitions

$$\mathbf{p}_\rho = \frac{\mathbf{p}_1'-\mathbf{p}_2'}{\sqrt{2}}, \qquad \mathbf{p}_\lambda = \frac{\mathbf{p}_1'+\mathbf{p}_2'-\frac{2m_c}{m_d}\mathbf{p}_3'}{\sqrt{2\frac{2m_c+m_d}{m_d}}},$$

$$\mathbf{P} = \mathbf{p}_1'+\mathbf{p}_2'+\mathbf{p}_3',$$

and the parameters α_ρ and α_λ reflect the nonperturbative effects and will be given in a later subsection.

In the center of mass frame of $\Xi_{cc}(3520)^+$, the hadronic matrix elements S_{fi} is written as

$$S_{fi} = (2\pi)^4\delta^4(p_1+p_2+p_3-M)T$$

with $T = (T_S + T_V)$.

For exchanging scalar unparticle, T_S matrix element is written as

$$\begin{aligned}T_S = \sum_{\text{spin}}\int d^3p_\rho d^3p_\lambda(2\pi)^3\frac{E_{p_3}}{m_{p_3}}[\bar{u}_u(p_1,s_1)\gamma_\mu(1-\gamma_5)\\ \times u_c(p_1',s_1')\bar{u}_u(p_2,s_2)\gamma_\nu(1-\gamma_5)u_c(p_2',s_2')]\\ \times\left(\frac{c_S^{cu}}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}}\right)^2\frac{A_{d_{\mathcal{U}}}}{2\sin(d_{\mathcal{U}}\pi)}\frac{iq^\mu q^\nu}{(-p^2-i\epsilon)^{2-d_{\mathcal{U}}}}\\ \times\mathcal{N}\psi_{\Xi_{cc}^+}(\mathbf{p}_\rho,\mathbf{p}_\lambda).\end{aligned} \tag{5}$$

For the vector unparticle exchange, T_V is

$$\begin{aligned}T_V = \sum_{\text{spin}}\int d^3p_\rho d^3p_\lambda(2\pi)^3\frac{E_{p_3}}{m_{p_3}}[\bar{u}_u(p_1,s_1)\gamma_\mu(1-\gamma_5)\\ \times u_c(p_1',s_1')\bar{u}_u(p_2,s_2)\gamma_\nu(1-\gamma_5)u_c(p_2',s_2')]\\ \times\left(\frac{c_V^{cu}}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}}\right)^2\frac{A_{d_{\mathcal{U}}}}{2\sin(d_{\mathcal{U}}\pi)}\frac{i(-g^{\mu\nu}+q^\mu q^\nu/q^2)}{(-p^2-i\epsilon)^{2-d_{\mathcal{U}}}}\\ \times\mathcal{N}\psi_{\Xi_{cc}^+}(\mathbf{p}_\rho,\mathbf{p}_\lambda).\end{aligned} \tag{6}$$

Here u_q and $\bar{u}_q$ ($q = c, u$) denote the Dirac spinors

$$u_q = \sqrt{\frac{E_q+m_q}{2m_q}}\begin{pmatrix}1\\ \frac{\sigma\cdot p}{E_q+m_q}\end{pmatrix}\chi, \tag{7}$$

$$\bar{u}_q = \sqrt{\frac{E_q+m_q}{2m_q}}\chi^\dagger\left(1, -\frac{\sigma\cdot p}{E_q+m_q}\right), \tag{8}$$

and we can use the expression [15]

$$\frac{|c_s^{cu}|^2}{\Lambda_{\mathcal{U}}^{2d_{\mathcal{U}}}} = \frac{6m\Delta m|\sin d_{\mathcal{U}}\pi|}{5f^2\hat{B}A_{d_{\mathcal{U}}}m^{2d_{\mathcal{U}}}}, \tag{9}$$

$$\frac{|c_s^{cu}|^2}{\Lambda_{\mathcal{U}}^{2d_{\mathcal{U}}-1}} = \frac{2m\Delta m|\sin d_{\mathcal{U}}\pi|}{f^2\hat{B}A_{d_{\mathcal{U}}}m^{2d_{\mathcal{U}}-2}} \tag{10}$$

to simplify T_V and T_S. One needs to sum over all possible spin assignments for the Dirac spinors.

B. The Z' scenario

The left-right models [16] are also a natural extension of the electroweak model. It has been widely applied to the analysis on high energy processes. For example, recently He and Valencia [17] employed this model with certain modifications to explain the anomaly in A_{FB}^b observed at LEP [18]. Barger *et al.* studied Z' mediated flavor-changing neutral currents in B-meson decays [19], $B_s-\bar{B}_s$ mixing [20] and $B\to K\pi$ puzzle [21].

The gauge group of the model [17] is $SU(3)\times SU(2)_L\times SU(2)_R\times U(1)_{B-L}$ where the four gauge couplings g_3, g_L, g_R, and g correspond to the four subgroups, respectively. The vacuum expectation values of the three Higgs bosons break the symmetry. The symmetry breaking patterns are depicted in the literature. The introduction of a scalar field ϕ causes Z_0 in the standard model to mix with a new gauge boson Z_R, then Z, Z' are the mass eigenstates.

For the neutral sector the Lagrangian is

$$\begin{aligned}\mathcal{L}=&-\frac{g_L}{2\cos\theta_W}\bar{q}\gamma^\mu(g_V-g_A\gamma_5)q(\cos\xi_Z Z_\mu-\sin\xi_Z Z'_\mu)\\&+\frac{g_Y}{2}\tan\theta_R\Big(\frac{1}{3}\bar{q}_L\gamma^\mu q_L+\frac{4}{3}\bar{u}_{Ri}\gamma^\mu u_{Ri}-\frac{2}{3}\bar{d}_{Ri}\gamma^\mu d_{Ri}\Big)\\&\times(\sin\xi_Z Z_\mu-\cos\xi_Z Z'_\mu)+\frac{g_Y}{2}(\tan\theta_R+\cot\theta_R)\\&\times(sin\xi_Z Z_\mu-cos\xi_Z Z'_\mu)(V^{d*}_{Rbi}V^d_{Rbj}\bar{d}_{Ri}\gamma^\mu d_{Rj}\\&-V^{u*}_{Rti}V^u_{Rtj}\bar{u}_{Ri}\gamma^\mu u_{Rj}).\end{aligned}$$

Here θ_W is the electroweak mixing angle ($\tan\theta_W=\frac{g_Y}{g_L}$), θ_R parametrizes the relative strength of the right-handed interaction ($\tan\theta_R=\frac{g}{g_R}$), ξ_Z is the $Z-Z'$ mixing angle and $V^{u,d}_{Rij}$ are two unitary matrices that rotate the right-handed up-(down)-type quarks from the weak eigenstates to the mass eigenstates. Note that we use current notation for the Pati-Salam model, and only the third family couples to $SU(2)_R$ in this model.

In the Z' model, inclusive decay of doubly charmed baryons into light quarks $ccq\to uuq$ occurs at tree level. The Feynman diagram (Fig. 2) is the same as that for the unparticle scenario, but only the exchanged agent is replaced by Z'.

In the center of mass frame of Ξ^+_{cc}, we can obtain

$$\begin{aligned}T=&\sum_{\rm spin}\int d^3p_\rho d^3p_\lambda(2\pi)^3\frac{E_{p_3}}{m_{p_3}}[\bar{u}_u(p_1,s_1)\gamma_\mu(1-\gamma_5)\\&\times u_c(p'_1,s'_1)\bar{u}_u(p_2,s_2)\gamma_\nu(1-\gamma_5)u_c(p'_2,s'_2)]\frac{-ig^{\mu\nu}}{(p^2+M'^2_Z)}\\&\times\Big[\frac{g_L\tan\theta_W(\tan\theta_R+\cot\theta_R)\cos\xi_Z}{2}V^{u*}_{Rti}V^u_{Rtj}\Big]^2\\&\times\mathcal{N}\psi_{\Xi^+_{cc}}(\mathbf{p}_\rho,\mathbf{p}_\lambda).\end{aligned}\tag{12}$$

The authors of Refs. [17,22,23] suggested that $\cot\theta_R$ is large, so that $\tan\theta_R$ can be ignored. They took approximations $\tan\theta_W\cot\theta_R\frac{M_W}{M'_Z}\sim1$ and $\cos\xi_Z\sim1$. Because $M_{Z'}$ is larger than 500 GeV [17], one has $\frac{1}{p^2+M'^2_Z}\sim\frac{1}{M'^2_Z}$.

Then we have the final expression as

$$\begin{aligned}T=&\sum_{\rm spin}\int d^3p_\rho d^3p_\lambda(2\pi)^3\frac{E_{p_3}}{m_{p_3}}[\bar{u}_u(p_1,s_1)\gamma_\mu(1-\gamma_5)\\&\times u_c(p'_1,s'_1)\bar{u}_u(p_2,s_2)\gamma_\nu(1-\gamma_5)u_c(p'_2,s'_2)]\\&\times(-ig^{\mu\nu})G_F\sqrt{2}(V^{u*}_{Rti}V^u_{Rtj})^2\mathcal{N}\psi_{\Xi^+_{cc}}(\mathbf{p}_\rho,\mathbf{p}_\lambda).\end{aligned}\tag{13}$$

C. The expression of the decay width

The inclusive decay rate would be obtained by integrating over the phase space which involves three free quarks and the procedure is standard [24],

$$\begin{aligned}\Gamma(\Xi_{cc}(3520)^+\to uud)=&\int_{a_1}^{a_2}dp_1^0\int_{b_1}^{b_2}dp_2^0\int_0^{2\pi}d\eta\\&\times\int_{-1}^1 d(\cos\theta)\frac{|T|^2}{16M_{\Xi^+_{cc}}(2\pi)^4},\end{aligned}\tag{14}$$

where a_1, a_2, b_1, and b_2 are defined as, respectively,

$$\begin{aligned}&a_1=0,\qquad a_2=\frac{M_{\Xi^+_{cc}}}{2}-\frac{(m_2+m_3)^2-m_1^2}{2M_{\Xi^+_{cc}}},\\&b_1=\frac{1}{2\tau}[\sigma(\tau+m_+m_-)-\sqrt{p_1^{02}(\tau-m_+^2)(\tau-m_-^2)}],\\&b_2=\frac{1}{2\tau}[\sigma(\tau+m_+m_-)+\sqrt{p_1^{02}(\tau-m_+^2)(\tau-m_-^2)}],\\&\sigma=M_{\Xi^+_{cc}}-p_1^0,\qquad\tau=\sigma^2-\sqrt{(p_1^0)^2},\\&m_\pm=m_2\pm m_3.\end{aligned}$$

Here $M_{\Xi^{++}_{cc}}$, m_1, m_2, and m_3 denote the masses of the doubly charmed baryon, up, and down quarks, respectively. In the following, for obtaining numerical results we use the Monte Carlo method to carry out this integral. For baryon Ξ^-_{bb}, the expression is the same but only the mass of the charm quark is replaced by that of the bottom quark.

III. NUMERICAL RESULTS

Now we present our numerical results.

Since only Ξ^{++}_{cc} has been measured, in the later calculation we use its measured mass as input, and for Ξ^-_{bb} we will only illustrate the dependence of its decay rate on the parameters. The input parameters include: $G_F=1.166\times10^{-5}$ GeV^{-2}, $m_c=1.60$ GeV, $m_u=m_d=0.3$ GeV, $m_s=0.45$ GeV, $m_b=4.87$ GeV. $M_{\Xi^+_{cc}}=3.519$ GeV. $\alpha_\rho=0.33$ GeV^2, $\alpha_\lambda=0.25$ GeV^2 [1,25–27]. Here, the light quark mass refers to the constituent mass.

A. The results in the unparticle scenario

For the unknown parameters $\Lambda_{\mathcal{U}}$ in the unparticle scenario, according to the general discussion, the energy scale may be at order of TeV, thus one can fix $\Lambda_{\mathcal{U}} = 1$ TeV. We choose $d_{\mathcal{U}} = 3/2$ in our calculation.

In this work, we also calculate the inclusive decay width of doubly bottomed baryon Ξ_{bb}^-. The mass of Ξ_{bb}^- is set as 10.09 GeV according to the estimate of Ref. [27], although there are no data available yet.

The numerical results are provided in Table I. Figure 3 illustrates the dependence of the decay widths of $\Xi_{cc}(3520)^+ \to uud$ and $\Xi_{bb}^- \to ddd$ on $d_{\mathcal{U}}$, the three lines (solid, dashed, and dotted) correspond to the contributions of scalar unparticle, vector unparticle, and both on $d_{\mathcal{U}}$. It is noted here that "both" means that at present we cannot determine whether the unparticle is a scalar or a vector and it is also possible that both scalar and vector exist simultaneously. Thus we assume both scalar and vector contribute and they interfere constructively. Definitely, it is worthy of further investigation.

B. The results for Z' exchange

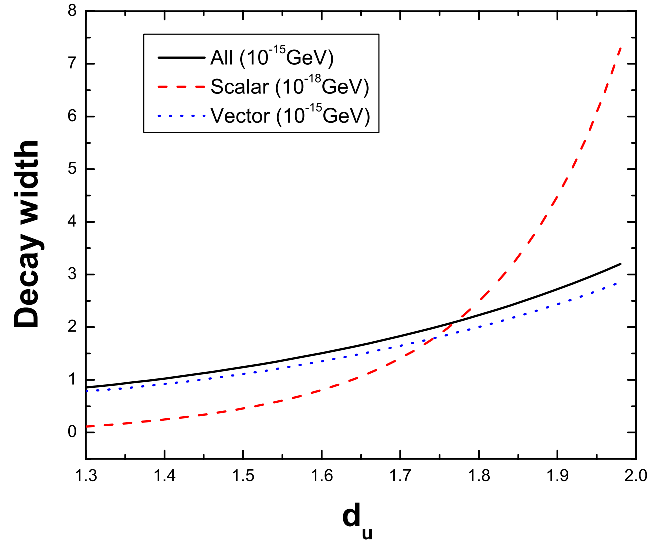
The earlier studies indicate that the mass of $M_{Z'}$ should be larger than 500 GeV [17] and $V^{u*}_{Rtc}V^{u*}_{Rtu}$ is bound no more than 2.0×10^{-4} [22]. In our calculation we take their extreme values as $M_{Z'} = 500$ GeV and $V^{u*}_{Rtc}V^{u*}_{Rtu} = 2.0 \times 10^{-4}$, thus we would obtain the upper limit of the decay width. It is estimated with all the input parameters as

$$\Gamma[\Xi_{cc}(3520)^+ \to uud] = 7.66 \times 10^{-21}\ \text{GeV}. \quad (15)$$

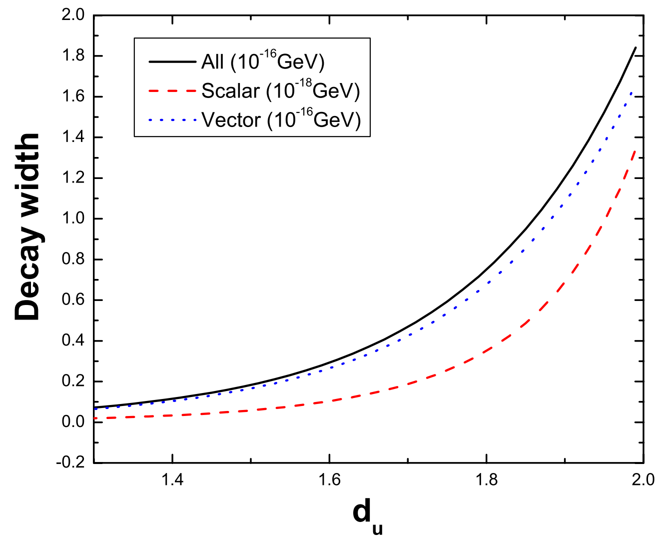
This is a too small numerical value compared with the width of Ξ_{cc}^+, therefore, it is hopeless to observe a nonzero branching ratio of Ξ_{cc}^+ into charmless final states if only Z' is applied.

C. Estimate the contribution from standard model

As indicated above, in the framework of the SM, Ξ_{cc}^+ can decay into two-body final states via Fig. 1(c). It would be interesting to compare the SM contribution with that from the two models. Thus, we would roughly estimate the ratio of the contribution of Fig. 1(c) to that of Fig. 2(b) for the Z' model. It is easier to compare them because the structures of two diagrams and the relevant effective vertices are similar.



(a)



(b)

FIG. 3 (color online). (a) and (b), respectively, show the dependences of the decay widths of $\Xi_{cc}(3520)^+ \to uud$ and $\Xi_{bb}^- \to ddd$ respectively coming from the contributions of scalar unparticle, vector unparticle, and both on $d_{\mathcal{U}}$.

By the order of magnitude estimation and with the $SU(3)$ symmetry, the ratio of the amplitude of Fig. 1(c) ($T_{\rm SM}$) to Fig. 2(b) (T_{un} or $T_{Z'}$) is

TABLE I. The decay widths of $\Xi_{cc}(3520)^+ \to uud$ or $\Xi_{cc}^{++} \to uuu$ and $\Xi_{bb}^- \to ddd(ssd)$ corresponding to $d_{\mathcal{U}} = 3/2$ (in units of GeV). In the table, the second, third, and fourth columns, respectively, correspond to the contributions from exchanging scalar unparticle, vector unparticle, and both.

	scalar	vector	scalar + vector
$\Gamma[\Xi_{cc}(3520)^+ \to uud]$	4.57×10^{-18}	1.11×10^{-15}	1.24×10^{-15}
$\Gamma[\Xi_{bb}^- \to ddd]$	5.85×10^{-20}	1.65×10^{-17}	1.83×10^{-17}
$\Gamma[\Xi_{bb}^- \to ssd]$	5.21×10^{-20}	1.23×10^{-17}	1.44×10^{-17}

$$\frac{T_{\rm SM}}{T_{Z'}} \approx \frac{8G_F}{\frac{4\pi\alpha_s}{q^2}\sqrt{2}(V^{u*}_{Rti}V^u_{Rtj})^2}, \tag{16}$$

where q is the momentum of the unparticle or Z' (in the case of the SM, $q^2 \ll M_W^2$) and can be neglected in the propagator. Because the contribution from smaller q^2 is dominant, in the estimation we set $q^2 = 0.5\ \mathrm{GeV}^2$. Since the whole case under consideration may fall in the nonperturbative QCD region, as a rough estimate we take $\alpha_s = 1$, $V^{u*}_{Rti}V^u_{Rtj} = 2 \times 10^{-4}$, and $G_F = 1.166 \times 10^{-5}\ \mathrm{GeV}^{-2}$, we can get $\frac{T_{\rm SM}}{T_{Z'}} \approx 66$ (i.e. $\frac{\Gamma_{\rm SM}}{\Gamma_{Z'}} \approx 4300$). Then we can obtain the ratio of decay widths $\frac{\Gamma_{\rm Unparticle}}{\Gamma_{\rm SM}} \approx 40$. This ratio indicates that for Ξ^+_{cc} the contribution of the SM is smaller than that of the unparticle scenario, but larger than that from the Z' model.

However, for Ξ^{++}_{cc}, Fig. 1(c) does not contribute at all, so that the decay of Ξ^{++}_{cc} into charmless final states (or Ξ^-_{bb} into bottomless final states) is more appropriate for exploring new physics than Ξ^+_{cc}.

It is worth noticing that the estimate of the contribution of the SM to the decay rate is very rough, thus what we can assure to ourselves is its order of magnitude. Indeed the magnitude contributed by the SM is very small and cannot produce sizable observational effects at all, even though it has a comparable order with that from the two sample models, the unparticle and Z'. In the future, if such mode were observed at LHCb we can definitely conclude that it is not caused by the SM but new physics.

IV. DISCUSSION AND CONCLUSION

In this work we propose to explore for new physics beyond the SM at LHCb by measuring direct decays of $\Xi^+_{cc}\Xi^{++}_{cc}$, (Ξ^-_{bb}, Ξ^0_{bb}) into charmless (bottomless) final states. Such decays can occur via the diagrams shown in Fig. 1 in the framework of the SM, but is much suppressed to be experimentally observed; therefore, if a sizable rate is measured it would be a clear signal for new physics beyond the SM. We use two models as examples, namely, the unparticle and Z' models to calculate the decay rates, because both of them allow a transition of $cc(bb) \to qq$ where q may be light quarks to occur at tree level. Thus one expects that these new models might result in nonzero observation.

Indeed, our work is motivated by three factors: first the great machine LHC will run next year and a remarkable amount of data will be available; second, the double-charmed baryon Ξ_{cc} which was observed by the SELEX collaboration provides us a possibility to probe new physics; and the last reason is that some models have been proposed and they may induce a flavor-changing neutral current—concretely the unparticle and Z' models are employed in this work. Definitely none of the two models are confirmed by either theory or experiment yet, and they still need further theoretical investigations; however, their framework is clear, so we may use them as examples to demonstrate how new physics may cause such decay modes and indicate that a sizable observational rate is a clear signature for new physics beyond the SM. Moreover, the double-charmed baryon has only been observed by the SELEX collaboration, but not at B-factories. It seems peculiar at first glimpse, but careful studies indicate that it is quite reasonable due to the fragmentation process of heavy quarks. The authors of Ref. [28] indicate that the meson B_c cannot be seen at any e^+e^- colliders because its production rate at such machines is too small, but contrarily, its production rate is greatly enhanced at hadron colliders. It was first observed at TEVATRON and its production rate at LHC would be much larger by several orders [28]. In analog, one can expect that such double-charmed baryons Ξ_{cc} or double-bottomed Ξ_{bb} can only be produced at LHC, but not at B-factories.

The inclusive decays of doubly charmed baryons $\Xi_{cc}(3520)^+$, Ξ^{++}_{cc}, and Ξ^0_{bb}, Ξ^-_{bb} are explored in unparticle and Z' scenarios. Our result indicates that the upper limit of the inclusive decay width of $\Xi^{++}_{cc} \to uuu$ is about 10^{-15} GeV with $d_{\mathcal{U}} = 3/2$. For inclusive decay $\Xi^-_{bb} \to ddd(ssd)$, the upper limit is at order of 10^{-17} GeV. It is learned that in the unparticle scenario, the contribution from exchanging a vector unparticle is much larger than that from exchanging a scalar unparticle, as shown in Table I.

The parameters which we employ in the numerical computations are obtained by fitting other experimental measurements, for example, if the recently observed $D^0 - \bar{D}^0$ can be interpreted by the unparticle model, an upper bound on the parameters in the model would be constrained. Indeed, all the present experimental data can only provide upper bounds on the model parameters no matter what the new physics model under consideration is.

So far it is hard to make an accurate estimate on the production rates of the heavy baryons which contain two heavy quarks at LHC yet, but one has reason to believe that the production rate would be roughly of the same order of the production rate of B_c which was evaluated by some authors [28], or even smaller by a factor of less than 10. The production rates indeed will be theoretically evaluated before or even after LHC begins running.

In Ref. [29], the authors estimate the number of Ξ_{cc} produced at LHCb as about 10^9. Since the available energy is much higher than the masses of Ξ_{cc} and Ξ_{bb}, one has strong reason to believe that their production rates are comparable. Unfortunately our numerical results indicate that the unparticle and Z' scenarios cannot result in sizable rates for $\Xi^{++}_{cc} \to uuu \to$ two hadrons and $\Xi^-_{bb} \to ddd(ssd) \to$ two hadrons which can be measured at LHCb and neither the SM. Even though the two sample models and SM cannot cause sufficiently large rates, the channels still may stand for a possible place to search for new physics. If a sizable rate is observed at LHCb, it would be a signal of new physics and indicate that the new physics is also not the unparticle and/or Z', but something else.

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