A Possibility of Search for New Physics at LHCb

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It is interesting to search for new physics beyond the standard model at LHCb. We suggest that weak decays of doubly charmed baryon such as \( \Xi_{cc}^{++}(3520) \), \( \Xi_{bb}^{++} \) 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 \( \Xi_{bb}^{++} \), \( \Xi_{bb}^{−} \) 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 SM. Many schemes have been proposed to reach the goal. Indeed, the LHCb detector, even though 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 existence of new physics.

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

In this work, we propose that direct decays of \( \Xi_{cc}^{++} \) with charmless final states or \( \Xi_{bb}^{++} \) with bottomless final states would be signals for new physics. By the quark-diagrams, one can easily notice that the main decay modes of \( \Xi_{cc}^{+} \) would be \( D^{+}\Lambda_0(\not\Sigma^0) \), \( \Lambda_c K^0 \), \( D^+PK^− \) and \( \Lambda_c K^−\pi^+ \). The later two modes are just the channels where the SELEX collaboration observed the baryon \( \Xi_{cc} \). While the direct decays of \( \Xi_{cc}^{+} \) into charmless (bottomless) final states are suppressed in the standard model, so that would be sensitive to new physics beyond the SM.

Since in \( \Xi_{cc} \) there are two identical charm quarks which can neither annihilate, nor exchange W-boson to convert into other quarks. In the SM, direct transition of \( \Xi_{cc} \) into charmless final states may realize 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 suppressed, so that cannot result in any observable effects and we can ignore them completely. If a non-zero rate is observed at LHCb, it should be a signal of new physics. Definitely the diagram of Fig. 1 (c) may cause a non-zero contribution and contaminate our situation for exploring new physics. If we consider the charmless decays of \( \Xi_{cc}^{++} \) or bottomless decays of \( \Xi_{bb}^{++} \) that diagram (Fig. 1 (c)) does not exist at all. Then, the first question is that can we distinguish such direct decays of \( \Xi_{cc}^{++} \) into charmless final states (or \( \Xi_{bb}^{++} \) into bottomless final states) from the secondary decays which result 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 non-charmed hadrons by whose momenta one can re-construct the invariant mass spectra of \( \Xi_{cc}^{++} \) (or \( \Xi_{bb}^{++} \)), whereas, in the regular modes with sequential decays, there are at least three hadrons in the final states.

The second question is that is there 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) \( SU(2)_L \times SU(2)_R \times U(1)_{B-L} \) 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 \( \Xi_{bb}^{++} \)), these direct transitions occur.

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

Below, we will investigate the processes caused by exchanging unparticles and \( Z' \) separately and then make a
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 \( R \), which is needed in later derivation. The effective Lagrangian describing the interaction of the unparticle with the SM quarks is

\[
\mathcal{L} = \frac{c_{qq'}}{\Lambda_{\text{ut}}^d} \bar{q} \gamma_{\mu}(1 - \gamma_5)q' \partial^\mu O_{\text{ut}} + \frac{c_{qq'}}{\Lambda_{\text{ut}}^d - 1} \bar{q} \gamma_{\mu}(1 - \gamma_5)q' O_{\text{ut}}^\nu + \text{h.c.,}
\]

where \( O_{\text{ut}} \) and \( O_{\text{ut}}^\nu \) are the scalar and vector unparticle fields respectively. \( q \) and \( q' \) denote the SM quark fields. Generally, the dimensionless coefficients \( c_{qq'} \) is related to the flavor of the quark field. This interaction induces a FCNC and contributes to the processes of concern.

For a scalar unparticle field, the propagator with momentum \( p \) and scale dimension \( d_{\text{ut}} \) is

\[
\int d^4x d^{p-x} \langle 0 | T O_{\text{ut}}(x) O_{\text{ut}}(0) | 0 \rangle = i \frac{A_{\text{ut}}}{2 \sin(d_{\text{ut}} \pi)} \frac{1}{(-p^2 - i\epsilon)^{2 - d_{\text{ut}}}}
\]

with

\[
A_{\text{ut}} = \frac{16\pi^{5/2}}{(2\pi)^{d_{\text{ut}}}} \frac{\Gamma(d_{\text{ut}} + 1/2)}{\Gamma(d_{\text{ut}} - 1)\Gamma(2d_{\text{ut}})}
\]

For the vector unparticle, the propagator reads

\[
\int d^4x e^{ip-x} \langle 0 | T O_{\text{ut}}^\mu O_{\text{ut}}^\nu(0) | 0 \rangle = i \frac{A_{\text{ut}}}{2 \sin(d_{\text{ut}} \pi)} \frac{-g_{\mu\nu} + p_{\mu} p_{\nu}/p^2}{(-p^2 - i\epsilon)^{2 - d_{\text{ut}}}}
\]

where the transverse condition \( \partial_{\mu} O_{\text{ut}}^\mu = 0 \) is required.

In the unparticle physics, the inclusive decay of doubly charmed baryons into light quarks \( ccq \rightarrow uuq \) 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.
and the parameters $\alpha_p$ and $\alpha_\lambda$ reflect the non-perturbative effects and will be given in later subsection.

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

$$S_{fi} = (2\pi)^3 \delta(p_1 + p_2 + p_3 - M)T$$

with $T = (T_S + T_V)$.

For exchanging scalar unparticle, $T_S$ matrix element is written as

$$T_S = \sum_{spin} \int d^3p_\rho d^3p_\lambda (2\pi)^3 \frac{E_{p_\rho}}{m_{p_\rho}} \times \bar{u}_s(p_1, s_1) \gamma_\mu (1 - \gamma_5) u_c(p_2', s_2') \times u_s(p_2, s_2) \gamma_\nu (1 - \gamma_5) u_c(p_2', s_2') \times \left( \frac{c^{u\nu}_{\lambda\mu}}{A_{d\lambda}} \right)^2 \frac{i(q^\mu q'^\nu - q^\nu q'^\mu)}{2 \sin(d\Delta_\pi)} \times (p_1 - p_2 - (p_1 - p_2 - p_3))^2 - d_{uv} \times N \psi_{\Xi_{cc}^+}^{p_\rho s_\rho}(p_\rho, p_\lambda).$$

For the vector unparticle exchange, $T_V$ is

$$T_V = \sum_{spin} \int d^3p_\rho d^3p_\lambda (2\pi)^3 \frac{E_{p_\rho}}{m_{p_\rho}} \times \bar{u}_s(p_1, s_1) \gamma_\mu (1 - \gamma_5) u_c(p_2', s_2') \times u_s(p_2, s_2) \gamma_\nu (1 - \gamma_5) u_c(p_2', s_2') \times \left( \frac{c^{u\nu}_{\lambda\mu}}{A_{d\lambda}} \right)^2 \frac{i(q^\mu q'^\nu - q^\nu q'^\mu)}{2 \sin(d\Delta_\pi)} \times (p_1 - p_2 - (p_1 - p_2 - p_3))^2 - d_{uv} \times N \psi_{\Xi_{cc}^+}^{p_\rho s_\rho}(p_\rho, p_\lambda).$$

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

$$u_q = \frac{\sqrt{E_q + m_q}}{2m_q} \left( \begin{array}{c} \sigma \cdot p \\ m_q \end{array} \right),$$

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

and we can use the expression

$$\frac{|c_{\lambda\mu}^{\nu}}{A_{d\lambda}} = \frac{6m_{\Delta M} m_{d\Delta_\pi}}{f^2 B A_{d\lambda} m_{2d\Delta_\pi}},$$

$$\frac{|c_{\lambda\mu}^{\nu}}{A_{d\lambda}^{2d\Delta_\pi}} = \frac{2m_{\Delta M} m_{d\Delta_\pi}}{f^2 B A_{d\lambda} m_{2d\Delta_\pi}^2},$$

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 is 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 employed this model with
certain modifications to explain the anomaly in $A_{FB}$ observed at LEP [18]. Barger et al. studied $Z'$ mediated flavor changing neutral currents in $B$-meson decays [19], $B_s - B_s$ mixing [20] and $B \rightarrow K \pi$ 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 sub-groups respectively. The vacuum expectation values of the three Higgs bosons break the symmetry. The symmetry breaking patterns are depicted in literature. The introduction of a scalar field $\phi$ causes $Z_0$ in the standard model to mix with a new gauge boson $Z_R$, then $Z$, $Z'$ are the mass eigen-states.

For the neutral sector the Lagrangian is

\[
\mathcal{L} = -\frac{g_L}{2 \cos \theta_W} \bar{\gamma}^\mu (g_V - g_A \gamma_5) q (\cos \xi Z \bar{Z}_\mu - \sin \xi Z \bar{Z}'_\mu) + \frac{g_V}{2} \tan \theta_R \left( \frac{\bar{Z}}{g_L} q L - \frac{1}{3} \bar{Z}_R \gamma^\mu u_R \right) - \frac{2}{3} \bar{d}_R \gamma^\mu d_R \\
+ \frac{g_V}{2} (\tan \theta_R + \cot \theta_R) (\sin \xi Z \bar{Z}_\mu - \cos \xi Z \bar{Z}'_\mu) (\gamma^\mu d_R) - V_{Rti}^{u} V_{Rj}^{u} \bar{Z}_R \gamma^\mu u_R. \tag{11}
\]

Here $\theta_W$ is the electroweak mixing angle ($\tan \theta_W = \frac{g_W}{g_L}$), $\theta_R$ parameterizes the relative strength of the right-handed interaction ($\frac{\tan \theta_R - \frac{1}{3} \tan \theta_R}{g_L}$), $\xi Z$ is the $Z - Z'$ mixing angle and $V_{Rti}^{u}$ are two unitary matrices that rotate the right-handed up-(down)-type quarks from the weak eigen-states to the mass eigen-states. Note that we use current notation for Pati-Salam model, and only third family couples to $SU(2)_R$ in this model.

In the $Z'$ model, inclusive decay of doubly charmed baryons into light quarks $ccq \rightarrow uuq$ 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 $\Xi_{cc}^+$, we obtain

\[
T = \sum_{\text{spin}} \int d^3p_d d^3p_\lambda (2\pi)^3 \frac{E_{pd}}{m_{pd}} \times [\bar{u}_u(p_1, s_1) \gamma_\mu (1 - \gamma_5) u_c(p'_1, s'_1) \times \bar{u}_u(p_2, s_2) \gamma_\nu (1 - \gamma_5) u_c(p'_2, s'_2)]
\times (\gamma^\mu \gamma^\nu + g_L \tan \theta_W (\tan \theta_R + \cot \theta_R) \cos \xi Z)
\frac{3 M_{Z'}^4}{(p^2 + M_{Z'}^2)^2} \times V_{Rti}^{u} V_{Rj}^{u} \bar{Z}_R \gamma^\mu u_R \times N_{\Xi_{cc}^+} (p_\rho, p_\lambda). \tag{12}
\]

Then we have the final expression as

\[
T = \sum_{\text{spin}} \int d^3p_d d^3p_\lambda (2\pi)^3 \frac{E_{pd}}{m_{pd}} \times [\bar{u}_u(p_1, s_1) \gamma_\mu (1 - \gamma_5) u_c(p'_1, s'_1) \times \bar{u}_u(p_2, s_2) \gamma_\nu (1 - \gamma_5) u_c(p'_2, s'_2)]
\times (-ig^{\mu\nu}) \frac{G_F \sqrt{2} (V_{Rti}^{u} V_{Rj}^{u})^2}{16 M_{\Xi_{cc}^+} (2\pi)^4} \times N_{\Xi_{cc}^+} (p_\rho, p_\lambda). \tag{13}
\]

\[\Gamma(\Xi_{cc}^+(3520)^+ \rightarrow uu d) = \frac{\int_{a_1}^{a_2} d\rho_1^R \int_{b_1}^{b_2} d\rho_2^R \int_0^{2\pi} \int_{-\tan \theta}^{\tan \theta} d(\cos \theta) |T|^2}{16 M_{\Xi_{cc}^+} (2\pi)^4}, \tag{14}\]

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

\[
a_1 = 0, \quad a_2 = \frac{M_{\Xi_{cc}^+}^2 - (m_2 + m_3)^2 - m_1^2}{2 M_{\Xi_{cc}^+}^2}.
\]

\[
b_1 = \frac{1}{2 \sigma} \left[ \sigma (\tau + m_+ m_-) - \sqrt{p_1^2 (\tau - m_+^2)(\tau - m_-^2)} \right],
\]

\[
b_2 = \frac{1}{2 \tau} \left[ \sigma (\tau + m_+ m_-) + \sqrt{p_1^2 (\tau - m_+^2)(\tau - m_-^2)} \right],
\]

\[
\sigma = M_{\Xi_{cc}^+}^2 - p_1^2, \quad \tau = \sigma^2 - \sqrt{p_1^2}, \quad m_+ = m_2 \pm m_3.
\]

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 $\Xi_{cc}^+$, the expression is the same but only the mass of charm quark is replaced by that of bottom quark.

### III. NUMERICAL RESULTS

Now we present our numerical results.

Since only $\Xi_{cc}^+$ has been measured, in the later calculation, we use its measured mass as input, and for $\Xi_{bb}$ we will only illustrate the dependence of its decay rate on the parameters. The input parameters include: $G_F = 1.166 \times 10^{-5}$ GeV$^{-2}$, $m_c = 1.60$ GeV, $m_d = 0.43$ 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$^2$. Here, the light quark mass refers to the constituent mass.
TABLE I: The decay widths of $\Xi_{cc}(3520)^+ \to uud$ or $\Xi_{cc}^{++} \to uuu$ and $\Xi_{bb} \to ddd(sbd)$ corresponding to $d\Lambda = 3/2$ (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.

| Decay Width | Scalar | Vector | Scalar + Vector |
|-------------|--------|--------|-----------------|
| $\Gamma[\Xi_{cc}(3520)^+ \to uud]$ | $4.57 \times 10^{-18}$ | $1.11 \times 10^{-15}$ | $1.24 \times 10^{-15}$ |
| $\Gamma[\Xi_{bb} \to ddd]$ | $5.85 \times 10^{-20}$ | $1.65 \times 10^{-17}$ | $1.83 \times 10^{-17}$ |
| $\Gamma[\Xi_{bb} \to sdd]$ | $5.21 \times 10^{-20}$ | $1.23 \times 10^{-17}$ | $1.44 \times 10^{-17}$ |

A. The results in the Unparticle scenario

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

In this work, we also calculate the inclusive decay width of doubly bottomed baryon $\Xi_{bb}$. The mass of $\Xi_{bb}$ 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. Fig. 3 illustrates the dependence of the decay widths of $\Xi_{cc}(3520)^+ \to uud$ and $\Xi_{bb} \to ddd$ on $d\Lambda$, the three lines (solid, dashed and dotted) correspond to the contributions of scalar unparticle, vector unparticle and both on $d\Lambda$. It is noted here "both" means that at present we cannot determine whether the unparticle is a scalar or vector and it is also possible that both scalar and vector exist simultaneously. Thus we assume both of scalar and vector contribute and they interfere constructively. Definitely, it is worth 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 $V_{u*c*}V_{u*c*}^*$ is bound no more than $2.0 \times 10^{-4}$ [22]. In our calculation, we take their extreme values as $M_{Z'} = 500$ GeV and $V_{u*c*}V_{u*c*}^* = 2.0 \times 10^{-4}$, 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 $\Xi_{cc}^{+}$, therefore, it is hopeless to observe a non-zero branching ratio of $\Xi_{cc}^{+}$ 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, $\Xi_{cc}^{+}$ can decay into two-body final states via Fig. 4 (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. 4 (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.

By the order of magnitude estimation and with the SU(3) symmetry, the ratio of the amplitude of Fig. 4 (c)
(T_{SM}) to Fig. 2(b) \((T_{\text{un}} \text{ or } T_{Z'})\) is

\[
\frac{T_{SM}}{T_{Z'}} \approx \frac{3\pi a}{q^2} \sqrt{2}(V_{uu}^* V_{uu}^{*\prime})^2,
\]

(16)

where \(q\) is the momentum of unparticle or \(Z'\) (in the case of the SM, \(q^2 \ll \mathcal{M}_q^2\)) and can be neglected in the propagator. Because the contribution from smaller \(q^2\) is dominant, in the estimation, we set \(q^2 = 0.5\text{GeV}^2\). Since the whole case under consideration, may fall in the non-perturbative QCD region, as a rough estimate, we take \(a_s = 1\), \(V_{R_{tt}}^{*\prime} V_{R_{ttj}} = 2 \times 10^{-4}\) and \(G_F = 1.166 \times 10^{-5}\text{GeV}^{-2}\), we can get \(\frac{T_{SM}}{T_{Z'}} \approx 66\) (i.e. \(\frac{\Gamma_{Z'}}{\Gamma_{SM}} \approx 4300\)). Then we can obtain the ratio of decay widths \(\frac{\Gamma_{\text{unparticle}}}{\Gamma_{SM}} \approx 40\). This ratio indicates that for \(\Xi^+_{cc}\) the contribution of the SM is smaller than that of the unparticle scenario, but larger than that from the \(Z'\) model.

However, for \(\Xi^+_{cc}\), Fig. 1(c) does not contribute at all, so that the decay of \(\Xi^+_{cc}\) into charmless final states (or \(\Xi_{bb}\) into bottomless final states) is more appropriate for exploring new physics than \(\Xi_{cc}\). 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 \(\Xi^+_{cc}, \Xi^{++}_{cc}\), \(\Xi^0_{bb}, \Xi^0_{bb}\) 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 \(cc (bb) \rightarrow qq\) where \(q\) may be light quarks to occur at tree level. Thus one expects that these new models might result in non-zero 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, then secondly, the double-charmed baryon \(\Xi_{cc}\) 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, but their framework is clear, so that 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 by contraries, 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 \(\Xi_{cc}\) or double-bottomed \(\Xi_{bb}\) can only be produced at LHC, but not at B-factories.

The inclusive decays of doubly charmed baryons \(\Xi_{cc}(3520)^+, \Xi_{cc}^{++}\) and \(\Xi_{bb}^0, \Xi_{bb}^0\) are explored in unparticle and \(Z'\) scenarios. Our result indicates that the upper limit of the inclusive decay width of \(\Xi_{cc}^{++} \rightarrow wu\) is about \(10^{-15}\) GeV with \(d_q = 3/2\). For inclusive decay \(\Xi_{bb} \rightarrow ddd(ssd), \) the upper limit is at order of \(10^{-17}\) GeV. It is learnt 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 1.

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 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. [23], the authors estimate the number of \(\Xi_{cc}\) produced at LHCb as about \(10^3\). Since the available energy is much higher than the masses of \(\Xi_{cc}\) and \(\Xi_{bb}\), 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 \(\Xi_{cc}^{++} \rightarrow wu \rightarrow two \text{ hadrons}\) and \(\Xi_{bb} \rightarrow ddd(ssd) \rightarrow two \text{ hadrons}\) 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 the new physics is also not the unparticle and/or \( Z' \), but something else.

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