Doubly-charged scalar in rare decays of the B_c meson

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In this paper, we study the lepton number violation processes of B_c meson induced by possible doubly-charged scalars. Both the three-body decay channels and the four-body decay channels are considered. For the former, the $B_c^- \to D_s^+ \mu^- \mu^-$ channel has the largest branching fraction 9.19×10^{-23}, and for the later channels, the $B_c^- \to \bar{B}_s^0 \pi^+ \mu^- \mu^-$ channel has the largest branching ratio 1.03×10^{-27}. These results are too small to be within the current experimental precision. However, they can provide a comparison in theory with the similar cases involving an off-shell Majorana neutrino.

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I. INTRODUCTION

Doubly-charged Higgs bosons ($\Delta^{\pm\pm}$) have been predicted by the left-right symmetric models [1–3] as the third component of scalar triplets. If one keeps only this triplet as the new physics beyond the standard model (without introducing the right-handed neutrinos), the Type-II seesaw models [4–7] are achieved. This particle is phenomenologically interesting as it can decay to two same-sign charged leptons which indicates the lepton number violation (LNV). It has been searched extensively at the Large Hadron Collider (LHC). Until now, there is no evidence to show their existence, which sets constraints on their masses. For example, the latest result of ATLAS Collaboration shows the lower limit on $m(\Delta^{\pm\pm})$ is 770–870 GeV for final states with 100% decay to ee, $e\mu$, and $\mu\mu$ [8]. And for CMS Collaboration, this lower bound is between 800–820 GeV [9].

It is also interesting to investigate the low energy processes with doubly-charged Higgs boson as the intermediate state. Experimentally, the final particles come from the same vertex because the masses of W and Δ^{++} bosons are very large. Theoretically, the heavy bosons cannot on the mass shell, their contribution is reflected in effective interaction vertices [10]. Such low energy processes include the rare decays of top quark [11], τ lepton [11–13], or charged mesons [10,14,15]. Surely the branching ratios of these decay modes will be very small due to the large Higgs mass and small coupling constant. However, by comparing the experimental results of the branching ratios of the LNV processes with the theoretical predictions, one can get the lower bound on the parameters involved in the effective vertices [16]. One may argue that the Majorana neutrino can also lead to the LNV processes, such as neutrinoless double beta decays in low energy processes. Especially for Majorana neutrinos with masses around GeV scale, as they could be on-shell, the narrow width approximation (NWA) can be applied, which greatly enhances the decay widths of these processes [17]. However, it may also be possible that there are only three generations of light Dirac neutrinos in nature. If so, one has to find other mechanisms which could give the same neutrinoless double beta decay signal, and doubly charged Higgs boson will be such a possible alternative. If there are only three generations of light Majorana neutrinos, these LNV processes induced by them are greatly suppressed [18,19] and may have the same order of magnitude as the contribution of the doubly-charged Higgs, which makes the latter case important.

In Refs. [14] and [15], the $M_1 \to M_2 l_1^\pm l_2^\pm$ processes induced by the Δ^{++} with $M_1 = B^-, D^-, K^-$ are considered. In this work, we will study such processes of B_c^- meson. Moreover, we notice that the LNV four-body decay processes of heavy mesons with Majorana neutrinos have been extensively studied in theory [20–24], while such processes within the doubly-charged Higgs boson formalism have not been investigated yet. So a careful calculation of such channels will be a great supplement for the three-body decay modes. Experimentally, as LHCb will produce more and more B_c mesons, searching such decay channels will setting an experimental upper limit for the branching

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ratios, which can also be used to constrain the parameters of doubly-charged Higgs boson.

This work is organized as follows. In Sec. II, we present the theoretical formalism. Three-body decay processes and four-body decay processes are both considered. In Sec. III, we give the numerical results and discussions. Finally, the conclusion is given in Sec. IV. And some details for the calculation of the hadronic transition matrix element are presented in the Appendix.

II. THEORETICAL FORMALISM

The Lagrangian describing the interaction of doubly-charged scalars with standard model fermions has the form [10]

$$\mathcal{L}_{\rm int} = i h_{ij} \psi^T_{iL} C \sigma_2 \Delta \psi_{jL} + \text{H.c.}, \tag{1}$$

where ψ_{iL} is the two-component leptonic doublet; h_{ij} is the leptonic Yukawa coupling constant; $C = i\gamma^2\gamma^0$ is the charge conjugation matrix; σ_2 is the second Pauli matrix; Δ is the complex triplet in the 2×2 representation which we have defined as

$$\Delta = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix}. \tag{2}$$

The Lagrangian which describes the interaction of $\Delta^{++}/\Delta^{\pm}$ with W^- gauge boson and quarks has the following form [10,14]

$$\mathcal{L}'_{\rm int} = -\sqrt{2} g m_W s_\Delta \Delta^{++} W^{-\mu} W^-_\mu + \frac{\sqrt{2}}{2} g c_\Delta W^{-\mu} \Delta^- \overleftrightarrow{\partial}_\mu \Delta^{++} + \frac{i g s_\Delta}{\sqrt{2} m_W c_\Delta} \Delta^+ (m_{q'} \bar{q}_R q'_R - m_q \bar{q}_L q'_L) + \text{H.c.}, \tag{3}$$

where $s_\Delta = \sin\theta_\Delta$ and $c_\Delta = \cos\theta_\Delta$ with θ_Δ is the mixing angle between the usual $\mathrm{SU}(2)_L$ Higgs doublet and the assumed Higgs triplet. This mixing happens when the electroweak gauge symmetry is spontaneously broken as both the neutral components of the Higgs doublet and triplet acquire corresponding vacuum expectation values (VEV) v_H and v_Δ. The mixing angle is related to the VEVs by the relation $s_\Delta = \sqrt{2x^2/(1+2x^2)}$ with $x = \frac{v_\Delta}{v_H}$ [25]. From the experimental result of the ρ parameter, which is defined as [25]

$$\rho = \frac{M_W^2}{M_Z^2 \cos^2\theta_W} = \frac{1+2x^2}{1+4x^2}, \tag{4}$$

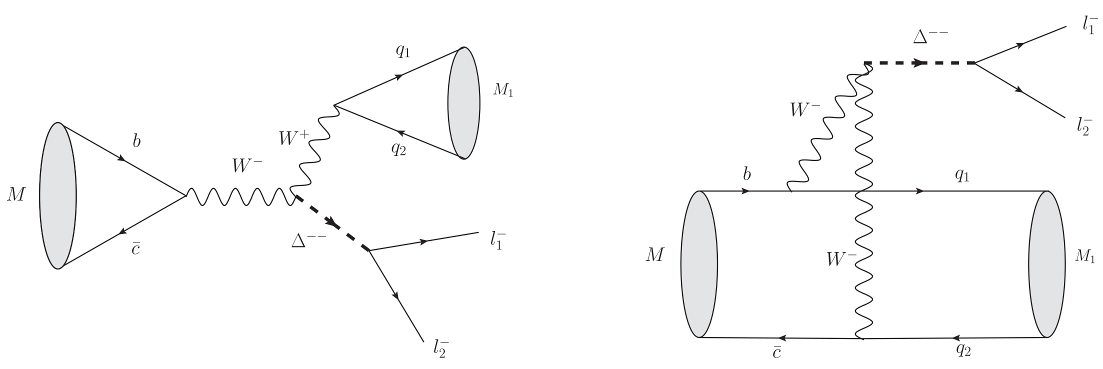
we can deduce the upper limit of x, and then set the upper limit of s_Δ.

A. The $B_c^- \to h^+ l_1^- l_2^-$ processes

The three-body decay process of B_c^- with lepton number violation is shown in Fig. 1. Actually, there are six other diagrams which contain $\Delta^{\pm}$. However, the contribution of those diagrams is very small compared with those of Fig. 1. This can be seen from that the parameters of the last two terms in Eq. (3) are very small compared with that of the first term. In Ref. [11], the ratio of the amplitudes with and without $\Delta^{\pm}$ is estimated to be less than 10^{-7}. So here we can safely neglect their contribution. The amplitude corresponding to the two diagrams in Fig. 1 is

$$\begin{aligned}\mathcal{M} &= \frac{g^3 s_\Delta h_{ij}}{8\sqrt{2} m_W^3 m_\Delta^2}\left(V^*_{cb} V_{q_1 q_2} + \frac{1}{3} V_{q_1 b} V^*_{c q_2}\right) \\ &\quad \times \langle h(p_1)|(\bar{c}b)_{V-A}(\bar{q}_1 q_2)_{V-A}|B_c^-(p)\rangle\langle\text{lepton}\rangle \\ &= \frac{g^3 s_\Delta h_{ij}}{8\sqrt{2} m_W^3 m_\Delta^2}\left(V^*_{cb} V_{q_1 q_2} + \frac{1}{3} V_{q_1 b} V^*_{c q_2}\right) \\ &\quad \times f_h f_{B_c} p \cdot p_1 \langle\text{lepton}\rangle,\end{aligned} \tag{5}$$

where we have used the definition $\langle h(p_1)|\bar{q}_1\gamma^\mu(1-\gamma_5) \times q_2|0\rangle = i f_h p_1^\mu$ with f_h being the decay constant of the final pseudoscalar meson. For the vector meson case, the definition $\langle h(p_1,\epsilon)|\bar{q}_1\gamma^\mu(1-\gamma_5)q_2|0\rangle = f_h M_1 \epsilon_1^\mu$ should



(a) (b)

FIG. 1. Feynman diagrams of the decay processes $B_c^- \to h^+ l_1^- l_2^-$.

be applied, and in Eq. (4), $p \cdot p_1$ should be changed to $M_1 p \cdot \epsilon_1$. We also defined $\langle \text{lepton} \rangle \equiv \bar{v}(k_2)(1-\gamma_5)u(k_1) - \bar{v}(k_1)(1-\gamma_5)u(k_2)$, where $u(k_i)$ and $v(k_i)$ are the spinors of charged leptons. The factor $\frac{1}{3}$ in the parentheses comes from the Fierz transformation. The squared amplitude can be written as

$$|\mathcal{M}|^2 = \sqrt{2}G_F^3 \left(\frac{s_\Delta h_{ij}}{M_\Delta^2}\right)^2 \left| V_{cb}^* V_{q_1 q_2} + \frac{1}{3} V_{q_1 b} V_{c q_2}^* \right|^2 \times f_h^2 f_{B_c}^2 |p \cdot p_1 \langle \text{lepton} \rangle|^2, \quad (6)$$

where we have used the definition $\frac{G_F}{\sqrt{2}} = \frac{g^2}{8 m_W^2}$.

The partial decay width can be achieved by finishing the phase space integral

$$\Gamma = \left(1 - \frac{1}{2}\delta_{l_1 l_2}\right) \frac{1}{512\pi^3 M^3} \int \frac{ds_{12}}{s_{12}} \lambda^{1/2}(M^2, s_{12}, M_1^2) \times \lambda^{1/2}(s_{12}, m_1^2, m_2^2) \int d\cos\theta_{12} |\mathcal{M}|^2, \quad (7)$$

where $s_{12} \equiv (k_1 + k_2)^2$; m_1 and m_2 are the masses of two charged lepton l_1 and l_2, respectively; the Källen function

$$\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz \quad (8)$$

is used; θ_{12} is the angle between the three-momenta $\vec{k}_{12} = \vec{k}_1 + \vec{k}_2$ and $\vec{K}_1$ (the later is the three-momentum of l_1 in the center-of-momentum frame of l_1 and l_2). $\delta_{l_1 l_2} = 0(1)$ when l_1 and l_2 are nonidentical (identical) leptons. The integral limits are

$$s_{12} \in [(m_1 + m_2)^2, \quad (M - M_1)^2], \quad \theta_{12} \in [0, \pi]. \quad (9)$$

B. The $B_c^- \to h_1^0 h_2^+ l_1^- l_2^-$ processes

For the $B_c^- \to J/\psi h_2^+ l_1^- l_2^-$ processes, when $h_2^+ = \pi^+$ or K^+, only the Feynman diagrams of Figs. 2(a) and 2(b) contribute; when $h_2^+ = D^+, D_s^+$, all the four diagrams of Fig. 2 give contribution, while that of Figs. 2(c) and 2(d) could be neglected as the $c\bar{c}$ pair production will be highly suppressed. So we only consider the contribution of Figs. 2(a) and 2(b). The corresponding amplitudes are written as

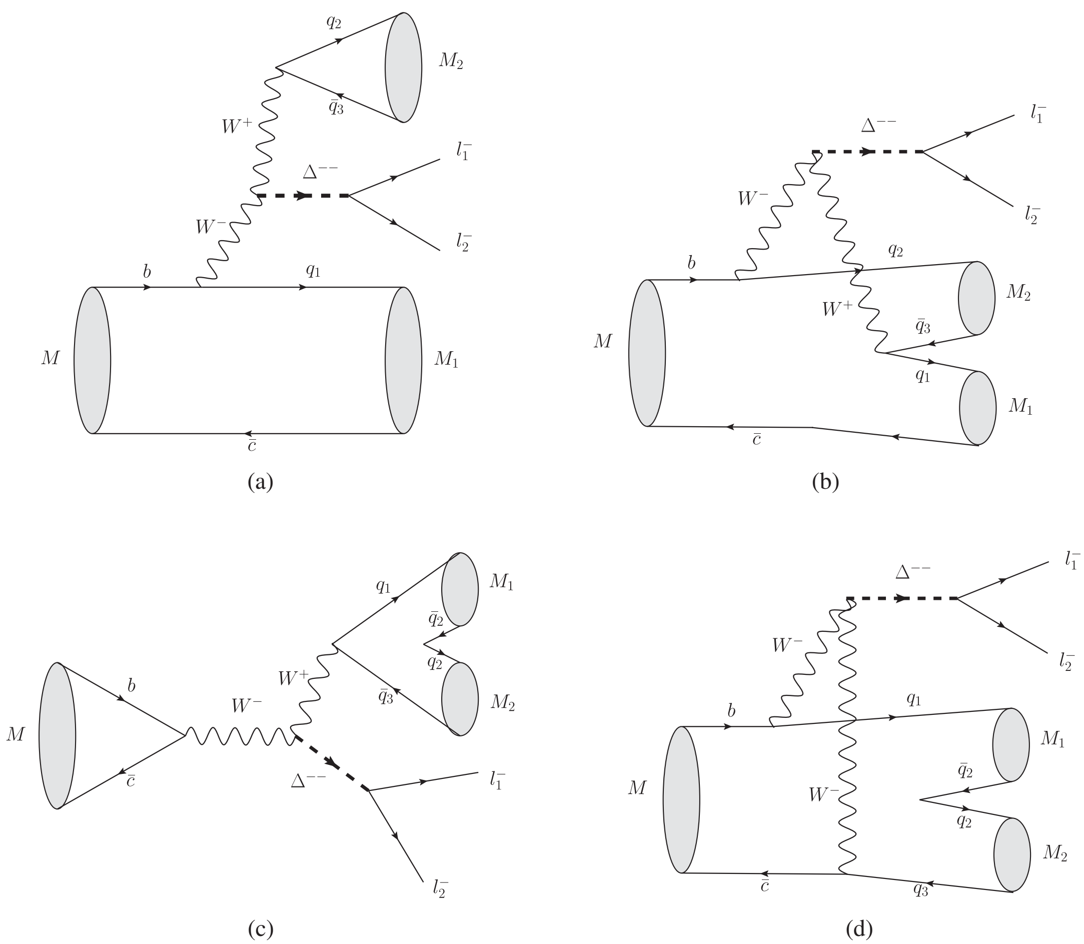


FIG. 2. Feynman diagrams of the decay processes $B_c^- \to h_1^0 h_2^+ l_1^- l_2^-$.

$$\begin{aligned}\mathcal{M}_A &= \frac{g^3}{8\sqrt{2}m_W^3} V_{cb} V_{q_2 q_3} \frac{s_\Delta h_{ij}}{m_\Delta^2} \langle J/\psi(p_1) h_2(p_2) | (\bar{c}b)_{V-A} (\bar{q}_2 q_3)_{V-A} | B_c^-(p) \rangle \langle \text{lepton} \rangle \\ &= \frac{g^3}{8\sqrt{2}m_W^3} V_{cb} V_{q_2 q_3} \frac{s_\Delta h_{ij}}{m_\Delta^2} f_{h_2} p_2^\mu \langle J/\psi(p_1) | \bar{c}\gamma_\mu (1-\gamma_5) b | B_c^-(p) \rangle \langle \text{lepton} \rangle, \end{aligned} \tag{10}$$

$$\begin{aligned}\mathcal{M}_B &= \frac{g^3}{8\sqrt{2}m_W^3} V_{q_2 b} V_{c q_3} \frac{s_\Delta h_{ij}}{m_\Delta^2} \langle J/\psi(p_1) h_2(p_2) | (\bar{q}_2 b)_{V-A} (\bar{c} q_3)_{V-A} | B_c^-(p) \rangle \langle \text{lepton} \rangle \\ &= \frac{g^3}{24\sqrt{2}m_W^3} V_{q_2 b} V_{c q_3} \frac{s_\Delta h_{ij}}{m_\Delta^2} f_{h_2} p_2^\mu \langle J/\psi(p_1) | \bar{c}\gamma_\mu (1-\gamma_5) b | B_c(p) \rangle \langle \text{lepton} \rangle, \end{aligned} \tag{11}$$

where we have used the Fierz transformation in $\mathcal{M}_B$. Here we only give the results when h_2 is a pseudoscalar meson. If h_2 is a vector meson, $f_{h_2} p_2^\mu$ should be replaced by $M_2 f_{h_2} \epsilon_2^\mu$ Finally, we get the transition amplitude

$$\mathcal{M} = \frac{g^3 s_\Delta h_{ij}}{8\sqrt{2}m_W^3 m_\Delta^2} \left(V_{cb} V_{q_2 q_3} + \frac{1}{3} V_{q_2 b} V_{c q_3} \right) \times f_{h_2} p_2^\mu \langle J/\psi(p_1) | \bar{c}\gamma_\mu (1-\gamma_5) b | B_c^-(p) \rangle \langle \text{lepton} \rangle. \tag{12}$$

The hadronic transition matrix can be expressed as [26]

$$\begin{aligned}\langle J/\psi(p_1) | V^\mu | B_c^-(p) \rangle &= -i \frac{2}{M+M_1} f_V(Q^2) \epsilon^{\mu \epsilon^* p p_1}, \\ \langle J/\psi(p_1) | A^\mu | B_c^-(p) \rangle &= f_1(Q^2) \frac{\epsilon^* \cdot p}{M+M_1} (p+p_1)^\mu \\ &\quad + f_2(Q^2) \frac{\epsilon^* \cdot p}{M+M_1} (p-p_1)^\mu \\ &\quad + f_0(Q^2)(M+M_1) \epsilon^{*\mu}, \end{aligned} \tag{13}$$

where $Q = p - p_1$, f_V and f_i ($i = 0, 1, 2$) are form factors.

For the $B_c^- \to \bar{D}^{(*)0} h_2^+ l_1^- l_2^-$ processes, h_2^+ can also be π^+, K^+, D^+, or D_s^+. For the same reason as the J/ψ case, the contribution of Figs. 2(c) and 2(d) or D^+ and D_s^+ situations are also neglected. The transition amplitude can be written as

$$\mathcal{M} = \frac{g^3 s_\Delta h_{ij}}{8\sqrt{2}m_W^3 m_\Delta^2} \left(V_{ub} V_{q_2 q_3} + \frac{1}{3} V_{q_2 b} V_{u q_3} \right) \times f_{h_2} p_2^\mu \langle \bar{D}^{(*)0}(p_1) | \bar{c}\gamma_\mu (1-\gamma_5) b | B_c^-(p) \rangle \langle \text{lepton} \rangle. \tag{14}$$

For $\bar{D}^{*0}$, the hadronic transition matrix element is parameterized in the same way as Eq. (12). For $\bar{D}^0$, it is parameterized as [26]

$$\langle \bar{D}^0(p_1) | V^\mu | B_c^-(p) \rangle = f_+(Q^2)(p+p_1)^\mu + f_-(Q^2)(p-p_1)^\mu, \tag{15}$$

where $f_\pm$ are form factors.

The phase space integral for four body decay processes can be expressed as

$$\Gamma = \left(1 - \frac{1}{2}\delta_{l_1 l_2}\right) \int \frac{ds_{12}}{s_{12}} \int \frac{ds_{34}}{s_{34}} \int d\cos\theta_{12} \int d\cos\theta_{34} \times \int d\phi \mathcal{K} |\mathcal{M}|^2, \tag{16}$$

where $s_{12} = (p_1 + p_2)^2$, $s_{34} = (k_1 + k_2)^2$. θ_{12} is the angle between the three-momenta $\vec{p}_{12} = \vec{p}_1 + \vec{p}_2$ and $\vec{P}_1$ (the later is the three-momentum of h_1 in the center-of-momentum frame of h_1 and h_2). θ_{34} is the angle between the three-momenta $\vec{k}_{12} = \vec{k}_1 + \vec{k}_2$ and $\vec{K}_1$ (the later is the three-momentum of l_1 in the center-of-momentum frame of l_1 and l_2). ϕ is the angle between the decay planes $\Sigma(h_1 h_2)$ and $\Sigma(l_1 l_2)$. The factor $\mathcal{K}$ has the expression

$$\mathcal{K} = \frac{1}{2^{15}\pi^6 M^3} \lambda^{1/2}(M^2, s_{12}, s_{34}) \lambda^{1/2}(s_{12}, M_1^2, M_2^2) \times \lambda^{1/2}(s_{34}, m_1^2, m_2^2). \tag{17}$$

The integral limits are

$$\begin{aligned} &s_{12} \in [(M_1 + M_2)^2, (M - m_1 - m_2)^2], \\ &s_{34} \in [(m_1 + m_2)^2, (M - \sqrt{s_{12}})^2], \\ &\phi \in [0, 2\pi], \theta_{12} \in [0, \pi], \theta_{34} \in [0, \pi]. \end{aligned} \tag{18}$$

III. NUMERICAL RESULTS

Here we present some parameters used in the calculation. The lifetime of B_c^- meson is 0.507×10^{-12} s [27]. The decay constants used here are as follows: $f_{B_c} = 0.322$ GeV [28], $f_\pi = 130.4$ MeV, $f_K = 156.2$ MeV, $f_D = 204.6$ MeV, and $f_{D_s} = 257.5$ MeV [27], $f_\rho = 0.205$ GeV, $f_{K^*} = 0.217$ GeV [29], $f_{D^*} = 0.340$ GeV,

TABLE I. The upper limit of the branching ratios of three-body decay channels of B_c^-.

Decay channel	Br	Decay channel	Br
$B_c^- \to \pi^+ e^- e^-$	7.50×10^{-24}	$B_c^- \to \rho^+ e^- e^-$	1.52×10^{-23}
$B_c^- \to \pi^+ \mu^- \mu^-$	1.92×10^{-23}	$B_c^- \to \rho^+ \mu^- \mu^-$	3.88×10^{-23}
$B_c^- \to \pi^+ e^- \mu^-$	2.47×10^{-33}	$B_c^- \to \rho^+ e^- \mu^-$	4.99×10^{-33}
$B_c^- \to K^+ e^- e^-$	5.94×10^{-25}	$B_c^- \to K^{*+} e^- e^-$	8.88×10^{-25}
$B_c^- \to K^+ \mu^- \mu^-$	1.52×10^{-24}	$B_c^- \to K^{*+} \mu^- \mu^-$	2.27×10^{-24}
$B_c^- \to K^+ e^- \mu^-$	1.95×10^{-34}	$B_c^- \to K^{*+} e^- \mu^-$	2.92×10^{-34}
$B_c^- \to D^+ e^- e^-$	1.21×10^{-24}	$B_c^- \to D^{*+} e^- e^-$	1.35×10^{-24}
$B_c^- \to D^+ \mu^- \mu^-$	3.10×10^{-24}	$B_c^- \to D^{*+} \mu^- \mu^-$	3.45×10^{-24}
$B_c^- \to D^+ e^- \mu^-$	3.98×10^{-34}	$B_c^- \to D^{*+} e^- \mu^-$	4.44×10^{-34}
$B_c^- \to D_s^+ e^- e^-$	3.59×10^{-23}	$B_c^- \to D_s^{*+} e^- e^-$	2.86×10^{-23}
$B_c^- \to D_s^+ \mu^- \mu^-$	9.19×10^{-23}	$B_c^- \to D_s^{*+} \mu^- \mu^-$	7.30×10^{-23}
$B_c^- \to D_s^+ e^- \mu^-$	1.18×10^{-32}	$B_c^- \to D_s^{*+} e^- \mu^-$	9.38×10^{-33}

and $f_{D_s^*} = 0.375$ GeV [30]. The quark masses used here are: $m_b = 4.96$ GeV, $m_c = 1.62$ GeV, $m_s = 1.50$ GeV, $m_d = 0.311$ GeV, and $m_u = 0.305$ GeV.

In the expressions of the transition amplitudes, a factor $\frac{s_\Delta h_{ij}}{M_\Delta^2}$ is contained. To estimate the upper limit of the decay width, we have to give the lower limit of the mass of the doubly-charged Higgs boson and the upper limit of the coupling constant h_{ij}. If we take the same values in Ref. [14], that is

$$\begin{aligned} h_{ee}^2 &= 9.7 \times 10^{-6}\ \mathrm{GeV}^{-2} M_\Delta^2, \\ h_{\mu\mu}^2 &= 2.5 \times 10^{-5}\ \mathrm{GeV}^{-2} M_\Delta^2, \\ h_{\mu e}^2 &= 1.6 \times 10^{-15}\ \mathrm{GeV}^{-2} M_\Delta^2, \\ s_\Delta &< 0.0056, \end{aligned} \tag{19}$$

and set $M_\Delta \sim 1000$ GeV considering the latest results in Refs. [8,9], then we can get $\left(\frac{s_\Delta h_{ij}}{M_\Delta^2}\right)^2 < 3.0 \times 10^{-16}\ \mathrm{GeV}^{-4}$ (for ee), $7.8 \times 10^{-16}\ \mathrm{GeV}^{-4}$ (for $\mu\mu$), $5.0 \times 10^{-26}\ \mathrm{GeV}^{-4}$ (for μe).

The branching ratios of three-body decay channels is presented in Table I. The $D_s^+ \mu^- \mu^-$ channel has the largest value which is 9.19×10^{-23}. One can see it is of the same order of magnitude as those of the D or B cases in Ref. [14]. The large difference between the channels with $e^- e^-$, $\mu^- \mu^-$, and $e^- \mu^-$ mainly comes from the difference of h_{ij}^2. If we set it to 1, then the ratios of these three channels will approximately be $1:1:2$, where 2 comes from that for the $e^- \mu^-$ channel we take $\delta_{l_1 l_2} = 0$.

In Table II, we give the branching ratios of four-body decay channels with J/ψ as one of the final mesons. Compared with the three-body decay channels, the branching ratios here are several orders smaller. Actually, most of the suppression comes from the phase space integral. We can estimate this as follows: from Eqs. (6) and (15) one can see that the ratio of the constants is $(2^6\pi^3)^{-1} = 5.0 \times 10^{-4}$, which provide most of the difference between Table I and II. The channels which have the largest upper limit are $J/\psi D_s^{*+} l^- l^-$, which are about 10^{-28}. One notices that, in Refs. [22,31], B_c four-body decays with a GeV scale Majorana neutrino are calculated. There Fig. 2(b) gives negligible contribution. Here this diagram is just color suppressed, while its contribution can have the same order of magnitude as that of Fig. 2(a).

For the $B_c^- \to \bar{D}^{(*)0} h_2^+ l_1^- l_2^-$ channels, the results are given in Tables III and IV. The largest upper limit of decay widths for these channels is of the order of 10^{-28}. As the final states contain $\bar{c}$, only Figs. 2(a) and 2(b) contribute to the channels with $h_2^+ = \pi^+, K^+$. For the channels with $h_2^+ = D^+, D_s^+$, (c) and (d) also give contribution, while they are neglected for the reason above. One notices that the decay widths in Table IV are about one order less than those in Table III. This is different with the semi-leptonic decay channels of B [27], where the $D^{*0} l^- \nu_l$ channel has larger width than that of $D^0 l^- \nu_l$. In Tables V and VI, we present

TABLE II. The upper limit of the branching ratios of B_c^- four-body decays induced by the $(\bar{c}b)_{V-A}(\bar{q}_1 q_2)_{V-A}$ current.

Decay channel	Br	Decay channel	Br
$B_c^- \to J/\psi \pi^+ e^- e^-$	6.81×10^{-30}	$B_c^- \to J/\psi \rho^+ e^- e^-$	2.81×10^{-29}
$B_c^- \to J/\psi \pi^+ \mu^- \mu^-$	1.86×10^{-29}	$B_c^- \to J/\psi \rho^+ \mu^- \mu^-$	7.94×10^{-29}
$B_c^- \to J/\psi \pi^+ e^- \mu^-$	2.30×10^{-39}	$B_c^- \to J/\psi \rho^+ e^- \mu^-$	9.67×10^{-39}
$B_c^- \to J/\psi K^+ e^- e^-$	1.02×10^{-30}	$B_c^- \to J/\psi K^{*+} e^- e^-$	2.12×10^{-30}
$B_c^- \to J/\psi K^+ \mu^- \mu^-$	2.82×10^{-30}	$B_c^- \to J/\psi K^{*+} \mu^- \mu^-$	5.98×10^{-30}
$B_c^- \to J/\psi K^+ e^- \mu^-$	3.46×10^{-40}	$B_c^- \to J/\psi K^{*+} e^- \mu^-$	7.27×10^{-40}
$B_c^- \to J/\psi D^+ e^- e^-$	2.94×10^{-30}	$B_c^- \to J/\psi D^{*+} e^- e^-$	5.92×10^{-30}
$B_c^- \to J/\psi D^+ \mu^- \mu^-$	8.40×10^{-30}	$B_c^- \to J/\psi D^{*+} \mu^- \mu^-$	1.70×10^{-29}
$B_c^- \to J/\psi D^+ e^- \mu^-$	1.00×10^{-39}	$B_c^- \to J/\psi D^{*+} e^- \mu^-$	2.02×10^{-39}
$B_c^- \to J/\psi D_s^+ e^- e^-$	8.18×10^{-29}	$B_c^- \to J/\psi D_s^{*+} e^- e^-$	1.21×10^{-28}
$B_c^- \to J/\psi D_s^+ \mu^- \mu^-$	2.34×10^{-28}	$B_c^- \to J/\psi D_s^{*+} \mu^- \mu^-$	3.48×10^{-28}
$B_c^- \to J/\psi D_s^+ e^- \mu^-$	2.78×10^{-38}	$B_c^- \to J/\psi D_s^{*+} e^- \mu^-$	4.12×10^{-38}

TABLE III. The upper limit of the branching ratios of B_c^- four-body decays induced by the $(\bar{u}b)_{V-A}(\bar{q}_1 q_2)_{V-A}$ current with $\bar{D}^0$ as one of the final mesons.

Decay channel	Br	Decay channel	Br
$B_c^- \to \bar{D}^0\pi^+e^-e^-$	6.73×10^{-30}	$B_c^- \to \bar{D}^0\rho^+e^-e^-$	1.43×10^{-29}
$B_c^- \to \bar{D}^0\pi^+\mu^-\mu^-$	1.85×10^{-29}	$B_c^- \to \bar{D}^0\rho^+\mu^-\mu^-$	3.96×10^{-29}
$B_c^- \to \bar{D}^0\pi^+\mu^-e^-$	2.29×10^{-39}	$B_c^- \to \bar{D}^0\rho^+\mu^-e^-$	4.88×10^{-39}
$B_c^- \to \bar{D}^0K^+e^-e^-$	5.39×10^{-31}	$B_c^- \to \bar{D}^0K^{*+}e^-e^-$	8.12×10^{-31}
$B_c^- \to \bar{D}^0K^+\mu^-\mu^-$	1.49×10^{-30}	$B_c^- \to \bar{D}^0K^{*+}\mu^-\mu^-$	2.25×10^{-30}
$B_c^- \to \bar{D}^0K^+\mu^-e^-$	1.84×10^{-40}	$B_c^- \to \bar{D}^0K^{*+}\mu^-e^-$	2.77×10^{-40}
$B_c^- \to \bar{D}^0D^+e^-e^-$	1.07×10^{-28}	$B_c^- \to \bar{D}^0D^{*+}e^-e^-$	1.03×10^{-28}
$B_c^- \to \bar{D}^0D^+\mu^-\mu^-$	3.01×10^{-28}	$B_c^- \to \bar{D}^0D^{*+}\mu^-\mu^-$	2.89×10^{-28}
$B_c^- \to \bar{D}^0D^+\mu^-e^-$	3.67×10^{-38}	$B_c^- \to \bar{D}^0D^{*+}\mu^-e^-$	3.53×10^{-38}
$B_c^- \to \bar{D}^0D_s^+e^-e^-$	2.50×10^{-29}	$B_c^- \to \bar{D}^0D_s^{*+}e^-e^-$	1.71×10^{-29}
$B_c^- \to \bar{D}^0D_s^+\mu^-\mu^-$	7.08×10^{-29}	$B_c^- \to \bar{D}^0D_s^{*+}\mu^-\mu^-$	4.81×10^{-29}
$B_c^- \to \bar{D}^0D_s^+\mu^-e^-$	8.61×10^{-39}	$B_c^- \to \bar{D}^0D_s^{*+}\mu^-e^-$	5.85×10^{-39}

TABLE IV. The upper limit of the branching ratios of B_c^- four-body decays induced by the $(\bar{u}b)_{V-A}(\bar{q}_1 q_2)_{V-A}$ current with $\bar{D}^{*0}$ as one of the final mesons.

Decay channel	Br	Decay channel	Br
$B_c^- \to \bar{D}^{*0}\pi^+e^-e^-$	9.78×10^{-32}	$B_c^- \to \bar{D}^{*0}\rho^+e^-e^-$	2.36×10^{-31}
$B_c^- \to \bar{D}^{*0}\pi^+\mu^-\mu^-$	2.69×10^{-31}	$B_c^- \to \bar{D}^{*0}\rho^+\mu^-\mu^-$	6.61×10^{-31}
$B_c^- \to \bar{D}^{*0}\pi^+\mu^-e^-$	3.32×10^{-41}	$B_c^- \to \bar{D}^{*0}\rho^+\mu^-e^-$	8.13×10^{-41}
$B_c^- \to \bar{D}^{*0}K^+e^-e^-$	1.38×10^{-32}	$B_c^- \to \bar{D}^{*0}K^{*+}e^-e^-$	2.06×10^{-32}
$B_c^- \to \bar{D}^{*0}K^+\mu^-\mu^-$	3.81×10^{-32}	$B_c^- \to \bar{D}^{*0}K^{*+}\mu^-\mu^-$	5.75×10^{-32}
$B_c^- \to \bar{D}^{*0}K^+\mu^-e^-$	4.70×10^{-42}	$B_c^- \to \bar{D}^{*0}K^{*+}\mu^-e^-$	7.05×10^{-42}
$B_c^- \to \bar{D}^{*0}D^+e^-e^-$	8.45×10^{-30}	$B_c^- \to \bar{D}^{*0}D^{*+}e^-e^-$	1.68×10^{-29}
$B_c^- \to \bar{D}^{*0}D^+\mu^-\mu^-$	2.39×10^{-29}	$B_c^- \to \bar{D}^{*0}D^{*+}\mu^-\mu^-$	4.77×10^{-29}
$B_c^- \to \bar{D}^{*0}D^+\mu^-e^-$	2.91×10^{-39}	$B_c^- \to \bar{D}^{*0}D^{*+}\mu^-e^-$	5.78×10^{-39}
$B_c^- \to \bar{D}^{*0}D_s^+e^-e^-$	2.02×10^{-30}	$B_c^- \to \bar{D}^{*0}D_s^{*+}e^-e^-$	3.08×10^{-30}
$B_c^- \to \bar{D}^{*0}D_s^+\mu^-\mu^-$	5.74×10^{-30}	$B_c^- \to \bar{D}^{*0}D_s^{*+}\mu^-\mu^-$	8.77×10^{-30}
$B_c^- \to \bar{D}^{*0}D_s^+\mu^-e^-$	6.97×10^{-40}	$B_c^- \to \bar{D}^{*0}D_s^{*+}\mu^-e^-$	1.06×10^{-39}

TABLE V. The upper limit of the branching ratios of B_c^- four-body decays induced by the $(\bar{q}_1 c)_{V-A}(\bar{q}_2 q_3)_{V-A}$ current with $\bar{B}_q^0$ as one of the final mesons.

Decay channel	Br	Decay channel	Br
$B_c^- \to \bar{B}^0\pi^+e^-e^-$	8.82×10^{-29}	$B_c^- \to \bar{B}^0\rho^+e^-e^-$	2.36×10^{-30}
$B_c^- \to \bar{B}^0\pi^+\mu^-\mu^-$	2.23×10^{-28}	$B_c^- \to \bar{B}^0\rho^+\mu^-\mu^-$	1.25×10^{-32}
$B_c^- \to \bar{B}^0\pi^+\mu^-e^-$	2.73×10^{-38}	$B_c^- \to \bar{B}^0\rho^+\mu^-e^-$	1.79×10^{-40}
$B_c^- \to \bar{B}^0K^+e^-e^-$	6.82×10^{-29}	$B_c^- \to \bar{B}^0K^{*+}e^-e^-$	9.97×10^{-32}
$B_c^- \to \bar{B}^0K^+\mu^-\mu^-$	1.60×10^{-28}	$B_c^- \to \bar{B}^0K^{*+}\mu^-\mu^-$	–
$B_c^- \to \bar{B}^0K^+\mu^-e^-$	1.95×10^{-38}	$B_c^- \to \bar{B}^0K^{*+}\mu^-e^-$	–
$B_c^- \to \bar{B}_s^0\pi^+e^-e^-$	4.29×10^{-28}	$B_c^- \to \bar{B}_s^0\rho+e^-e^-$	1.96×10^{-30}
$B_c^- \to \bar{B}_s^0\pi^+\mu^-\mu^-$	1.03×10^{-27}	$B_c^- \to \bar{B}_s^0\rho^+\mu^-\mu^-$	–
$B_c^- \to \bar{B}_s^0\pi^+\mu^-e^-$	1.28×10^{-37}	$B_c^- \to \bar{B}_s^0\rho^+\mu^-e^-$	4.96×10^{-42}
$B_c^- \to \bar{B}_s^0K^+e^-e^-$	3.09×10^{-29}	$B_c^- \to \bar{B}_s^0K^*+e^-e^-$	2.51×10^{-35}
$B_c^- \to \bar{B}_s^0K^+\mu^-\mu^-$	6.12×10^{-29}	$B_c^- \to \bar{B}_s^0K^{*+}\mu^-\mu^-$	–
$B_c^- \to \bar{B}_s^0K^+\mu^-e^-$	7.93×10^{-39}	$B_c^- \to \bar{B}_s^0K^{*+}\mu^-e^-$	–

TABLE VI. The upper limit of the branching ratios of B_c^- four-body decays induced by the $(\bar{q}_1 c)_{V-A}(\bar{q}_2 q_3)_{V-A}$ current with $\bar{B}_q^{*0}$ as one of the final mesons.

Decay channel	Br	Decay channel	Br
$B_c^- \to \bar{B}^{*0}\pi^+e^-e^-$	3.15×10^{-30}	$B_c^- \to \bar{B}_s^{*0}\pi^+e^-e^-$	3.65×10^{-29}
$B_c^- \to \bar{B}^{*0}\pi^+\mu^-\mu^-$	8.69×10^{-30}	$B_c^- \to \bar{B}_s^{*0}\pi^+\mu^-\mu^-$	1.01×10^{-28}
$B_c^- \to \bar{B}^{*0}\pi^+\mu^-e^-$	1.03×10^{-39}	$B_c^- \to \bar{B}_s^{*0}\pi^+\mu^-e^-$	1.19×10^{-38}
$B_c^- \to \bar{B}^{*0}K^+e^-e^-$	9.68×10^{-30}	$B_c^- \to \bar{B}_s^{*0}K^+e^-e^-$	1.31×10^{-29}
$B_c^- \to \bar{B}^{*0}K^+\mu^-\mu^-$	2.28×10^{-29}	$B_c^- \to \bar{B}_s^{*0}K^+\mu^-\mu^-$	3.07×10^{-29}
$B_c^- \to \bar{B}^{*0}K^+\mu^-e^-$	2.75×10^{-39}	$B_c^- \to \bar{B}_s^{*0}K^+\mu^-e^-$	3.71×10^{-39}

the branching ratios of channels with $\bar{B}^0_{(s)}h_1^+l_1^-l_2^-$ and $\bar{B}^{*0}_{(s)}h_1^+l_1^-l_2^-$ as the final states. The decay width of the $\bar{B}_s^0\pi^+l^-l^-$ channel has the largest upper limit of 10^{-27}, which mainly due to the large Cabibbo-Kobayashi-Maskawa (CKM) matrix elements.

Here three things should be mentioned to the four-body decay channels. First, except the channels calculated here, there are also some other channels which can only be realized through Figs. 2(c) and 2(d), such as $D^{(*)0}h_2^+l_1^-l_2^-$ and $\pi^+\pi^0l_1^-l_2^-$ channels. They are not considered here. Second, the QCD corrections are not included. But it is easy to be added if only Figs. 2(a) and 2(b) contribute, which is similar to that of the two-body nonleptonic decay channels of the B_c meson. Third, the final state interactions (FSI) are not considered, since it will not greatly change the results' order of magnitude.

IV. CONCLUSIONS

We have studied the doubly-charged Higgs boson induced lepton number violation processes of B_c meson. Both the three-body decay channels and four-body decay channels are considered. For the former, the largest value of the branching fraction is 9.19×10^{-23}, which comes from the $D_s^+\mu^-\mu^-$ channel. For the later ones, the branching fractions are of the order of $10^{-27}\sim10^{-42}$. The largest value comes from the $\bar{B}_s^0\pi^+\mu^-\mu^-$ channel. But they are still three orders smaller than the smallest value of three-body decay channels. The branching ratios of these channels are much smaller than the experimental precision, which makes them impossible to be achieved in the current experiments. However, on the one hand, our work is a helpful supplement to such studies in other works, such as the neutrinoless double beta decay processes of K, D_q, and B_q. Our results show that the partial widths of the three-body decay channels of the B_c meson have the same order with those of other mesons, which makes such studies share the same interests with other works. On the other hand, the four-body decays of such mesons with the doubly-charged Higgs boson have not been investigated before. Although their branching fractions are even smaller due to the phase space integral, they have the similar order of magnitude with those involving a off-shell Majorana neutrino. This will be important in theory if no on-shell Majorana neutrino is found in the future.

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APPENDIX: THE HADRONIC TRANSITION AMPLITUDE

The hadronic transition amplitude can be written as [32]

$$\langle h_1^0(p_1)|\bar{q}_1\gamma^\mu(1-\gamma_5)b|B_c^-(p)\rangle = \int\frac{d^3q}{(2\pi)^3}\mathrm{Tr}\left[\frac{\not{P}}{M}\overline{\varphi_{p_1}^{++}}(\vec{q}_1)\gamma_\mu(1-\gamma_5)\varphi_p^{++}(\vec{q})\right], \quad \text{(A1)}$$

where $\vec{q}$ and $\vec{q}_1$ are the relative momenta of B_c^- and h_1^0 mesons, respectively. $\varphi_p^{++}(\vec{q})$ and $\varphi_{p_1}^{++}(\vec{q}_1)$ are the positive energy parts of the wave functions of the initial and final heavy mesons, respectively, which have the following forms [33]

$$\varphi_{0^-}^{++}(q_\perp) = \left[A_1(q_\perp)+\frac{\not{P}}{M}A_2(q_\perp) + \frac{\not{q}_\perp}{M}A_3(q_\perp)+\frac{\not{P}\not{q}_\perp}{M^2}A_4(q_\perp)\right]\gamma_5, \quad \text{(A2)}$$

where the coefficients are

$$A_1 = \frac{M}{2}\left[\frac{\omega_1+\omega_2}{m_1+m_2}f_1+f_2\right],$$
$$A_2 = \frac{M}{2}\left[f_1+\frac{m_1+m_2}{\omega_1+\omega_2}f_2\right],$$
$$A_3 = -\frac{M(\omega_1-\omega_2)}{m_1\omega_2+m_2\omega_1}A_1,$$
$$A_4 = -\frac{M(m_1+m_2)}{m_1\omega_2+m_2\omega_1}A_1. \quad \text{(A3)}$$

In the above equation, m_1 and m_2 are respectively the masses of quark and antiquark inside the meson. ω_i is defined as $\sqrt{m_i^2+\vec{q}^2}$. f_1 and f_2 are functions of $\vec{q}^2$.

For the 1^- state, the positive energy part of the wave function has the form [33]

$$\varphi_{1^-}^{++}(q_\perp) = (q_\perp\cdot\epsilon)\left[B_1(q_\perp)+\frac{\not{P}}{M}B_2(q_\perp)+\frac{\not{q}_\perp}{M}B_3(q_\perp)+\frac{\not{P}\not{q}_\perp}{M^2}B_4(q_\perp)\right] + M\not{\epsilon}\left[B_5(q_\perp)+\frac{\not{P}}{M}B_6(q_\perp)+\frac{\not{q}_\perp}{M}B_7(q_\perp)+\frac{\not{P}\not{q}_\perp}{M^2}B_8(q_\perp)\right], \tag{A4}$$

where the coefficients are

$$\begin{aligned}
B_1 &= \frac{1}{2M(m_1\omega_2+m_2\omega_1)}[(\omega_1+\omega_2)q_\perp^2f_3+(m_1+m_2)q_\perp^2f_4+2M^2\omega_2f_5-2M^2m_2f_6],\\
B_2 &= \frac{1}{2M(m_1\omega_2+m_2\omega_1)}[(m_1-m_2)q_\perp^2f_3+(\omega_1-\omega_2)q_\perp^2f_4-2M^2m_2f_5+2M^2\omega_2f_6],\\
B_3 &= \frac{1}{2}\left[f_3+\frac{m_1+m_2}{\omega_1+\omega_2}f_4-\frac{2M^2}{m_1\omega_2+m_2\omega_1}f_6\right], \qquad B_4=\frac{1}{2}\left[\frac{\omega_1+\omega_2}{m_1+m_2}f_3+f_4-\frac{2M^2}{m_1\omega_2+m_2\omega_1}f_5\right],\\
A_5 &= \frac{1}{2}\left[f_5-\frac{\omega_1+\omega_2}{m_1+m_2}f_6\right], \qquad A_6=\frac{1}{2}\left[-\frac{m_1+m_2}{\omega_1+\omega_2}f_5+f_6\right],\\
B_7 &= \frac{M}{2}\frac{\omega_1-\omega_2}{m_1\omega_2+m_2\omega_1}\left[f_5-\frac{\omega_1+\omega_2}{m_1+m_2}f_6\right], \qquad B_8=\frac{M}{2}\frac{m_1+m_2}{m_1\omega_2+m_2\omega_1}\left[-f_5+\frac{\omega_1+\omega_2}{m_1+m_2}f_6\right].
\end{aligned} \tag{A5}$$

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