CP violation with a GeV-scale Majorana neutrino in $\Lambda_b \to (\Lambda_c^+, p^+)\pi^+\mu^-\mu^-$ decays

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We explore the possibility of CP violation in baryonic $\Lambda_b \to (\Lambda_c^+, p^+)\pi^+\mu^-\mu^-$ decays which are mediated by two Majorana sterile neutrinos and are $|\Delta L| = 2$ lepton-number-violating processes. Appreciable CP asymmetry can be obtained if there are two on-shell Majorana neutrinos that are quasidegenerate in mass with the mass difference of the order of average decay widths. We find that, given the present constraints on the heavy to light mixing element $|V_{\mu N}|$, the $\Lambda_b \to p^+\pi^+\mu^-\mu^-$ and $\Lambda_b \to \Lambda_c^+\pi^+\mu^-\mu^-$ decay rates are suppressed but could be within the experimental reach at the LHC. If searches of the modes are performed, then experimental limits on the rates can be translated to constraints on the Majorana neutrino mass m_N and heavy to light mixing element squared $|V_{\mu N}|^2$. We show that the constraints on the $(m_N, |V_{\mu N}|^2)$ parameter space coming from the $|\Delta L| = 2$ baryonic decays are complementary to the bounds coming from other processes.

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I. INTRODUCTION

The neutrino oscillation experiments confirm that at least two of the three active light neutrinos are massive [1–3]. This opens up the possibility of CP violation in the leptonic interactions which can be searched in neutrino oscillation experiments [4]. Leptonic CP violation can arise in the same manner as in the quark sector, namely, complex phases in the leptonic mixing matrix. Whether the neutrinos are of the Dirac or Majorana type, CP violation is expected in both cases. But two additional sources of CP-violating phases can arise if the neutrinos are Majorana rather than if they are Dirac. Majorana character plays an important role as far as the origin of the smallness of the active neutrino masses is concerned. If N_R is a Standard Model right-handed gauge-singlet (and, hence, sterile) neutrino, then the Standard Model allows both a Dirac mass term of the type $m_D(\bar{\nu}_L N_R + \mathrm{H.c.})$ and a Majorana term of the type $m_N N_R N_R$. Then, via a "seesaw" mechanism, one can have small active neutrino mass $m_\nu \sim m_D^2/m_N$ if m_D is at the electroweak scale or lower [5–11]. In the simplest version of the mechanism, the so-called type-I seesaw, the heavy electroweak singlets N_R of a few TeV are introduced that give rise to the light eigenstates $m_\nu \lesssim 1$ eV. However, low-energy seesaw mechanisms, where the sterile states N_R are in the range of a few hundreds of MeV to a few GeV, also have been proposed [12–18]. These so-called GeV-scale sterile neutrinos have several advantages: They could simultaneously explain the baryon asymmetry of the Universe [13,19–22] and can be experimentally searched at both the intensity and the energy frontiers.

An important distinguishing feature between Dirac and Majorana sterile neutrinos is that the latter participates in $|\Delta L| = 2$ lepton-number-violating (LNV) decays. For a light Majorana exchange, the neutrinoless double-beta decay ($0\nu\beta\beta$) [23–26] is one of the most sensitive probes of lepton number violation. But it was recently pointed out that, with the exchange of heavy Majorana neutrinos at the GeV scale, this rate can be enhanced [27,28]. Unfortunately, the $0\nu\beta\beta$ process is yet to be experimentally verified, and the best limits on the half-lives of different isotopes (^{76}Ge, ^{136}Xe, and ^{130}Te) come from several different experiments [29–33]. Because of the lack of evidence of LNV decay so far, it is imperative to pursue complementary search strategies. This is further reinforced by the fact that observation of $0\nu\beta\beta$ confirms lepton number violation only in the first family of neutrinos, and, to observe the same in other families, alternative processes must be investigated. Lepton-number-violating rare decays of mesons and baryons which are mediated by Majorana neutrinos are important in this regard. For light or heavy Majorana neutrino exchange, decay rates are too suppressed to be accessed by current experiments. But, if the Majorana mass is within a few

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hundred MeV to a few GeV, then the decay rates can be within the sensitivity reach of future experiments [34,35]. Because of ongoing searches of LNV processes at flavor factories including the LHC and Belle-II, there has been theoretical interest in the LNV decays of hadrons [34–66], τ-lepton decays [67–71], and different scattering processes [72–81]. The LHCb has searched for the process $B^- \to \pi^+\mu^-\mu^-$ [82] and the NA48/2 has searched for $K^- \to \pi^+\mu^-\mu^-$ [83], and these experiments provide stringent constraints on the heavy to light mixing matrix elements. With large integrated luminosity coming from Belle-II as well as the upgrade of the LHCb, sensitivity to $|\Delta L| = 2$ processes in mesons and baryons is expected to increase.

In this paper, we study lepton-number-violating four-body $\mathcal{B}_1 \to \mathcal{B}_2^{\mp}\pi^{\mp}\ell_1^{\pm}\ell_2^{\pm}$ decay, where $\mathcal{B}_1^0$ is Λ_b, $\mathcal{B}_2^+$ is either a Λ_c^+ or p^+, and ℓ_1 and ℓ_2 can, in general, be of different flavors. Previously, in Ref. [48], these decays were considered in a model involving single on-shell Majorana exchange at the GeV scale. We are interested in a scenario where the decays are mediated by the exchange of two almost degenerate Majorana neutrinos of mass in the range between a few hundred MeV to a few GeV so that they can be on shell. An interesting consequence of two-Majorana exchange is the possibility of CP violation. We show that the CP violation can be appreciable if the two Majoranas are almost degenerate with the mass difference of the order of decay widths, $\Delta m_N \sim \Gamma_N$. There are well-motivated models where quasidegenerate Majorana neutrinos in the range of a few hundreds of MeV to a few GeV are predicted [84]. We calculate the branching ratios for $\Delta m_N \sim \Gamma_N$ and find that, for the present experimental bound on $|V_{\mu N}|^2$, the $\Lambda_b \to (\Lambda_c^+, p^+)\pi^+\mu\mu$ rates might be within the reach of the LHC in the future. Even if the modes are not immediately seen, experimental limits on the decay rates can be used to obtain constraints on the neutrino mass m_N and the neutrino mixing matrix elements $|V_{\mu N}|^2$. Lepton-number-violating decays in baryons have been searched by the HyperCP, E653, and BESIII Collaborations in three-body $\Xi^- \to p\mu^-\mu^-$ [85], $\Lambda_c^+ \to \Sigma^-\mu^+\mu^+$ [86], and $\Sigma^- \to pe^-e^-$ [87] decays, respectively. To the best of our knowledge, LNV processes in four-body baryonic decays have not yet been searched. The LHCb has observed the $\Lambda_b \to \Lambda\mu^+\mu^-$ decay where the Λ is reconstructed in the $\Lambda \to p\pi^-$ [88,89] processes. The final state of this decay is similar to one of the mode considered in this paper. Therefore, it is possible in the near future for the LHCb to search for four-body LNV decays.

The paper is organized as follows. In Sec. II, we work out the formalism for a generic $\mathcal{B}_1 \to \mathcal{B}_2^{\mp}\pi^{\mp}\ell_1^{\pm}\ell_2^{\pm}$ decay mediated by on-shell Majorana neutrino. In Sec. III, we perform a numerical analysis of the CP asymmetry for $\Lambda_b \to (\Lambda_c, p)\pi\mu\mu$ and discuss the constraint on the $(m_N, |V_{\mu N}|^2)$ parameter space assuming experimental upper limits. We summarize our results in Sec. IV. Some details of our derivations are given in the Appendixes.

II. $\mathcal{B}_1 \to \mathcal{B}_2^{\mp}\pi^{\mp}\ell_1^{\pm}\ell_2^{\pm}$ FORMALISM

We consider a model scenario where, in addition to the components $\nu_{\ell L}$ of the left-handed $SU(2)_L$ doublets of the Standard Model, there are two right-handed singlet sterile neutrinos denoted by N_1 and N_2. The flavor eigenstates $\nu_{\ell L}$ can be written in terms of the mass eigenstates as

$$\nu_{\ell L} = \sum_{i=1}^{3} U_{\ell i}\nu_{iL} + V_{\ell N_1}N_1 + V_{\ell N_2}N_2, \qquad (2.1)$$

where ν_{iL} are the light mass eigenstates. We assume that the heavy to light mixing elements $V_{\ell N_1}$ and $V_{\ell N_2}$ are free parameters and can be constrained by experiments. They, in general, can be complex:

$$V_{\ell N_j} = |V_{\ell N_j}|e^{i\phi_{\ell j}} \quad (j = 1, 2), \qquad (2.2)$$

where $\phi_{\ell j}$ is a CP-odd phase. According to our convention, $V_{\ell N}$ is the mixing element between negatively charged lepton ℓ and Majorana neutrino N.

We are interested to calculate the decay widths of $\mathcal{B}_1(p_{\mathcal{B}_1}) \to \mathcal{B}_2(p_{\mathcal{B}_2})\pi^+(p_\pi)\ell_1^-(p_1)\ell_2^-(p_2)$ and its CP conjugate mode $\bar{\mathcal{B}}_1(p_{\mathcal{B}_1}) \to \bar{\mathcal{B}}_2(p_{\mathcal{B}_2})\pi^-(p_\pi)\ell_1^+(p_1)\ell_2^+(p_2)$ in this model. The decays can be viewed as a two-step processes: First, the $\mathcal{B}_1$ decays via a charged current interaction $\mathcal{B}_1 \to \mathcal{B}_2^{\mp}N_j\ell_1^{\pm}$, followed by the decay of the heavy neutrino $N_j \to \ell_2^{\pm}\pi^{\mp}$. For these processes there are two dominant “s-channel” topologies, the direct channel (D) and the crossed channel (C), as shown in Fig. 1. A “t-channel” Feynman diagram is also shown in Fig. 2. Here, the neutrino is off shell and the amplitude is suppressed, so we neglect the t-channel diagrams in our calculations.

Appreciable decay rates can be obtained if the neutrinos have kinematically allowed mass

$$\begin{aligned} &m_\pi + \ell_1 < m_{N_j} < (m_{\mathcal{B}_1} - m_{\mathcal{B}_2} - \ell_2) \quad \text{or/and} \\ &m_\pi + \ell_2 < m_{N_j} < (m_{\mathcal{B}_1} - m_{\mathcal{B}_2} - \ell_1). \end{aligned} \qquad (2.3)$$

We denote the momentum of the heavy neutrino in the D channel by $p_N = p_{\mathcal{B}_1} - p_{\mathcal{B}_2} - p_1$, and for the C channel by $p'_N = p_{\mathcal{B}_1} - p_{\mathcal{B}_2} - p_2$. Defining $\Gamma_{\mathcal{B}_1} \equiv \Gamma(\mathcal{B}_1 \to \mathcal{B}_2\pi^+\ell_1^-\ell_2^-)$ and $\Gamma_{\bar{\mathcal{B}}_1} \equiv \Gamma(\bar{\mathcal{B}}_1 \to \bar{\mathcal{B}}_2\pi^-\ell_1^+\ell_2^+)$ the decay widths can be written as

$$\Gamma_{\mathcal{B}_1(\bar{\mathcal{B}}_1)} = (2 - \delta_{\ell_1\ell_2})\frac{1}{2!}\frac{1}{2m_{\mathcal{B}_1}}\int d_4^{\rm PS}|\overline{\mathcal{M}}_{\rm tot}^{+(-)}|^2. \qquad (2.4)$$

The symmetry factor $1/2!$ comes because the two charged leptons can be the same. The $|\overline{\mathcal{M}}_{\rm tot}^{+}|^2$ ($|\overline{\mathcal{M}}_{\rm tot}^{-}|^2$) is the total

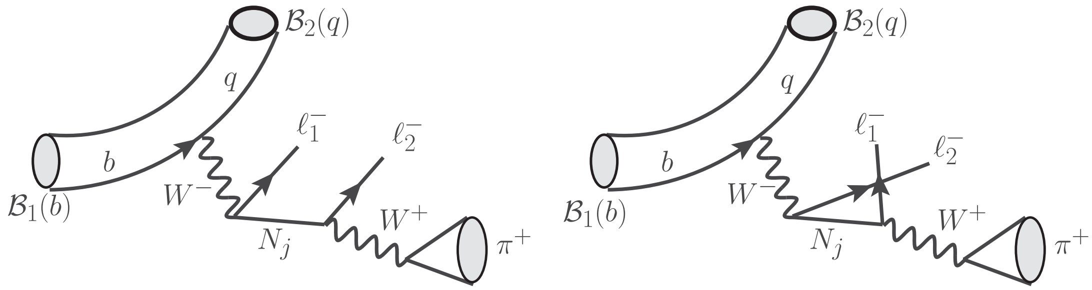

FIG. 1. The direct (D) and cross (C) channel Feynman diagrams for $\mathcal{B}_1 \to \mathcal{B}_2^{\mp}\pi^{\mp}\ell_1^{\pm}\ell_2^{\pm}$ decay.

matrix element mod squared of $\mathcal{B}_1 \to \mathcal{B}_2\pi^+\ell_1^-\ell_2^-$ ($\bar{\mathcal{B}}_1 \to \bar{\mathcal{B}}_2\pi^-\ell_1^+\ell_2^+$) after averaging over the initial spin and summing over the final spins

$$
\begin{aligned}
|\bar{\mathcal{M}}_{\rm tot}^{\pm}|^2 &= \frac{1}{2}\sum_{\rm spins}|\mathcal{M}_{\rm tot}^{\pm}|^2 \\
&= \frac{1}{2}\sum_{\rm spins}\left|\sum_{j=1}^{2}(\mathcal{M}_{D_j}^{\pm}+\mathcal{M}_{C_j}^{\pm})\right|^2 \\
&= \frac{1}{2}\sum_{\rm spins}\left[\sum_{i,j=1}^{2}\mathcal{M}_{D_i}^{\pm}(\mathcal{M}_{D_j}^{\pm})^* + \sum_{i,j=1}^{2}\mathcal{M}_{C_i}^{\pm}(\mathcal{M}_{C_j}^{\pm})^* + \sum_{i,j=1}^{2}\mathcal{M}_{D_i}^{\pm}(\mathcal{M}_{C_j}^{\pm})^* + \sum_{i,j=1}^{2}\mathcal{M}_{C_i}^{\pm}(\mathcal{M}_{D_j}^{\pm})^*\right] \\
&= \mathcal{N}\left[\sum_{i,j=1}^{2} v_i^{\pm}(v_j^{\pm})^* m_{N_i}m_{N_j}P_{D_i}P_{D_j}^{*}T_{\pm}(DD^*) + \sum_{i,j=1}^{2} v_i^{\pm}(v_j^{\pm})^* m_{N_i}m_{N_j}P_{D_i}P_{C_j}^{*}T_{\pm}(DC^*) \right. \\
&\qquad \left. + (D \leftrightarrow C)\right].
\end{aligned}
\tag{2.5}
$$

In the second line of Eq. (2.5), the suffix D_j (C_j) stand for the direct (cross) channel with jth neutrino exchange, and in the last line we have introduced the following notations:

$$
\mathcal{N} = \frac{1}{2}G_F^4|V_{ud}|^2|V_{qb}|^2 f_\pi^2, \qquad v_i^+ = V_{\ell_1 N_i}V_{\ell_2 N_i},
$$
$$
v_i^- = (v_i^+)^*, \tag{2.6}
$$

where $V_{qb} = V_{ub}$ for $\Lambda_b \to p^+\pi^+\ell^-\ell^-$, $V_{qb} = V_{cb}$ for $\Lambda_b \to \Lambda_c^+\pi^+\ell^-\ell^-$, and f_π is the pion decay constant. In the last line of Eq. (2.5), the spin summed and averaged matrix element mod squared splits into universal functions $T_{\pm}(XY^*)$, where $X(Y) = D$, C, and the functions P_{X_j}, which are functions of the masses m_{N_1} and m_{N_2} and decay widths Γ_{N_1} and Γ_{N_2} of the exchanged neutrinos:

$$
\begin{aligned}
P_{D_j} &= \frac{1}{(p_N^2 - m_{N_j}^2) + i\Gamma_{N_j}m_{N_j}}, \\
P_{C_j} &= \frac{1}{(p_N'^2 - m_{N_j}^2) + i\Gamma_{N_j}m_{N_j}}.
\end{aligned}
\tag{2.7}
$$

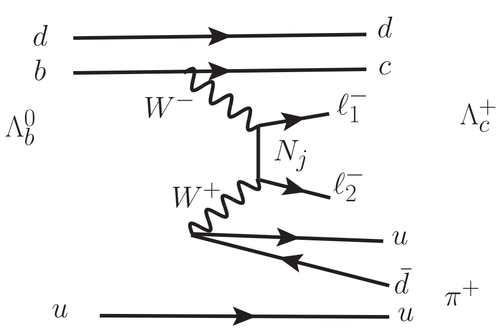

FIG. 2. A t-channel diagram for $\Lambda_b \to \Lambda_c^+\pi^+\ell_1^-\ell_2^-$ decay.

Using Eq. (2.5), the total decay widths can be conveniently written as

$$
\begin{aligned}
\Gamma_{\mathcal{B}_1} = (2-\delta_{\ell_1\ell_2})\sum_{i,j=1}^{2} v_i^+(v_j^+)^*(&\hat{\Gamma}(DD^*)_{ij} + \hat{\Gamma}(CC^*)_{ij} \\
&+ \hat{\Gamma}_+(DC^*)_{ij} + \hat{\Gamma}_+(D^*C)_{ij}),
\end{aligned}
\tag{2.8}
$$

$$
\begin{aligned}
\Gamma_{\bar{\mathcal{B}}_1} = (2-\delta_{\ell_1\ell_2})\sum_{i,j=1}^{2} v_i^-(v_j^-)^*(&\hat{\Gamma}(DD^*)_{ij} + \hat{\Gamma}(CC^*)_{ij} \\
&+ \hat{\Gamma}_-(DC^*)_{ij} + \hat{\Gamma}_-(D^*C)_{ij}),
\end{aligned}
\tag{2.9}
$$

where the quantities $\hat{\Gamma}_{\pm}$ are

$$\hat{\Gamma}_{\pm}(XY^*)_{ij} = \frac{\mathcal{N}}{2m_{\mathcal{B}_1}2!}\int m_{N_i}m_{N_j}P_{X_i}P^*_{Y_j}T_{\pm}(XY^*)d\Phi_4,$$
$$X, Y = C, D. \tag{2.10}$$

The expressions of $T_{\pm}(XY^*)$ and the requisite kinematics to evaluate these expressions are given in Appendixes A and B, respectively, and the four-body phase space $d\Phi_4$ is given in Appendix C. In Eqs. (2.8) and (2.9), using the relation $T_+(XX^*) = T_-(XX^*)$ (see Appendix A), we have defined

$$\hat{\Gamma}(XX^*)_{ij} \equiv \hat{\Gamma}_+(XX^*)_{ij} = \hat{\Gamma}_-(XX^*)_{ij}, \quad X = D, C. \tag{2.11}$$

To physically interpret the terms, $\hat{\Gamma}(XX^*)_{ij}$ are the contributions of N_i exchange in the X channel and the conjugate of N_j exchange in the X^* channel. The interference terms $\hat{\Gamma}(XY^*)_{ij}$ are the contributions of N_i exchange in the X channel and the conjugate of N_j exchange the Y channel. Numerically, $D-C$ channel interference contributions $\hat{\Gamma}(XY^*)_{ij}$ for $X \neq Y$ are insignificant compared to $\hat{\Gamma}(XX^*)_{ij}$ and are ignored in our calculations.

In addition to the decay rates, the quantities of interest are their sum and differences:

$$\Gamma_{\mathcal{B}_1} + \Gamma_{\bar{\mathcal{B}}_1} = 2(2-\delta_{\ell_1\ell_2})[|V_{\ell_1N_1}|^2|V_{\ell_2N_1}|^2(\hat{\Gamma}(DD^*)_{11} + \hat{\Gamma}(CC^*)_{11}) + |V_{\ell_1N_2}|^2|V_{\ell_2N_2}|^2(\hat{\Gamma}(DD^*)_{22} + \hat{\Gamma}(CC^*)_{22})$$
$$+ 2\cos(\theta_{21})|V_{\ell_1N_1}||V_{\ell_2N_1}||V_{\ell_1N_2}||V_{\ell_2N_2}|(\mathrm{Re}\hat{\Gamma}(DD^*)_{12} + \mathrm{Re}\hat{\Gamma}(CC^*)_{12})], \tag{2.12}$$

$$\Gamma_{\mathcal{B}_1} - \Gamma_{\bar{\mathcal{B}}_1} = 4(2-\delta_{\ell_1\ell_2})|V_{\ell_1N_1}||V_{\ell_2N_1}||V_{\ell_1N_2}||V_{\ell_2N_2}|[\sin(\theta_{21})(\mathrm{Im}\hat{\Gamma}(DD^*)_{12} + \mathrm{Im}\hat{\Gamma}(CC^*)_{12})], \tag{2.13}$$

where the CP-odd phase, based on the convention adopted in Eq. (2.2), is

$$\theta_{ij} = \arg(V_{\ell_1N_i}) + \arg(V_{\ell_2N_i}) - \arg(V_{\ell_1N_j}) - \arg(V_{\ell_2N_j})$$
$$= (\phi_{1i} + \phi_{2i} - \phi_{1j} - \phi_{2j}), \quad i, j = 1, 2. \tag{2.14}$$

A CP-even phase $\Delta\xi = \xi_1 - \xi_2$ essential for CP violation is also present in the interference of N_1 and N_2 contributions:

$$\mathrm{Re}\hat{\Gamma}(XX^*)_{12} = \frac{\mathcal{N}}{2m_{\mathcal{B}_1}2!}\int m_{N_1}m_{N_2}|P_{X_1}||P_{X_2}|\cos(\Delta\xi)T(XX^*)d_4^{\mathrm{PS}}, \qquad X = C, D, \tag{2.15}$$

$$\mathrm{Im}\hat{\Gamma}(XX^*)_{12} = \frac{\mathcal{N}}{2m_{\mathcal{B}_1}2!}\int m_{N_1}m_{N_2}|P_{X_1}||P_{X_2}|\sin(\Delta\xi)T(XX^*)d_4^{\mathrm{PS}}, \qquad X = D, C, \tag{2.16}$$

where $\xi_{1,2}$ are given as

$$\tan\xi_1 = \frac{m_{N_1}\Gamma_{N_1}}{k_N^2 - m_{N_1}^2}, \qquad \tan\xi_2 = \frac{m_{N_2}\Gamma_{N_2}}{k_N^2 - m_{N_2}^2}, \tag{2.17}$$

and $k_N^2 = p_N^2$ for D channel and $k_N^2 = (p'_N)^2$ for C channel.

III. RESULTS

Following the formalism in the previous section, we turn to numerical analysis with specific decay modes. At the LHC, about 5% of the total b hadrons produced are Λ_b baryons, and both at the LHCb and CMS the muon reconstruction efficiency is comparatively higher than the other two charged leptons. We therefore are interested in the modes $\Lambda_b \to \Lambda_c\pi\mu\mu$ and $\Lambda_b \to p\pi\mu\mu$ channels. Since $\ell_1 = \ell_2 = \mu$, the CP-odd phase is $\theta_{21} = 2(\phi_{\mu2} - \phi_{\mu1})$. For numerical analysis, form factors parameterizing the $\Lambda_b^0 \to \Lambda_c^+$ and $\Lambda_b^0 \to p^+$ hadronic matrix elements are taken from the lattice QCD calculations [90], and we take the decay constant of pion $f_\pi = 130.2(0.8)$ MeV from Ref. [91].

We also need to know the total decay widths of the heavy neutrinos $\Gamma_{N_{1,2}}$ as a function of their masses. For a Majorana neutrino mass between $m_\pi + m_\mu < m_N < (m_{\mathcal{B}_1} - m_{\mathcal{B}_2} - m_\mu)$, both purely leptonic as well as semihadronic decays may be relevant. For $m_N < 1$ GeV, the decays to leptonic modes as well as to light pseudoscalar and vector mesons have been calculated in Ref. [92]. For higher values of m_N, decays to semihadronic modes are increasingly difficult due to the limited knowledge of the resonances. An inclusive approach based on quark-hadron duality was adopted in Refs. [35,93] to calculate the widths of the semihadronic channel. For this analysis, we leave the decay width as a phenomenological parameter that can be measured by experiments. Following the analysis of Ref. [48], we take the neutrino the lifetimes $\tau_N = \hbar/\Gamma_N = [10, 100, 1000]$ ps for numerical illustration.

We are interested in the signal of leptonic CP asymmetry

$$\mathcal{A}_{CP} = \frac{\Gamma_{\mathcal{B}_1} - \Gamma_{\bar{\mathcal{B}}_1}}{\Gamma_{\mathcal{B}_1} + \Gamma_{\bar{\mathcal{B}}_1}}. \tag{3.1}$$

The reason this asymmetry will be present in the decay can be understood as follows. There are two interfering amplitudes coming from the two intermediate neutrinos N_1 and N_2. The interfering amplitudes have CP-odd phase θ_{21} that changes sign for the conjugate process. A CP-even phase $\Delta\xi$ comes from an absorptive part that is generated due to the interference of the two neutrino contributions and does not change sign in the conjugate process. In general, θ_{21} can be anything, but a maximal $\mathcal{A}_{CP}$ can be obtained for $\theta_{21} = \pi/2$ as can be seen from Eq. (2.13). To understand the behavior of $\mathcal{A}_{CP}$ with the neutrino mass, we note that $\mathrm{Im}[\hat{\Gamma}(DD^*)_{ij}] \propto \mathrm{Im}[P_{D_i} P^*_{D_j}]$. For our choices of the neutrino lifetime τ_N and the kinematically allowed neutrino mass m_N, the approximation $\Gamma_{N_j} \ll m_{N_j}$ is always valid so that

$$|P_{D_j}|^2 = \frac{\pi}{m_{N_j}\Gamma_{N_j}} \delta(p_N^2 - m_{N_j}^2),$$
$$|P_{C_j}|^2 = \frac{\pi}{m_{N_j}\Gamma_{N_j}} \delta(p_N'^2 - m_{N_j}^2), \tag{3.2}$$

which yields

$$\frac{\hat{\Gamma}(XX^*)_{ii}}{\hat{\Gamma}(XX^*)_{jj}} = \frac{\Gamma_{N_j}}{\Gamma_{N_i}}. \tag{3.3}$$

When the mass difference between the neutrinos is such that $\Gamma_{N_j} \ll \Delta m_N$, then we can write

$$\begin{aligned}
\mathrm{Im}[P_{D_1} P^*_{D_2}]|_{\Gamma_{N_j} \ll \Delta m_N} &= \mathcal{P}\left(\frac{1}{p_N^2 - m_{N_1}^2}\right)\pi\delta(p_N^2 - m_{N_2}^2) - \mathcal{P}\left(\frac{1}{p_N^2 - m_{N_2}^2}\right)\pi\delta(p_N^2 - m_{N_1}^2) \\
&= \frac{\pi}{m_{N_2}^2 - m_{N_1}^2}(\delta(p_N^2 - m_{N_1}^2) + \delta(p_N^2 - m_{N_2}^2)) \\
&= \frac{1}{y}\frac{2\pi}{(m_{N_1} + m_{N_2})(\Gamma_{N_1} + \Gamma_{N_1})}(\delta(p_N^2 - m_{N_1}^2) + \delta(p_N^2 - m_{N_2}^2)),
\end{aligned} \tag{3.4}$$

where $y = \Delta m_N/\Gamma_N$ and $\Gamma_N = (\Gamma_{N_1} + \Gamma_{N_2})/2$. This yields

$$\begin{aligned}
\frac{\mathrm{Im}\hat{\Gamma}(XX^*)_{12}}{\hat{\Gamma}(XX^*)_{jj}} &= \frac{\mathrm{Im}\hat{\Gamma}(XX^*)_{12}|_{\Gamma_{N_j} \ll \Delta m_N}}{\hat{\Gamma}(XX^*)_{jj}} \frac{\mathrm{Im}\hat{\Gamma}(XX^*)_{12}}{\mathrm{Im}\hat{\Gamma}(XX^*)_{12}|_{\Gamma_{N_j} \ll \Delta m_N}} \\
&= \frac{1}{y}\frac{4\pi}{(m_{N_1} + m_{N_2})(\Gamma_{N_1} + \Gamma_{N_1})}\frac{m_{N_j}\Gamma_{N_j}}{\pi}\eta,
\end{aligned} \tag{3.5}$$

where the suppression factor

$$\eta \equiv \frac{\mathrm{Im}\hat{\Gamma}(XX^*)_{12}}{\mathrm{Im}\hat{\Gamma}(XX^*)_{12}|_{\Gamma_{N_j} \ll \Delta m_N}} \tag{3.6}$$

accounts for the departure from the approximation $\Gamma_{N_j} \ll \Delta m_N$ in the term $\mathrm{Im}\hat{\Gamma}(XX^*)$. Assuming the neutrinos to be almost degenerate $m_{N_1} \sim m_{N_2} = m_N$,

$$\frac{\mathrm{Im}\hat{\Gamma}(XX^*)_{12}}{\hat{\Gamma}(XX^*)_{jj}} = \frac{2\Gamma_{N_j}}{(\Gamma_{N_1} + \Gamma_{N_1})}\frac{\eta(y)}{y}. \tag{3.7}$$

We define another factor

$$\delta_j(y) = \frac{\mathrm{Re}\hat{\Gamma}(XX^*)_{12}}{\hat{\Gamma}(XX^*)_{jj}}, \tag{3.8}$$

which measures the interference of the two neutrinos in the real part of $\hat{\Gamma}(XX^*)$. Following Eq. (3.3), we get

$$\frac{\delta_1}{\delta_2} = \frac{\Gamma_{N_1}}{\Gamma_{N_2}}. \tag{3.9}$$

Since we are considering a same-sign dimuon in the final state, η and δ_j are the same for the D and C channels, which follows from the fact that

$$\hat{\Gamma}(DD^*)_{jj} = \hat{\Gamma}(CC^*)_{jj}, \quad \mathrm{Re}\hat{\Gamma}(DD^*)_{12} = \mathrm{Re}\hat{\Gamma}(CC^*)_{12},$$
$$\mathrm{Im}\hat{\Gamma}(DD^*)_{12} = \mathrm{Im}\hat{\Gamma}(CC^*)_{12}. \tag{3.10}$$

The CP asymmetry can now be written in a convenient form as

$$\mathcal{A}_{CP} = \frac{4\sin\theta_{21}}{\frac{|V_{\ell N_1}||V_{\ell N_1}|}{|V_{\ell N_2}||V_{\ell N_2}|}(1 + \frac{\Gamma_{N_2}}{\Gamma_{N_1}}) + \frac{|V_{\ell N_2}||V_{\ell N_2}|}{|V_{\ell N_1}||V_{\ell N_1}|}(1 + \frac{\Gamma_{N_1}}{\Gamma_{N_2}}) + 4\delta(y)\cos\theta_{21}}\frac{\eta(y)}{y}, \tag{3.11}$$

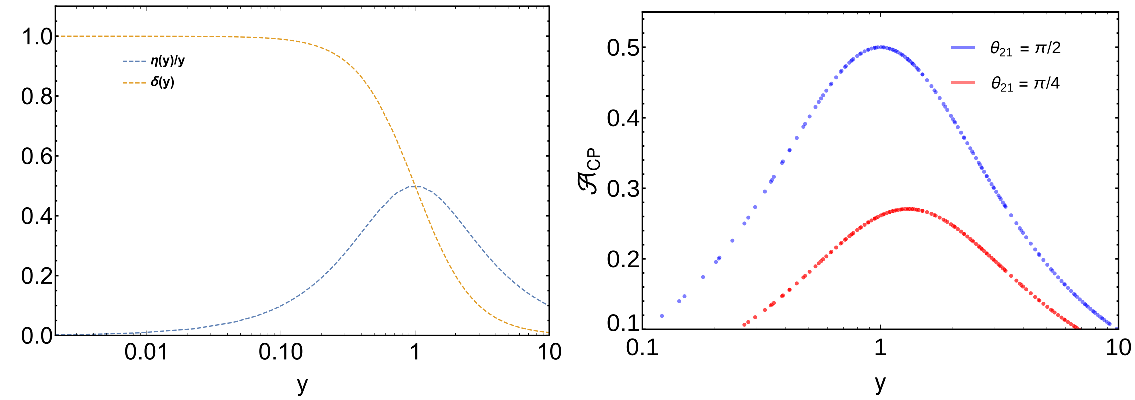

FIG. 3. The factors $\delta(y)$ and $\eta(y)/y$ as a function of $y = \Delta m_N/\Gamma_N$. The CP asymmetry observable $\mathcal{A}_{CP}$ for $\Lambda_b \to \Lambda_c \pi\mu\mu$ is shown as a function of y for different values of the weak phase. An identical plot is obtained for $\Lambda_b \to p\pi\mu\mu$. In these plots, we have taken $|V_{\mu N_1}|^2 = |V_{\mu N_2}|^2 = 1$.

where we define

$$\delta(y) = \frac{\delta_1(y) + \delta_2(y)}{2}. \tag{3.12}$$

For nearly degenerate neutrinos, it is natural to assume $|V_{\mu N_1}| \sim |V_{\mu N_2}| = |V_{\mu N}|$. This further simplifies the expression of $\mathcal{A}_{CP}$:

$$\mathcal{A}_{CP} = \frac{4\sin\theta_{21}}{(1+\frac{\Gamma_{N_2}}{\Gamma_{N_1}}) + (1+\frac{\Gamma_{N_1}}{\Gamma_{N_2}}) + 4\delta(y)\cos\theta_{21}} \frac{\eta(y)}{y}. \tag{3.13}$$

In Fig. 3, we show the suppression factor $\eta(y)/y$ and $\delta(y)$ as a function of y. This figure demonstrates that the $\mathcal{A}_{CP}$ will be maximum for $y \sim 1$, i.e., when $\Delta m_N \sim \Gamma_N$. In Fig. 3, we also show the $\mathcal{A}_{CP}$ for the $\Lambda_b \to \Lambda\pi\mu\mu$ mode for different values of θ_{21} and as a function of y. An identical plot is obtained for $\Lambda_b \to p\pi\mu\mu$. For a particular mode, the possibility to observe $\mathcal{A}_{CP}$ does not depend entirely on its size but also depends on the decay rates. In Fig. 4, we show the CP-averaged branching ratios

$$\mathcal{B}r(\mathcal{B}_1 \to \mathcal{B}_2\pi\mu\mu) = \frac{1}{2}(\mathcal{B}r(\mathcal{B}_1 \to \mathcal{B}_2\pi^+\mu^-\mu^-) + \mathcal{B}r(\bar{\mathcal{B}}_1 \to \bar{\mathcal{B}}_2\pi^-\mu^+\mu^+)) \tag{3.14}$$

of $\Lambda_b \to \Lambda_c\pi\mu\mu$ and $\Lambda_b \to p\pi\mu\mu$ as a function of sterile neutrino mass m_{N_1} for neutrino lifetimes $\tau_{N_1} \sim \tau_{N_2} \sim 100$ and 1000 ps, $|V_{\mu N}|^2 \sim 10^{-5}$, $\theta_{21} = \pi/4$, and the neutrino mass difference $\Delta m_N = 10^{-15}$ GeV. We find that the branching ratio of $\Lambda_b \to \Lambda_c\pi\mu\mu$ can be within the $10^{-10} - 10^{-9}$ range, whereas $\mathcal{B}(\Lambda_b \to p\pi\mu\mu) \sim 10^{-12} - 10^{-11}$ is suppressed due to small Cabibbo-Kobayashi-Maskawa (CKM) element V_{ub}. The LHCb has already observed the lepton-number-conserving mode $\Lambda_b \to \Lambda(\to p\pi)\mu^+\mu^-$ [88,89]. In the next LHCb upgrade, about $\sim 10^{12}$ number of Λ_b is expected to be produced [94,95]. Hence, the LNV rates could be within the reach of the future LHC sensitivity. For a detailed discussion of the number of expected events at the LHCb and CMS, please see Refs. [48,49].

Even if the decays are not fully observed, upper limits can be translated to bounds on the m_N vs $|V_{\mu N}|^2$ parameter

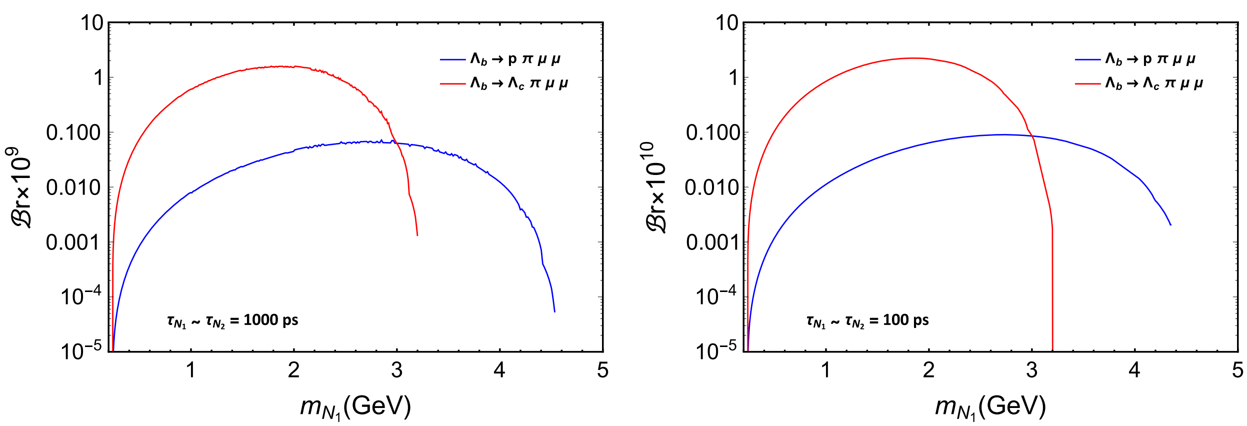

FIG. 4. The branching ratios $\mathcal{B}(\Lambda_b \to \Lambda_c\pi\mu\mu)$ and $\mathcal{B}(\Lambda_b \to p\pi\mu\mu)$ for $|V_{\mu N_1}|^2 \sim |V_{\mu N_2}|^2 = |V_{\mu N}|^2 = 10^{-5}$, the weak phase $\theta_{21} = \pi/4$, the mass difference $\Delta m_N = 10^{-15}$ GeV, and the neutrino lifetimes $\tau_N = [100, 1000]$ ps.

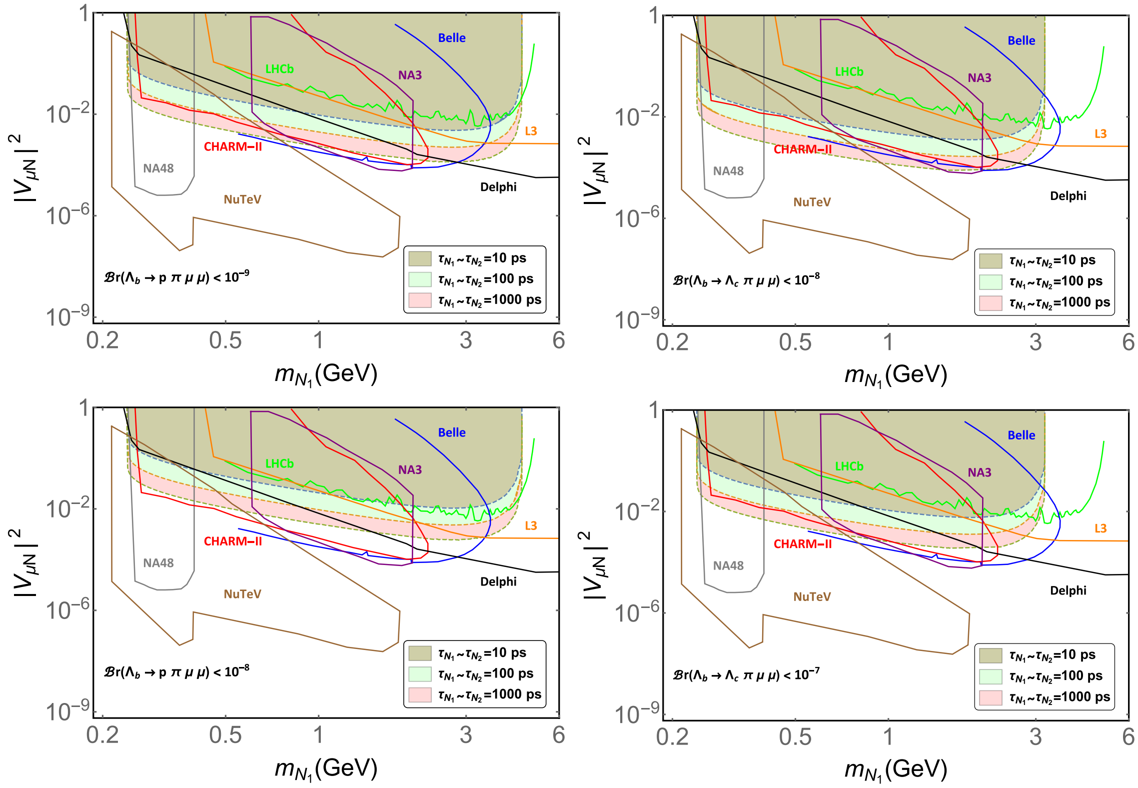

FIG. 5. Exclusion regions on the $(m_{N_1}, |V_{\mu N}|^2)$ parameter space for $\mathcal{B}r(\Lambda_b \to p\pi\mu\mu) < 10^{-8}, 10^{-9}$ and $\mathcal{B}r(\Lambda_b \to \Lambda_c\pi\mu\mu) < 10^{-7}, 10^{-8}$ for different values of τ_N, $\theta_{21} = \pi/4$, and $\Delta m_N = 10^{-15}$ GeV.

space. In Fig. 5, we show the exclusion region in the $(m_N, |V_{\mu N}|^2)$ plane obtained by assuming upper bounds $\mathcal{B}r(\Lambda_b \to p\pi\mu\mu) < 10^{-8}$, 10^{-9} and $\mathcal{B}r(\Lambda_b \to \Lambda_c\pi\mu\mu) < 10^{-7}$, 10^{-8} for different choices of the heavy neutrino lifetimes. The regions shown in brown, light green, and light red correspond to exclusion regions obtained for $\tau_N = 10$, 100, and 1000 ps, respectively. To compare our bounds, in Fig. 5, we also show the exclusion limits from LHCb [43,82], Belle [96], L3 [97], Delphi [98], NA3 [99], CHARM [100], NuTeV [101], and NA48 [83] experiments. These comparisons show that the LNV modes $\Lambda_b \to p\pi\mu\mu$ and $\Lambda_b \to p\pi\mu\mu$ can give complementary bounds on the sterile neutrino parameters. And, with the possibility to observe *CP* asymmetry, the modes should be searched at the LHC.

IV. SUMMARY

In this paper, we have studied lepton-number-violating baryonic decays $\Lambda_b \to \Lambda_c\pi\mu\mu$ and $\Lambda_b \to p\pi\mu\mu$ that are mediated by on-shell sterile Majorana neutrinos. The decays are studied in a model where there are two Majorana neutrinos. An interesting consequence of considering two Majorana neutrinos is that it gives rise to the possibility of *CP* violation in these modes. We find that appreciable *CP* asymmetry can be achieved if neutrinos are quasidegenerate and the mass difference is of the order of decay widths. We have shown that, in the absence of observation, upper limits to the branching ratios can give limits on the m_N vs $|V_{\mu N}|^2$ parameter space that is comparable to limits obtained by other methods.

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APPENDIX A: $\mathcal{B}_1 \to \mathcal{B}_2^{\mp}\pi^{\mp}\ell_1^{\pm}\ell_2^{\pm}$ AMPLITUDES

The effective Hamiltonian for the $\mathcal{B}_1 \to \mathcal{B}_2\ell N$ decay and the subsequent decay of the intermediate neutrino $N \to \ell\pi$ are

$$\mathcal{H}_{\mathrm{eff}}^{b\to q\ell N}=\frac{G_F}{\sqrt{2}}V_{qb}\bar{q}\gamma_\mu(1-\gamma_5)b\left(\sum_{\ell=e}^{\tau}\sum_{i=1}^{3}U_{\ell i}\bar{\nu}_i\gamma^\mu(1-\gamma_5)\ell+\sum_{\ell=e}^{\tau}\sum_{j=1}^{n}V_{\ell N_j}\bar{N}_j^c\gamma^\mu(1-\gamma_5)\ell\right)+\mathrm{H.c.},\tag{A1}$$

$$\mathcal{H}_{\mathrm{eff}}^{N_j\to\ell\pi}=\frac{G_F}{\sqrt{2}}V_{ud}\bar{d}\gamma_\mu(1-\gamma_5)u\sum_{\ell=e}^{\tau}\sum_{j=1}^{n}V_{\ell N_j}\bar{\ell}\gamma^\mu(1-\gamma_5)N_j^c+\mathrm{H.c.}\tag{A2}$$

The $\mathcal{B}_1\to\mathcal{B}_2\pi\ell\ell$ amplitudes for the D- and C-channel diagrams are

$$\mathcal{M}_{D_j}^{\pm}=(G_F^2V_{qb}M_{N_j})(V_{\ell_1N_j}V_{\ell_2N_j})P_{D_j}H_\nu^{\pm}L_D^{\nu\alpha\pm}J_\alpha^{\pm},\tag{A3}$$

$$\mathcal{M}_{C_j}^{\pm}=(G_F^2V_{qb}M_{N_j})(V_{\ell_1N_j}V_{\ell_2N_j})P_{C_j}H_\nu^{\pm}L_C^{\alpha\nu\pm}J_\alpha^{\pm},\tag{A4}$$

where the $+$ corresponds to the modes $\Lambda_b^0\to(\Lambda_c^+,p^+)\pi^+\ell_1^-\ell_2^-$ and the $-$ corresponds to the conjugate modes. According to our convention, $H_\nu^+=H_\nu$ and $H_\nu^-=H_\nu^*$. The CKM elements $V_{qb}=V_{ub}$ for $\Lambda_b^0\to p^+\pi\ell\ell$ and $V_{qb}=V_{cb}$ for $\Lambda_b^0\to\Lambda_c^+\pi\ell\ell$ modes. The leptonic parts of the amplitudes are

$$L_D^{\nu\alpha\pm}=\bar{u}_{\ell_1}(p_1)\gamma^\nu\gamma^\alpha(1\pm\gamma_5)v_{\ell_2}(p_2),\qquad L_C^{\alpha\nu\pm}=\bar{u}_{\ell_1}(p_1)\gamma^\alpha\gamma^\nu(1\pm\gamma_5)v_{\ell_2}(p_2).\tag{A5}$$

The hadronic amplitudes H^μ are calculated using the form factor parametrization of $\mathcal{B}_1\to\mathcal{B}_2$ transitions from Ref. [90]:

$$\begin{aligned}\langle\mathcal{B}_2(k,s_k)|\bar{s}\gamma^\mu b|\mathcal{B}_1(p,s_p)\rangle=\bar{u}(k,s_k)\Bigg[&f_t^V(q^2)(m_{\mathcal{B}_1}-m_{\mathcal{B}_2})\frac{q^\mu}{q^2}+f_0^V(q^2)\frac{m_{\mathcal{B}_1}+m_{\mathcal{B}_2}}{s_+}\left\{p^\mu+k^\mu-\frac{q^\mu}{q^2}(m_{\mathcal{B}_1}^2-m_{\mathcal{B}_2}^2)\right\}\\&+f_\perp^V(q^2)\left\{\gamma^\mu-\frac{2m_{\mathcal{B}_2}}{s_+}p^\mu-\frac{2m_{\mathcal{B}_1}}{s_+}k^\mu\right\}\Bigg]u(p,s_p),\end{aligned}\tag{A6}$$

$$\begin{aligned}\langle\mathcal{B}_2(k,s_k)|\bar{s}\gamma^\mu\gamma_5 b|\mathcal{B}_1(p,s_p)\rangle=-\bar{u}(k,s_k)\gamma_5\Bigg[&f_t^A(q^2)(m_{\mathcal{B}_1}+m_{\mathcal{B}_2})\frac{q^\mu}{q^2}+f_0^A(q^2)\frac{m_{\mathcal{B}_1}-m_{\mathcal{B}_2}}{s_-}\left\{p^\mu+k^\mu-\frac{q^\mu}{q^2}(m_{\mathcal{B}_1}^2-m_{\mathcal{B}_2}^2)\right\}\\&+f_\perp^A(q^2)\left\{\gamma^\mu+\frac{2m_{\mathcal{B}_2}}{s_-}p^\mu-\frac{2m_{\mathcal{B}_1}}{s_-}k^\mu\right\}\Bigg]u(p,s_p).\end{aligned}\tag{A7}$$

Using Eqs. (A6) and (A7), we can write the expression of H^μ as

$$\begin{aligned}H^\mu&=\langle\mathcal{B}_2(k,s_k)|\bar{c}\gamma^\mu(1-\gamma_5)b|\mathcal{B}_1(p,s_p)\rangle\\&=\bar{u}(k,s_k)(A_1q^\mu+A_2k^\mu+A_3\gamma^\mu\\&\quad+\gamma_5\{A_4q^\mu+A_5k^\mu+A_6\gamma^\mu\})u(p,s_p),\end{aligned}\tag{A8}$$

where the q^2-dependent functions A_i can be written in terms of the form factors as

$$\begin{aligned}A_1=f_t^V\frac{m_{\mathcal{B}_1}-m_{\mathcal{B}_2}}{q^2}+f_0^V\frac{m_{\mathcal{B}_1}+m_{\mathcal{B}_2}}{s_+}\left(1-\frac{m_{\mathcal{B}_2}^1-m_{\mathcal{B}_2}^2}{q^2}\right)\\-f_\perp^V\frac{2m_{\mathcal{B}_2}}{s_+},\end{aligned}\tag{A9}$$

$$A_2=2f_0^V\frac{m_{\mathcal{B}_1}+m_{\mathcal{B}_2}}{s_+}-f_\perp^V\left(\frac{2m_{\mathcal{B}_2}}{s_+}+\frac{2m_{\mathcal{B}_1}}{s_+}\right),\tag{A10}$$

$$A_3=f_\perp^V,\tag{A11}$$

$$\begin{aligned}A_4=f_t^A\frac{m_{\mathcal{B}_1}+m_{\mathcal{B}_2}}{q^2}+f_0^A\frac{m_{\mathcal{B}_1}-m_{\mathcal{B}_2}}{s_-}\left(1-\frac{m_{\mathcal{B}_2}^1-m_{\mathcal{B}_2}^2}{q^2}\right)\\+f_\perp^A\frac{2m_{\mathcal{B}_2}}{s_-},\end{aligned}\tag{A12}$$

$$A_5=2f_0^A\frac{m_{\mathcal{B}_1}-m_{\mathcal{B}_2}}{s_-}+f_\perp^A\left(\frac{2m_{\mathcal{B}_2}}{s_-}-\frac{2m_{\mathcal{B}_1}}{s_-}\right),\tag{A13}$$

$$A_6=f_\perp^A.\tag{A14}$$

Finally, the amplitudes for the pion production are

$$\begin{aligned}J_\alpha^+&=\langle\pi^+(k)|\bar{u}\gamma_\mu(1-\gamma_5)d|0\rangle=ik_\mu V_{ud}f_\pi,\\J_\alpha^-&=-ik_\mu V_{ud}^*f_\pi=(J_\alpha^+)^\dagger.\end{aligned}\tag{A15}$$

The matrix element mod squared after summing over the final spins and averaging over the initial spin is given in Eq. (2.5). The quadratic terms $T_\pm(XY^*)$ and $T(XY^*)$ given in Eq. (2.5) are

$$T_\pm(DD^*) = \sum_{\text{spins}}[H_\nu^\pm(H_\rho^\pm)^*]\sum_{\text{spins}}[L_D^{\nu\alpha\pm}(L_D^{\rho\beta\pm})^*]k_{\pi\alpha}k_{\pi\beta}, \tag{A16}$$

$$T_\pm(CC^*) = \sum_{\text{spins}}[H_\nu^\pm(H_\rho^\pm)^*]\sum_{\text{spins}}[L_C^{\alpha\nu\pm}(L_C^{\beta\rho\pm})^*]k_{\pi\alpha}k_{\pi\beta}, \tag{A17}$$

$$T_\pm(DC^*) = \sum_{\text{spins}}[H_\nu^\pm(H_\rho^\pm)^*]\sum_{\text{spins}}[L_D^{\nu\alpha\pm}(L_C^{\beta\rho\pm})^*]k_{\pi\alpha}k_{\pi\beta}, \tag{A18}$$

$$T_\pm(D^*C) = \sum_{\text{spins}}[(H_\nu^\pm)^*H_\rho^\pm]\sum_{\text{spins}}[(L_D^{\nu\alpha\pm})^*L_C^{\beta\rho\pm}]k_{\pi\alpha}k_{\pi\beta}, \tag{A19}$$

$$\begin{aligned} T(DD^*) &\equiv T_+(DD^*) = T_-(DD^*), \\ T(CC^*) &\equiv T_+(CC^*) = T_-(CC^*), \end{aligned} \tag{A20}$$

$$T_+(DC^*) = T_-(D^*C), \qquad T_-(DC^*) = T_+(D^*C). \tag{A21}$$

In Appendix B, we describe the kinematics required to calculate the momentum dot products required to evaluate the quadratic terms $T_\pm(XY^*)$.

APPENDIX B: KINEMATICS

As mentioned in the text, the contributions coming from the interference of the direct and cross channel diagrams are negligibly small and neglected in our calculations. Therefore, kinematics of the direct and cross channel can be evaluated independently. In this section, we work out the kinematics for the direct channel. The cross channel can be obtained trivially from the results presented here.

Referring to the diagrams in Fig. 1, in this section, we work out the kinematics of $\mathcal{B}_1^0(p_{\mathcal{B}_1}) \to \mathcal{B}_2(p_{\mathcal{B}_2})\pi^+(p_\pi)\ell_1^-(p_1)\ell_2^-(p_2)$ decay in the $\mathcal{B}_1(p_{\mathcal{B}_1})$ rest frame ($\mathcal{B}_1$-RF). In this frame, the four momentum of $\mathcal{B}_2(p_{\mathcal{B}_2})$ and $W_1(q)$ are

$$p_{\mathcal{B}_2}^{\mathcal{B}_1\text{-RF}} \equiv (m_{\mathcal{B}_1} - E_q^{\mathcal{B}_1\text{-RF}}, 0, 0, \mathbf{p}_{\mathcal{B}_2}^{\mathcal{B}_1\text{-RF}}), \tag{B1}$$

$$q^{\mathcal{B}_1\text{-RF}} \equiv (E_q^{\mathcal{B}_1\text{-RF}}, 0, 0, -\mathbf{p}_{\mathcal{B}_2}^{\mathcal{B}_1\text{-RF}}), \tag{B2}$$

where the q^0 and the $\mathbf{p}_{\mathcal{B}_2}^{\mathcal{B}_1\text{-RF}}$ are

$$\begin{aligned} E_q^{\mathcal{B}_1\text{-RF}} &= \frac{m_{\mathcal{B}_2}^1 + q^2 - m_{\mathcal{B}_2}^2}{2m_{\mathcal{B}_1}}, \\ |\mathbf{p}_{\mathcal{B}_2}^{\mathcal{B}_1\text{-RF}}| &= \frac{\sqrt{\lambda(m_{\mathcal{B}_2}^1, m_{\mathcal{B}_2}^2, q^2)}}{2m_{\mathcal{B}_1}}, \end{aligned} \tag{B3}$$

and $\lambda(a,b,c) = a^2 + b^2 + c^2 - 2(ab + bc + ca)$. In the W^--RF, we define θ_1 as the angle made by ℓ_1 with respect to the $\mathcal{B}_2$, i.e., in the $+\hat{z}$ direction. The four momentum of $\ell_1(p_1)$ and $N(p_N)$ read

$$p_1^{W^-\text{-RF}} = (E_1^{W^-\text{-RF}}, |\mathbf{p}_1^{W^-\text{-RF}}|\sin\theta_1, 0, |\mathbf{p}_1^{W^-\text{-RF}}|\cos\theta_1), \tag{B4}$$

$$\begin{aligned} p_N^{W^-\text{-RF}} = \Big(&\sqrt{q^2} - E_1^{W^-\text{-RF}}, -|\mathbf{p}_1^{W^-\text{-RF}}|\sin\theta_1, 0, \\ &- |\mathbf{p}_1^{W^-\text{-RF}}|\cos\theta_1\Big), \end{aligned} \tag{B5}$$

where $E_1^{W^-\text{-RF}}$ and $\mathbf{p}_1^{W^-\text{-RF}}$ are given, respectively, as

$$E_1^{W^-\text{-RF}} = \frac{q^2 + m_1^2 - p_N^2}{2\sqrt{q^2}}, \quad |\mathbf{p}_1^{W^-\text{-RF}}| = \frac{\sqrt{\lambda(q^2, m_1^2, p_N^2)}}{2\sqrt{q^2}}. \tag{B6}$$

The Lorentz boost matrix to transform four vectors from the W^--RF to the $\mathcal{B}_1$-RF reads

$$\Lambda_{W^-\to\mathcal{B}_1} = \begin{pmatrix} \gamma_1 & -\gamma_1\beta_{1x} & -\gamma_1\beta_{1y} & -\gamma_1\beta_{1z} \\ -\gamma_1\beta_{1x} & 1 + (\gamma_1 - 1)\frac{\beta_{1x}^2}{\vec{\beta}_1^2} & (\gamma_1 - 1)\frac{\beta_{1x}\beta_{1y}}{\vec{\beta}_1^2} & (\gamma_1 - 1)\frac{\beta_{1x}\beta_{1z}}{\vec{\beta}_1^2} \\ -\gamma_1\beta_{1y} & (\gamma_1 - 1)\frac{\beta_{1x}\beta_{1y}}{\vec{\beta}_1^2} & 1 + (\gamma_1 - 1)\frac{\beta_{1y}^2}{\vec{\beta}_1^2} & (\gamma_1 - 1)\frac{\beta_{1y}\beta_{1z}}{\vec{\beta}_1^2} \\ -\gamma_1\beta_{1z} & (\gamma_1 - 1)\frac{\beta_{1x}\beta_{1z}}{\vec{\beta}_1^2} & (\gamma_1 - 1)\frac{\beta_{1y}\beta_{1z}}{\vec{\beta}_1^2} & 1 + (\gamma_1 - 1)\frac{\beta_{1z}^2}{\vec{\beta}_1^2} \end{pmatrix}, \tag{B7}$$

where the velocity $\vec{\beta}_1$ is the velocity of the $W^-(q)$ as seen in the $\mathcal{B}_1$-RF and

$$\gamma_1 = \frac{1}{\sqrt{1-\vec{\beta}_1^2}}, \qquad \beta_{1x} = 0, \qquad \beta_{1y} = 0, \qquad \beta_{1z} = \frac{|\mathbf{p}_{\mathcal{B}_2}^{\mathcal{B}_1\text{-RF}}|}{E_q^{\mathcal{B}_1\text{-RF}}}. \tag{B8}$$

In the heavy neutrino rest frame N-RF, the four momentum of $\ell_2(p_2)$ and the $W^+(p_\pi)$ are given as

$$p_2^{N\text{-RF}} \equiv (E_2^{N\text{-RF}}, |\mathbf{p}_2^{N\text{-RF}}|\sin\theta_2\cos\phi, |\mathbf{p}_2^{N\text{-RF}}|\sin\theta_2\sin\phi, |\mathbf{p}_2^{N\text{-RF}}|\cos\theta_2), \tag{B9}$$

$$p_\pi^{N\text{-RF}} \equiv (E_\pi^{N\text{-RF}}, -|\mathbf{p}_2^{N\text{-RF}}|\sin\theta_2\cos\phi, -|\mathbf{p}_2^{N\text{-RF}}|\sin\theta_2\sin\phi, -|\mathbf{p}_2^{N\text{-RF}}|\cos\theta_2), \tag{B10}$$

where θ_2 is the angle made by the lepton ℓ_2 with respect to the $+\hat{z}$ direction. The angle ϕ is made by the plane containing ℓ_2 and W^+ with respect to the plane containing ℓ_1 and N as seen in the $\mathcal{B}_1$-RF and is defined as

$$\cos\phi = (\hat{p}_1^{\mathcal{B}_1\text{-RF}} \times \hat{p}_N^{\mathcal{B}_1\text{-RF}}).(\hat{p}_2^{\mathcal{B}_1\text{-RF}} \times \hat{p}_{W^+}^{\mathcal{B}_1\text{-RF}}), \tag{B11}$$

$$\sin\phi = -[(\hat{p}_1^{\mathcal{B}_1\text{-RF}} \times \hat{p}_N^{\mathcal{B}_1\text{-RF}}) \times (\hat{p}_2^{\mathcal{B}_1\text{-RF}} \times \hat{p}_{W^+}^{\mathcal{B}_1\text{-RF}})].\hat{p}_N^{\mathcal{B}_1\text{-RF}}. \tag{B12}$$

The energies $E_2^{N\text{-RF}}$ and $E_\pi^{N\text{-RF}}$ and three momentum $|\mathbf{p}_2^{N\text{-RF}}|$ are

$$E_2^{N\text{-RF}} = \frac{p_N^2 + m_2^2 - p_\pi^2}{2\sqrt{p_N^2}}, \qquad E_\pi^{N\text{-RF}} = \sqrt{p_N^2} - E_2^{N\text{-RF}}, \qquad |\mathbf{p}_2^{N\text{-RF}}| = \frac{\sqrt{\lambda(p_N^2, m_2^2, p_\pi^2)}}{2\sqrt{p_N^2}}. \tag{B13}$$

The Lorentz boost required to go from the heavy neutrino rest frame to the W^- rest frame is given as

$$\Lambda_{N\to W^-} = \begin{pmatrix} \gamma_2 & -\gamma_2\beta_{2x} & -\gamma_2\beta_{2y} & -\gamma_2\beta_{2z} \\ -\gamma_2\beta_{2x} & 1+(\gamma_2-1)\frac{\beta_{2x}^2}{\vec{\beta}_2^2} & (\gamma_2-1)\frac{\beta_{2x}\beta_{2y}}{\vec{\beta}_2^2} & (\gamma_2-1)\frac{\beta_{2x}\beta_{2z}}{\vec{\beta}_2^2} \\ -\gamma_2\beta_{2y} & (\gamma_2-1)\frac{\beta_{2x}\beta_{2y}}{\vec{\beta}_2^2} & 1+(\gamma_2-1)\frac{\beta_{2y}^2}{\vec{\beta}_2^2} & (\gamma_2-1)\frac{\beta_{2y}\beta_{2z}}{\vec{\beta}_2^2} \\ -\gamma_2\beta_{2z} & (\gamma_2-1)\frac{\beta_{2x}\beta_{2z}}{\vec{\beta}_2^2} & (\gamma_2-1)\frac{\beta_{2y}\beta_{2z}}{\vec{\beta}_2^2} & 1+(\gamma_2-1)\frac{\beta_{2z}^2}{\vec{\beta}_2^2} \end{pmatrix}, \tag{B14}$$

where the velocity $\vec{\beta}_2$ is the velocity of the N as seen in the W^--RF

$$\gamma_2 = \frac{1}{\sqrt{1-\vec{\beta}_2^2}}, \qquad \beta_{2x} = \frac{|\mathbf{p}_N^{W^-\text{-RF}}|}{E_N^{W^-\text{-RF}}}\sin\theta_1, \qquad \beta_{2y} = 0, \qquad \beta_{2z} = \frac{|\mathbf{p}_N^{W^-\text{-RF}}|}{E_N^{W^-\text{-RF}}}\cos\theta_1. \tag{B15}$$

APPENDIX C: PHASE SPACE

The differential width for the four-body final state is

$$d\Gamma = \frac{1}{2m_{\mathcal{B}_1}}|\mathcal{M}|^2 d\Phi_4(\mathcal{B}_1 \to \mathcal{B}_2\ell_1\ell_2\pi). \tag{C1}$$

The four-body phase space can be split up as

$$d\Phi_4(\mathcal{B}_1 \to \mathcal{B}_2\ell_1\ell_2\pi) = d\Phi_3(\mathcal{B}_1 \to \mathcal{B}_2\ell_1 N)\frac{dp_N^2}{2\pi}d\Phi_2(N \to \ell_2\pi), \quad \text{for } D \text{ channel}, \tag{C2}$$

$$= d\Phi_3(\mathcal{B}_1 \to \mathcal{B}_2\ell_2 N)\frac{dp_N^2}{2\pi}d\Phi_2(N \to \ell_1\pi), \quad \text{for } C\,\text{channel}. \tag{C3}$$

The expressions of different phase spaces are given below:

$$d\Phi_3(\mathcal{B}_1 \to \mathcal{B}_2 \ell_1 N) = \frac{\sqrt{\lambda(1, m_{\mathcal{B}_2}^2/m_{\mathcal{B}_2}^1, q^2/m_{\mathcal{B}_2}^1)}}{4\pi} \frac{\sqrt{\lambda(1, m_{\ell_1}^2/q^2, p_N^2/q^2)}}{(8\pi)^2} \int dq^2 d\cos\theta_1, \tag{C4}$$

$$d\Phi_2(N \to \ell_2 \pi) = \frac{\sqrt{\lambda(1, m_{\ell_2}^2/p_N^2, p_\pi^2/p_N^2)}}{8\pi} \int \frac{d\cos\theta_2}{2} \frac{d\phi}{2\pi}. \tag{C5}$$

The limits of the integration of the angles are as follows:

$$-1 \le \cos\theta_1 \le 1, \qquad -1 \le \cos\theta_2 \le 1, \qquad 0 \le \phi \le 2\pi. \tag{C6}$$

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