Bottomed baryon decays with invisible Majorana fermions

Postdoc. Geng Li 李更

Co-authors: Chia-Wei Liu, Chao-Qiang Geng

arxiv: 2206.01575, PRD publishing

December 9, 2022

國科大杭州高等研究院
Hangzhou Institute for Advanced Study,UCAS

Motivation

- Flavor-changing neutral current (FCNC) processes of long-lived particles provide a window to observe new physics (NP) beyond the standard model (SM).
  $c \to u$, $s \to d$, and $b \to d(s)$ transitions.

- Experiments have given the constraints on hadronic FCNC decays with missing energy ($\not{E}$).
  $\mathcal{B}(K_L \to \pi^0 \bar{\nu}\nu)_{\rm KOTO} < 3.0 \times 10^{-9}$ at 90% C.L.
  $\mathcal{B}(K_L \to \pi^0 \bar{\nu}\nu)_{\rm SM} = (3.4 \pm 0.6) \times 10^{-11}$
  $\mathcal{B}(K^+ \to \pi^+ \bar{\nu}\nu)_{\rm NA62} = (11.0^{+4.0}_{-3.5}({\rm stat}) \pm 0.3({\rm syst})) \times 10^{-11}$ at 68% C.L.
  $\mathcal{B}(K^+ \to \pi^+ \bar{\nu}\nu)_{\rm E949} = (17.3^{+11.5}_{-10.5}) \times 10^{-11}$
  $\mathcal{B}(K^+ \to \pi^+ \bar{\nu}\nu)_{\rm SM} = (8.4 \pm 1.0) \times 10^{-11}$

- The upper bounds of branching ratios of $B \to K^{(*)}, \pi, \rho$ modes have been given by the CLEO, BarBar, Belle and Belle II collaborations.

- Models of fermionic dark matter particles, such as sterile neutrino, neutralino, Higgs-portal, Z-portal, and singlet-doublet.

- Particularly, the Belle II: sensitivity for measurement of $B^{0(+)} \to K^{(*)0(+)} \bar{\nu}\nu$ can be increased by 25−30%. Future e^+e^- colliders, such as the FCC-ee experiment, have shown the ability of precise measurements of FCNC processes.

Experimental bounds

Table 1: The branching ratios ($\mathcal{B}$) (in units of 10^{-6}) of B decays involving missing energy.

Experimental bound[a]	SM prediction	Invisible particles bound
$\mathcal{B}(B^\pm \to K^\pm \not{E}) < 16$	$\mathcal{B}(B^\pm \to K^\pm \nu\bar{\nu}) = 4.73 \pm 0.56$	$\mathcal{B}(B^\pm \to K^\pm \chi\chi) < 11.8$
$\mathcal{B}(B^\pm \to \pi^\pm \not{E}) < 14$	$\mathcal{B}(B^\pm \to \pi^\pm \nu\bar{\nu}) = 8.12 \pm 0.01$	$\mathcal{B}(B^\pm \to \pi^\pm \chi\chi) < 5.89$
$\mathcal{B}(B^\pm \to K^{*\pm} \not{E}) < 40$	$\mathcal{B}(B^\pm \to K^{*\pm} \nu\bar{\nu}) = 8.93 \pm 1.07$	$\mathcal{B}(B^\pm \to K^{*\pm} \chi\chi) < 32.1$
$\mathcal{B}(B^\pm \to \rho^\pm \not{E}) < 30$	$\mathcal{B}(B^\pm \to \rho^\pm \nu\bar{\nu}) = 0.48 \pm 0.18$	$\mathcal{B}(B^\pm \to \rho^\pm \chi\chi) < 29.7$

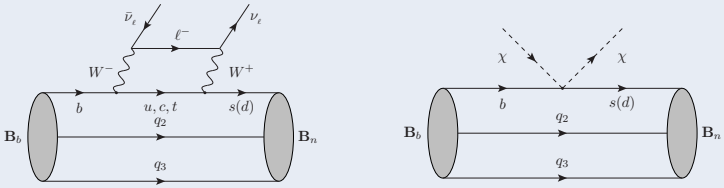
Phenomenologically, these new invisible fermions of χ can weakly interact with the SM fermions via a mediator, which can be a scalar, pseudoscalar, vector or axial-vector particle.

In our study, we will concentrate on a general model-independent approach to introduce the effective Lagrangian, which contains all possible currents involving the invisible fermions with the coupling constants extracted from the experiments.

[a] These experimental bounds are adopted by PDG-live, which are not certainly the latest or strictest constraints.

SM expectations

The FCNC decay processes of bottomed baryons with missing energy are described as



a: Within the Standard Model

b: Beyond the Standard Model

Figure 1: Feynman diagrams of bottomed baryon FCNC decays with missing energy.

$$\mathcal{L}_{\bar{\nu}\nu} = \frac{4G_F}{\sqrt{2}} \frac{\alpha}{2\pi \sin^2\theta_W} \sum_{\ell=e,\mu,\tau} \sum_{q=u,c,t} V_{bq} V_{sq} X^\ell(x_q) (\bar{s}_L \gamma^\mu b_L)(\bar{\nu}_{\ell L} \gamma_\mu \nu_{\ell L}),$$

The loop is calculated by Inami-Lim function:

$$X^\ell(x_q) = \frac{x_q}{8} \left[\frac{x_q + 2}{x_q - 1} + \frac{3(x_q - 2)}{(x_q - 1)^2} \ln x_q \right],$$

where G_F represents the Fermi coupling constant, α corresponds to the fine structure constant, θ_W stands for the Weinberg angle, V_{ij} are the Cabibbo–Kobayashi–Maskawa (CKM) matrix elements, and $x_q = m_q^2/M_W^2$ with m_q (M_W) being the mass of the quark (W-boson).

Baryonic amplitude

- Consequently, the transition amplitude is given by

$$\langle \mathbf{B}_n \bar{\nu}\nu | \mathcal{L}_{\bar{\nu}\nu} | \mathbf{B}_b \rangle = \frac{\sqrt{2} G_F \alpha}{4\pi \sin^2\theta_W} V_{bt} V_{st} X^{\ell}(x_t) \langle \mathbf{B}_n | \bar{s}\gamma^{\mu}(1-\gamma^5) b | \mathbf{B}_b \rangle \times \bar{u}_{\nu_\ell} \gamma_\mu (1-\gamma^5) v_{\nu_\ell}.$$

Baryonic form factors

- The baryonic transition matrix elements can be parameterized by the form factors (FFs) of $f_i^{V,A}$ $(i = 1, 2, 3)$, f^S and f^P, defined by

$$\langle \mathbf{B}_n(P_f, s_f) | (\bar{q}_f \gamma_\mu q) | \mathbf{B}_b(P, s) \rangle = \bar{u}_{\mathbf{B}_n}(P_f, s_f) \left[\gamma_\mu f_1^V(q^2) + i\sigma_{\mu\nu} \frac{q^\nu}{M} f_2^V(q^2) + \frac{q^\mu}{M} f_3^V(q^2) \right] u_{\mathbf{B}_b}(P, s),$$

$$\langle \mathbf{B}_n(P_f, s_f) | (\bar{q}_f q) | \mathbf{B}_b(P, s) \rangle = \bar{u}_{\mathbf{B}_n}(P_f, s_f) f^S(q^2) u_{\mathbf{B}_b}(P, s),$$

$$\langle \mathbf{B}_n(P_f, s_f) | (\bar{q}_f \gamma_\mu \gamma^5 q) | \mathbf{B}_b(P, s) \rangle = \bar{u}_{\mathbf{B}_n}(P_f, s_f) \left[\gamma_\mu f_1^A(q^2) + i\sigma_{\mu\nu} \frac{q^\nu}{M} f_2^A(q^2) + \frac{q^\mu}{M} f_3^A(q^2) \right] \gamma^5 u_{\mathbf{B}_b}(P, s),$$

$$\langle \mathbf{B}_n(P_f, s_f) | (\bar{q}_f \gamma^5 q) | \mathbf{B}_b(P, s) \rangle = \bar{u}_{\mathbf{B}_n}(P_f, s_f) f^P(q^2) \gamma^5 u_{\mathbf{B}_b}(P, s),$$

Modified bag model

- We will evaluate these elements in terms of the MBM, which works well for the heavy baryonic decays. In the MBM, the baryon wave functions at rest are read as

$$\Psi(x_{q_1}, x_{q_2}, x_{q_3}) = \mathcal{N} \int d^3\vec{x} \prod_{i=1,2,3} \phi_{q_i}(\vec{x}_{q_i} - \vec{x}) e^{-iE_{q_i} t_{q_i}} ,$$

  where q_i are the quark components of the baryons, $\mathcal{N}$ the overall normalization constant, x_{q_i} (E_{q_i}) the spacetime coordinates (energies) of q_i.

- $\phi_{q_i}(x)$ the quark wave functions inside a static bag, located at the center, given by

$$\phi_q(\vec{x}) = \begin{pmatrix} \omega_{q+} j_0(p_q r)\chi_q \\ i\omega_{q-} j_1(p_q r)\hat{r} \cdot \vec{\sigma}\chi_q \end{pmatrix} .$$

  Here, $j_{0,1}$ represent the spherical Bessel functions, $\omega_{q\pm} = \sqrt{T_q \pm M_q}$ with T_q the kinematic energies, and χ_q are the two component spinors.

- By demanding that quark currents shall not penetrate the boundary of bags, we have the boundary condition

$$\tan(p_q R) = \frac{p_q R}{1 - M_q R - E_q R} ,$$

  where R is the bag radius, resulting in that the magnitudes of 3-momenta are quantized, which can be analogous to the well-know infinite square well.

Modified bag model

- By sandwiching the operators, we arrive

$$\int \langle \Lambda | \bar{s}\Gamma b(x) e^{iqx} | \Lambda_b \rangle d^4x = \mathcal{Z} \int d^3\vec{x}_\Delta \Gamma_{sb}(\vec{x}_\Delta) \prod_{q_j=u,d} D_{q_j}(\vec{x}_\Delta) \,,$$

with

$$\mathcal{Z} \equiv (2\pi)^4 \delta^4(p_{\Lambda_b} - p_\Lambda - q) \mathcal{N}_{\Lambda_b} \mathcal{N}_\Lambda \,,$$

$$D_{q_j}(\vec{x}_\Delta) \equiv \sqrt{1-v^2} \int d^3\vec{x} \phi^\dagger_{q_j}\left(\vec{x} + \frac{1}{2}\vec{x}_\Delta\right) \phi_{q_j}\left(\vec{x} - \frac{1}{2}\vec{x}_\Delta\right) e^{-2iE_{q_j}\vec{v}\cdot\vec{x}} \,,$$

$$\Gamma_{sb}(\vec{x}_\Delta) = \int d^3\vec{x} \phi_s\left(\vec{x} + \frac{1}{2}\vec{x}_\Delta\right) \gamma^0 S_{-\vec{v}} \Gamma S_{-\vec{v}} \phi_b\left(\vec{x} - \frac{1}{2}\vec{x}_\Delta\right) e^{i(M_\Lambda + M_{\Lambda_b} - E_s - E_b)\vec{v}\cdot\vec{x}} \,,$$

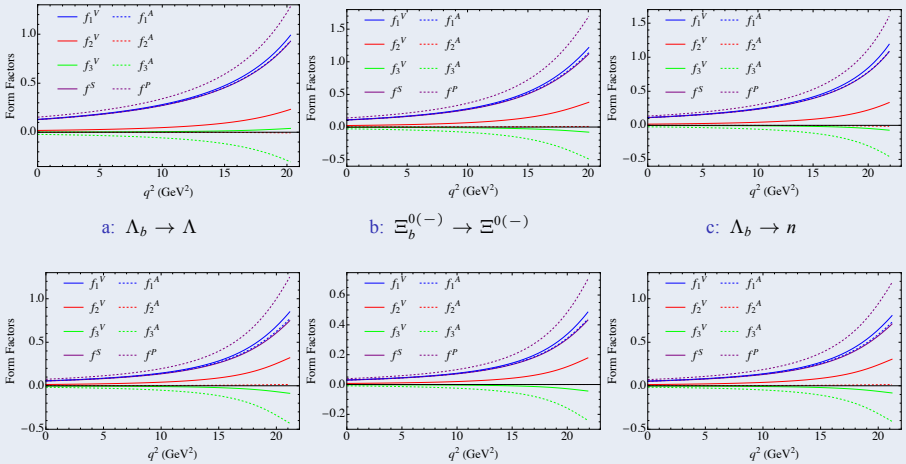
where Γ are arbitrary Dirac matrices, and $S_{\vec{v}}$ the Lorentz boost matrix of Dirac spinors. We have taken the initial (final) state as Λ_b (Λ) for an concrete example. To simplify the algebra, the Briet frame is chosen, where Λ_b and Λ have the velocity $-\vec{v}$ and $\vec{v}$, respectively.

- Notably, all the parameters of the model are extracted from the mass spectra, given as

$$R = 4.8 \text{ GeV}^{-1} \,, \quad M_{u,d} = 0 \,, \quad M_s = 0.28 \text{ GeV} \,, \quad M_b = 5.093 \text{ GeV} \,.$$

Form Factors from MBM

We consider the bottomed baryon decays of ($\Lambda_b \to \Lambda\bar{\nu}\nu$ and $\Xi_b^{0(-)} \to \Xi^{0(-)}\bar{\nu}\nu$) and ($\Lambda_b \to n\bar{\nu}\nu$, $\Xi_b^{0(-)} \to \Sigma^{0(-)}\bar{\nu}\nu$, and $\Xi_b^0 \to \Lambda\bar{\nu}\nu$), due to the ($b \to s$) and ($b \to d$) transitions at quark level.



a: $\Lambda_b \to \Lambda$ b: $\Xi_b^{0(-)} \to \Xi^{0(-)}$ c: $\Lambda_b \to n$

d: $\Xi_b^- \to \Sigma^-$ e: $\Xi_b^0 \to \Lambda$ f: $\Xi_b^0 \to \Sigma^0$

Figure 2: Form factors as function of q^2

SM expectations

- By integrating the three-body phase space, we obtain the decay branching ratio to be

$$\mathcal{B}(\mathbf{B}_b \to \mathbf{B}_n \bar{\nu}\nu) = \frac{1}{512\pi^3 M^3 \Gamma_{\mathbf{B}_b}} \int \frac{dq^2}{q^2} \lambda^{1/2}(M^2, q^2, M_f^2) \lambda^{1/2}(q^2, m_1^2, m_2^2) \int d\cos\theta \sum |\mathcal{M}|^2,$$

  where $\lambda(x,y,z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz$ is the Källen function. As the three generations of neutrinos are indistinguishable experimentally, the final results need to be multiplied by three.

- For the $b \to s$ transition, the decay branching ratios associated with $\bar{\nu}\nu$ are as follows:

$$\mathcal{B}(\Lambda_b \to \Lambda \bar{\nu}\nu) = 5.52^{+0.28}_{-0.28} \times 10^{-6},$$
$$\mathcal{B}(\Xi_b^{0(-)} \to \Xi^{0(-)} \bar{\nu}\nu) = 7.80^{+0.71}_{-0.67} \times 10^{-6}.$$

  Here, due to the SU(3) flavor symmetry, the branching ratios of Ξ_b^0 and Ξ_b^- are considered approximately to be equal.

- Similarly, for the $b \to d$ transition we have that

$$\mathcal{B}(\Lambda_b \to n \bar{\nu}\nu) = 2.76^{+0.17}_{-0.16} \times 10^{-7},$$
$$\mathcal{B}(\Xi_b^- \to \Sigma^- \bar{\nu}\nu) = 2.65^{+0.29}_{-0.26} \times 10^{-7},$$
$$\mathcal{B}(\Xi_b^0 \to \Sigma^0 \bar{\nu}\nu) = 1.24^{+0.13}_{-0.12} \times 10^{-7},$$
$$\mathcal{B}(\Xi_b^0 \to \Lambda \bar{\nu}\nu) = 3.88^{+0.37}_{-0.40} \times 10^{-8},$$

Effective Lagrangian with invisible particles

- Under the low energy scale, the model-independent effective Lagrangian is given by

$$\mathcal{L}_{eff} = \sum_{i=1}^{6} g_{mi} Q_i,$$

where g_{fi} are the phenomenological coupling constants. There are 6 independent dimension-six effective operators

$$Q_1 = (\bar{q}_f q)(\chi\chi), \qquad Q_2 = (\bar{q}_f \gamma^5 q)(\chi\chi), \qquad Q_3 = (\bar{q}_f q)(\chi\gamma^5\chi),$$
$$Q_4 = (\bar{q}_f \gamma^5 q)(\chi\gamma^5\chi), \quad Q_5 = (\bar{q}_f \gamma_\mu q)(\chi\gamma^\mu\gamma^5\chi), \quad Q_6 = (\bar{q}_f \gamma_\mu \gamma^5 q)(\chi\gamma^\mu\gamma^5\chi),$$

where the invisible particles of χ have been assumed to be the Majorana type. Since $\chi\gamma^\mu\chi = 0$ and $\chi\sigma^{\mu\nu}\chi = 0$, there is no contribution from the vector or tensor current.

- The upper limits of the coupling constants in the effective Lagrangian can be extracted from $B^- \to K^-(K^{*-}) + \not{E}$ and $B^- \to \pi^-(\rho^-) + \not{E}$ modes

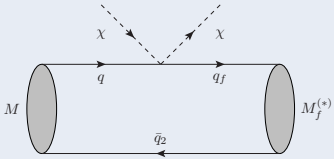


Figure 3: Feynman diagram of bottomed meson FCNC decays with invisible particles.

Mesonic amplitude

- For the $0^- \to 0^-$ meson decays of $M^- \to M_f^- \chi\chi$, only operators $Q_{1,3,5}$ give the contributions. The amplitudes of the $0^- \to 0^-$ decays can be simplified as

$$\langle M_f^- \chi\chi|\mathcal{L}_{eff}|M^-\rangle = 2g_{m1}\langle M_f^-|(\bar{q}_f q)|M^-\rangle \bar{u}_\chi v_\chi + 2g_{m3}\langle M_f^-|(\bar{q}_f q)|M^-\rangle \bar{u}_\chi \gamma^5 v_\chi + 2g_{m5}\langle M_f^-|(\bar{q}_f \gamma_\mu q)|M^-\rangle \bar{u}_\chi \gamma^\mu \gamma^5 v_\chi,$$

- For the $0^- \to 1^-$ meson decays of $M^- \to M_f^{*-} \chi\chi$ only operators $Q_{2,4,5,6}$ give the contributions. The amplitudes of the $0^- \to 1^-$ decays can be simplified as

$$\langle M_f^{*-} \chi\chi|\mathcal{L}_{eff}|M^-\rangle = 2g_{m2}\langle M_f^{*-}|(\bar{q}_f \gamma^5 q)|M^-\rangle \bar{u}_\chi v_\chi + 2g_{m4}\langle M_f^{*-}|(\bar{q}_f \gamma^5 q)|M^-\rangle \bar{u}_\chi \gamma^5 v_\chi + 2g_{m5}\langle M_f^{*-}|(\bar{q}_f \gamma_\mu q)|M^-\rangle \bar{u}_\chi \gamma^\mu \gamma^5 v_\chi + 2g_{m6}\langle M_f^{*-}|(\bar{q}_f \gamma_\mu \gamma^5 q)|M^-\rangle \bar{u}_\chi \gamma^\mu \gamma^5 v_\chi.$$

Mesonic form factors

- The hadronic transition matrix elements can be expressed as

$$\langle M_f^- | (\bar{q}_f q) | M^- \rangle = \frac{M^2 - M_f^2}{m_q - m_{q_f}} f_0(q^2),$$

$$\langle M_f^- | (\bar{q}_f \gamma_\mu q) | M^- \rangle = (P + P_f)_\mu f_+(q^2) + (P - P_f)_\mu \frac{M^2 - M_f^2}{q^2} \left[f_0(q^2) - f_+(q^2) \right],$$

$$\langle M_f^- | (\bar{q}_f \sigma_{\mu\nu} q) | M^- \rangle = i \left[P_\mu (P - P_f)_\nu - P_\nu (P - P_f)_\mu \right] \frac{2}{M + M_f} f_T(q^2),$$

and

$$\langle M_f^{*-} | (\bar{q}_f \gamma^5 q) | M^- \rangle = -i \left[\epsilon \cdot (P - P_f) \right] \frac{2M_f}{m_q + m_{q_f}} A_0(q^2),$$

$$\langle M_f^{*-} | (\bar{q}_f \gamma_\mu \gamma^5 q) | M^- \rangle = i \Big\{ \epsilon_\mu (M + M_f) A_1(q^2) - (P + P_f)_\mu \frac{\epsilon \cdot (P - P_f)}{M + M_f} A_2(q^2)$$
$$- (P - P_f)_\mu \left[\epsilon \cdot (P - P_f) \right] \frac{2M_f}{q^2} \left[A_3(q^2) - A_0(q^2) \right] \Big\},$$

$$\langle M_f^{*-} | (\bar{q}_f \gamma_\mu q) | M^- \rangle = \varepsilon_{\mu\nu\rho\sigma} \epsilon^\nu P^\rho (P - P_f)^\sigma \frac{2}{M + M_f} V(q^2),$$

where f_j $(j = 0, +, T)$, A_k $(k = 0 - 3)$ and V are the FFs, which are evaluated from the method of the LCSR, and ϵ is the polarization vector of the final meson with the convention of $\varepsilon^{0123} = 1$.

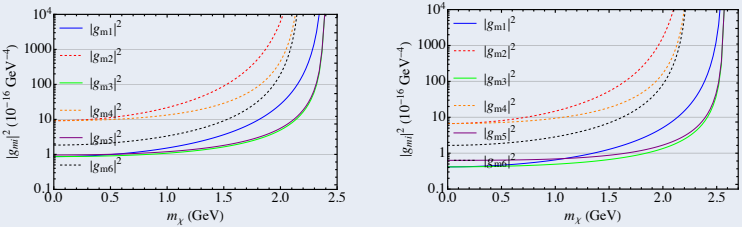
Coupling constants

- In our calculation, we assume that only one operator contributes to the process at a time. The upper limits of the coupling constants g_{mi} can be obtained from Table 1, given by

$$\mathcal{B}(M \to M_f^{(*)} \not{E})_{\rm exp} - \mathcal{B}(M \to M_f^{(*)} \bar{\nu}\nu)_{\rm SM} \geq \mathcal{B}(M \to M_f^{(*)} \chi\chi)_{Q_i} = \frac{|g_{mi}|^2 \widetilde{\Gamma}_{ii}}{\Gamma_{M_B}},$$

  Notably, the partial decay width should be divided by two since the Majorana fermion is identical to its antiparticle.

- The upper limits of $|g_{mi}|^2$ on $(bs\chi\chi)$ and $(bd\chi\chi)$ vertices are shown as functions of m_χ in Fig. 4 with m_χ is the mass of χ.



a: $(bs\chi\chi)$ vertex

b: $(bd\chi\chi)$ vertex

Figure 4: Upper limits of $|g_{mi}|^2$ as functions of m_χ

Coupling constants

- For the baryonic decays of $\mathbf{B}_b \to \mathbf{B}_n \chi\chi$, all operators should be considered.

$$\mathcal{L}_{eff} = \sum_{i=1}^{6} g_{mi} Q_i,$$

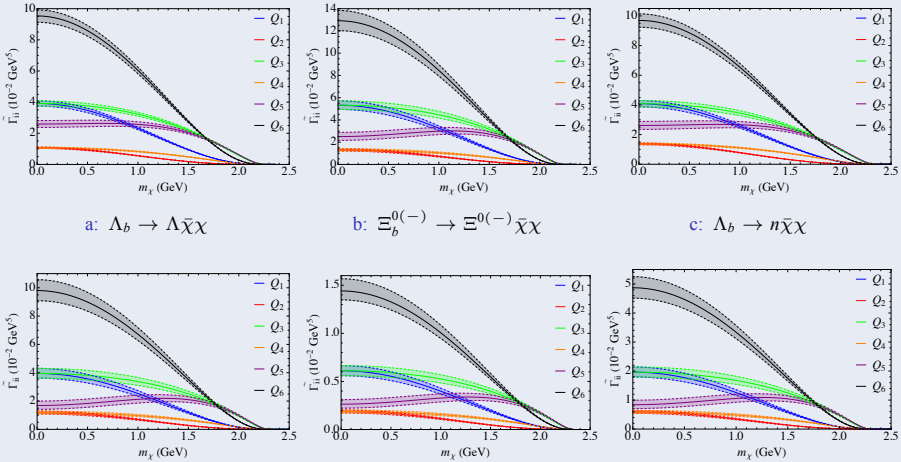
$$Q_1 = (\bar{q}_f q)(\chi\chi), \qquad Q_2 = (\bar{q}_f \gamma^5 q)(\chi\chi), \qquad Q_3 = (\bar{q}_f q)(\chi\gamma^5\chi),$$
$$Q_4 = (\bar{q}_f \gamma^5 q)(\chi\gamma^5\chi), \quad Q_5 = (\bar{q}_f \gamma_\mu q)(\chi\gamma^\mu\gamma^5\chi), \quad Q_6 = (\bar{q}_f \gamma_\mu \gamma^5 q)(\chi\gamma^\mu\gamma^5\chi),$$

- The decay amplitude can be expressed as

$$\begin{aligned}\langle \mathbf{B}_n \chi\chi | \mathcal{L}_{eff} | \mathbf{B}_b \rangle &= 2g_{m1} \langle \mathbf{B}_n | (\bar{q}_f q) | \mathbf{B}_b \rangle \bar{u}_\chi v_\chi + 2g_{m2} \langle \mathbf{B}_n | (\bar{q}_f \gamma^5 q) | \mathbf{B}_b \rangle \bar{u}_\chi v_\chi \\ &+ 2g_{m3} \langle \mathbf{B}_n | (\bar{q}_f q) | \mathbf{B}_b \rangle \bar{u}_\chi \gamma^5 v_\chi + 2g_{m4} \langle \mathbf{B}_n | (\bar{q}_f \gamma^5 q) | \mathbf{B}_b \rangle \bar{u}_\chi \gamma^5 v_\chi \\ &+ 2g_{m5} \langle \mathbf{B}_n | (\bar{q}_f \gamma_\mu q) | \mathbf{B}_b \rangle \bar{u}_\chi \gamma^\mu \gamma^5 v_\chi + 2g_{m6} \langle \mathbf{B}_n | (\bar{q}_f \gamma_\mu \gamma^5 q) | \mathbf{B}_b \rangle \bar{u}_\chi \gamma^\mu \gamma^5 v_\chi.\end{aligned}$$

Partial width

$\widetilde{\Gamma}_{ii}$ defined above are obtained with the numerical results



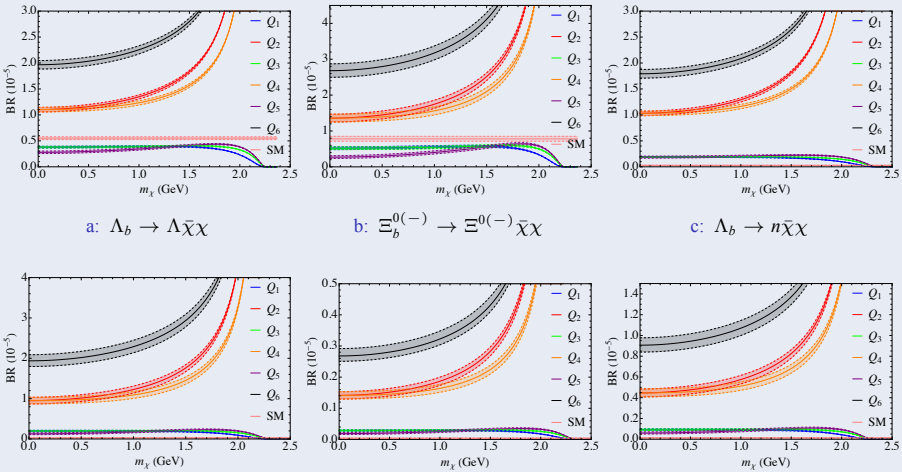
a: $\Lambda_b \to \Lambda \bar{\chi}\chi$ b: $\Xi_b^{0(-)} \to \Xi^{0(-)} \bar{\chi}\chi$ c: $\Lambda_b \to n \bar{\chi}\chi$

d: $\Xi_b^- \to \Sigma^- \bar{\chi}\chi$ e: $\Xi_b^0 \to \Lambda \bar{\chi}\chi$ f: $\Xi_b^0 \to \Sigma^0 \bar{\chi}\chi$

Figure 5: $\widetilde{\Gamma}_{ij}$ as function of m_χ in $b \to d$ decays

Branching ratios

The upper limits of the branching ratios are as shown



a: $\Lambda_b \to \Lambda\bar{\chi}\chi$ b: $\Xi_b^{0(-)} \to \Xi^{0(-)}\bar{\chi}\chi$ c: $\Lambda_b \to n\bar{\chi}\chi$

d: $\Xi_b^- \to \Sigma^-\bar{\chi}\chi$ e: $\Xi_b^0 \to \Lambda\bar{\chi}\chi$ f: $\Xi_b^0 \to \Sigma^0\bar{\chi}\chi$

Figure 6: The upper limits of the $\mathcal{BR}$ as function of m_χ in bottomed baryon FCNC decays

Branching ratios

Table 2: Upper limits of $\mathcal{B}(\mathbf{B}_b \to \mathbf{B}_n \chi\chi)$ when $m_\chi = 2$ GeV (in units of 10^{-5})

Operator	$\Lambda_b \to \Lambda\chi\chi$	$\Xi_b^{0(-)} \to \Xi^{0(-)}\chi\chi$	$\Lambda_b \to n\chi\chi$	$\Xi_b^- \to \Sigma^-\chi\chi$	$\Xi_b^0 \to \Sigma^0\chi\chi$	$\Xi_b^0 \to \Lambda\chi\chi$
Q_1	0.22	0.33	0.096	0.091	0.042	0.017
Q_2	5.3	7.3	5.2	4.4	2.0	0.93
Q_3	0.32	0.49	0.14	0.15	0.071	0.026
Q_4	3.6	5.4	3.6	3.3	1.6	0.61
Q_5	0.38	0.57	0.19	0.20	0.091	0.032
Q_6	6.3	9.2	5.7	5.8	2.7	1.0

Summary

- We have studied the light invisible Majorana fermions in the FCNC processes of the long-lived bottomed baryons.
- The model-independent effective Lagrangian which contains six operators has been introduced to describe the couplings between the quarks and invisible Majorana fermions.
- The bounds of the coupling constants have been extracted from the differences between the experimental upper limits and SM predictions of the relevant B meson FCNC decays.
- Based on these bounds, we have predicted the upper limits of $\mathcal{B}(\mathbf{B}_b \to \mathbf{B}_n \chi\chi)$.
- We are looking forward to the future experiments, such as those at Belle II, to get more measurements on bottomed baryons to find signs of new particles.

Prospect

- Invisible particles in $c \to u$ FCNC transition?
- First experimental bound on charmed hadron by BES III in April 2022.
  $\mathcal{B}(D^0 \to \pi^0 \bar{\nu}\nu) < 2.1 \times 10^{-4}$
- In the SM, strongly suppressed by the Glashow-Iliopoulos-Maiani (GIM) mechanism.
  $\mathcal{B}(D^0 \to \pi^0 \bar{\nu}\nu) \sim 10^{-15}$

Thanks for your attention!

Let's have a further discussion, please.

國科大杭州高等研究院
Hangzhou Institute for Advanced Study,UCAS

December 9, 2022