

Chiral perturbation theory for baryon-number-violating nucleon decay into a vector meson

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(Received 14 June 2025; accepted 11 August 2025; published 28 August 2025)

In a recent work [New chiral structures for baryon number violating nucleon decays, [arXiv:2504.14855](https://arxiv.org/abs/2504.14855).] we identified generic baryon-number-violating (BNV) structures containing triple light quarks and achieved their leading-order chiral realizations involving octet pseudoscalars and baryons. Although many two-body nucleon decays into a vector meson have been experimentally searched for and stringently constrained, a consistent theoretical framework for their calculation is still lacking. In this Letter, we fill the gap by implementing chiral matching of all these triple-quark interactions onto hadronic interactions involving octet vector mesons, baryons, and pseudoscalars. This paves the way for a consistent and comprehensive study of all relevant BNV processes. As an illustration of application, we show how degeneracy in the parameter space of Wilson coefficients can be broken by synthesizing experimental constraints on nucleon decays into a vector or pseudoscalar meson when relevant hadronic low-energy constants can be reasonably determined.

DOI: [10.1103/fzq9-tfcp](https://doi.org/10.1103/fzq9-tfcp)

Introduction. Baryon-number-violating (BNV) interactions are of fundamental importance in particle physics and cosmology due to their connections to major unresolved questions, such as the origin of matter-antimatter asymmetry in the Universe [1], the nature of dark matter, and the mechanism for tiny neutrino mass. Investigating these interactions can provide critical insights into physics beyond the standard model and guide the search for new particles and forces.

Nucleon decay that violates baryon number by one unit ($\Delta B = 1$) serves as a golden channel for probing BNV interactions, owing to the vast number of nucleons available in terrestrial detectors. The two-body nucleon N (either a proton p or a neutron n) decays to an octet pseudoscalar meson M and a lepton l (either a charged lepton ℓ or a neutrino ν) have been extensively searched for by large-fiducial-mass experiments, including IMB [2], SNO+ [3], KamLAND [4], Kamiokande [5], and Super-Kamiokande [6]. Although no positive signals have been observed, these

experiments have established stringent limits on the occurrence of such decays. Besides the pseudoscalar modes, significant experimental endeavors have been undertaken in nucleon decays into an octet vector meson [7–11]. Their experimental limits are comparably stringent to pseudoscalar channels; see the second column of Table I. The next generation of neutrino experiments, such as DUNE [12], Hyper-Kamiokande [13], JUNO [14], and THEIA [15], are expected to further improve sensitivity to these and other nucleon decay modes.

Theoretically, effective field theory (EFT) has emerged as an increasingly vital tool for investigating nucleon decay in a model-independent manner. In a recent work [16], we have examined the most general $\Delta B = 1$ interactions in the low-energy effective field theory (LEFT) that involve triple fields of light quarks. We found that they are classified into four irreducible representations (irreps) under the QCD chiral group $G_\chi \equiv \text{SU}(3)_L \times \text{SU}(3)_R$: $\mathbf{8}_L \otimes \mathbf{1}_R$, $\bar{\mathbf{3}}_L \otimes \mathbf{3}_R$, $\mathbf{6}_L \otimes \mathbf{3}_R$, $\mathbf{10}_L \otimes \mathbf{1}_R$, plus their chirality partners under the interchange of chiralities $L \leftrightarrow R$. Restricting to dimension-6 (dim-6) operators in LEFT, only the first two irreps (and their chirality partners) are possible, and correspond to the ones known since the 1980s [17–23], while the other two are new and can only be realized by higher dimensional operators. We have also achieved in that work the leading-order (LO) chiral realizations for all of the above irreps that involve the nucleon and octet pseudoscalar mesons. Again, while the chiral realizations for the dim-6 operators

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TABLE I. Derived new bounds on the nucleon two-body decay modes involving a vector meson based on those of the corresponding two-body pseudoscalar modes, under the assumption of single operator dominance. The modes marked with a $\times$ are suppressed by additional weak interactions or dim-9 interactions and are therefore neglected.

		Derived new bound	
Mode	Γ_{exp}^{-1} (10^{30} yr)	Γ_{new}^{-1} (10^{30} yr)	LEFT operator
$\Delta(B-L) = 0$			
$p \rightarrow \bar{\nu} \rho^+$	162	2.09×10^4	$\mathcal{O}_{\nu d u d}^{\text{LR},x}, \mathcal{O}_{\nu d u d}^{\text{LL},x}$
		6.04×10^7	$\mathcal{O}_{\nu d s u}^{\text{LR},x}, \mathcal{O}_{\nu d s u}^{\text{LL},x}$
$p \rightarrow \bar{\nu} K^{*+}$	51	9.67×10^5	$\mathcal{O}_{\nu s u d}^{\text{LR},x}, \mathcal{O}_{\nu s u d}^{\text{LL},x}$
		2.85×10^5	$\mathcal{O}_{\nu u d s}^{\text{LR},x}$
$p \rightarrow e^+ \rho^0$	720	9.10×10^5	$\mathcal{O}_{\ell u u d}^{\text{LR/RL},e}, \mathcal{O}_{\ell u u d}^{\text{LL/RR},e}$
$p \rightarrow \mu^+ \rho^0$	570	1.37×10^6	$\mathcal{O}_{\ell u u d}^{\text{LR/RL},\mu}, \mathcal{O}_{\ell u u d}^{\text{LL/RR},\mu}$
$p \rightarrow e^+ \omega$	1600	9.98×10^5	$\mathcal{O}_{\ell u u d}^{\text{LR/RL},e}, \mathcal{O}_{\ell u u d}^{\text{LL/RR},e}$
$p \rightarrow \mu^+ \omega$	2800	1.57×10^6	$\mathcal{O}_{\ell u u d}^{\text{LR/RL},\mu}, \mathcal{O}_{\ell u u d}^{\text{LL/RR},\mu}$
$p \rightarrow e^+ K^{*0}$	84	$1.50(0.59) \times 10^5$	$\mathcal{O}_{\ell u s u}^{\text{LR/RL},e} (\mathcal{O}_{\ell u s u}^{\text{LL/RR},e})$
$p \rightarrow e^+ \bar{K}^{*0}$	0.6	$\times$ (weak/5-quark operators)	
$n \rightarrow e^+ \rho^-$	217	4.48×10^5	$\mathcal{O}_{\ell u u d}^{\text{LR/RL},e}, \mathcal{O}_{\ell u u d}^{\text{LL/RR},e}$
$n \rightarrow \mu^+ \rho^-$	228	6.69×10^5	$\mathcal{O}_{\ell u u d}^{\text{LR/RL},\mu}, \mathcal{O}_{\ell u u d}^{\text{LL/RR},\mu}$
$n \rightarrow e^+ K^{*-}$	0.6	$\times$ (weak/5-quark operators)	
$n \rightarrow \bar{\nu} \rho^0$	19	4.13×10^4	$\mathcal{O}_{\nu d u d}^{\text{LR},x}, \mathcal{O}_{\nu d u d}^{\text{LL},x}$
$n \rightarrow \bar{\nu} \omega$	108	4.53×10^4	$\mathcal{O}_{\nu d u d}^{\text{LR},x}, \mathcal{O}_{\nu d u d}^{\text{LL},x}$
		8.46×10^4	$\mathcal{O}_{\nu d s u}^{\text{LR},x}, \mathcal{O}_{\nu d s u}^{\text{LL},x}$
$n \rightarrow \bar{\nu} K^{*0}$	78	1.22×10^6	$\mathcal{O}_{\nu s u d}^{\text{LR},x}, \mathcal{O}_{\nu s u d}^{\text{LL},x}$
		3.23×10^8	$\mathcal{O}_{\nu u d s}^{\text{LR},x}$
$\Delta(B+L) = 0$			
$p \rightarrow \nu \rho^+$	162	2.09×10^4	$\mathcal{O}_{\nu d u d}^{\text{RL},x}, \mathcal{O}_{\nu d u d}^{\text{RR},x}$
		6.04×10^7	$\mathcal{O}_{\nu d s u}^{\text{RL},x}, \mathcal{O}_{\nu d s u}^{\text{RR},x}$
$p \rightarrow \nu K^{*+}$	51	9.67×10^5	$\mathcal{O}_{\nu s u d}^{\text{RL},x}, \mathcal{O}_{\nu s u d}^{\text{RR},x}$
		2.84×10^5	$\mathcal{O}_{\nu u d s}^{\text{RL},x}$
$n \rightarrow e^- \rho^+$	62	$\times$ ($\Delta I = 3/2$)	
$n \rightarrow \mu^- \rho^+$	7		
$n \rightarrow e^- K^{*+}$	...	$3.10(1.61) \times 10^3$	$\mathcal{O}_{\ell d d s}^{\text{LR/RL},e} (\mathcal{O}_{\ell d d s}^{\text{LL/RR},e})$
$n \rightarrow \nu \rho^0$	19	4.13×10^4	$\mathcal{O}_{\nu d u d}^{\text{RL},x}, \mathcal{O}_{\nu d u d}^{\text{RR},x}$
$n \rightarrow \nu \omega$	108	4.53×10^4	$\mathcal{O}_{\nu d u d}^{\text{RL},x}, \mathcal{O}_{\nu d u d}^{\text{RR},x}$
		8.46×10^4	$\mathcal{O}_{\nu d s u}^{\text{RL},x}, \mathcal{O}_{\nu d s u}^{\text{RR},x}$
$n \rightarrow \nu K^{*0}$	78	1.22×10^6	$\mathcal{O}_{\nu s u d}^{\text{RL},x}, \mathcal{O}_{\nu s u d}^{\text{RR},x}$
		3.23×10^8	$\mathcal{O}_{\nu u d s}^{\text{RL},x}$

reproduce the known result in [24], the chiral realizations of the operators in the two new irreps are new and turn out to be nontrivial.

With the above experimental advances we think it is timely to develop a chiral perturbation theory for octet vector mesons and baryons. As we will analyze later, such a

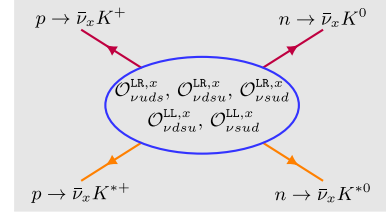


FIG. 1. An example of complementarity between pseudoscalar and vector meson modes of nucleon decay in constraining dim-6 BNV interactions involving a strange quark.

consistent theory is not available in the literature. As a bonus, this will facilitate a complementary exploration of BNV interactions through both pseudoscalar and vector meson channels. An example is shown in Fig. 1 for the dim-6 operators containing a strange quark that can be probed by two-body nucleon decays into either a pseudoscalar or a vector kaon.

General BNV triple quark interactions. To set the stage, we recall our recent results on BNV interactions involving triple quarks [16]. In LEFT or its extension by a new light field, a general BNV operator involving triple light quarks without being acted upon by a derivative always fall into one of the following four independent structures and their chirality partners,

$$\begin{aligned} \mathcal{O}_a^{yzw} &= \bar{\Psi}_a \mathcal{N}_{yzw}^{\text{LL}}, & \mathcal{O}_b^{yzw} &= \bar{\Psi}_b \mathcal{N}_{yzw}^{\text{RL}}, \\ \mathcal{O}_c^{yzw} &= \bar{\Psi}_{c,\mu} \mathcal{N}_{yzw}^{\text{LR},\mu}, & \mathcal{O}_d^{yzw} &= \bar{\Psi}_{d,\mu\nu} \mathcal{N}_{yzw}^{\text{LL},\mu\nu}, \end{aligned} \quad (1)$$

where the Ψ s are nonquark fields, and the $\mathcal{N}$ s are formed with three chiral quark fields in a definite irrep of G_χ ,

$$\mathcal{N}_{yzw}^{\text{LL}} = q_{L,y}^\alpha (\overline{q_{L,z}^{\beta C}} q_{L,w}^\gamma) \epsilon_{\alpha\beta\gamma} \in \mathbf{8}_L \otimes \mathbf{1}_R, \quad (2a)$$

$$\mathcal{N}_{yzw}^{\text{RL}} = q_{R,y}^\alpha (\overline{q_{L,z}^{\beta C}} q_{L,w}^\gamma) \epsilon_{\alpha\beta\gamma} \in \bar{\mathbf{3}}_L \otimes \mathbf{3}_R, \quad (2b)$$

$$\mathcal{N}_{yzw}^{\text{LR},\mu} = q_{L,\{y}^\alpha (\overline{q_{L,z}^{\beta C}} \gamma^\mu q_{R,w}^\gamma) \epsilon_{\alpha\beta\gamma} \in \mathbf{6}_L \otimes \mathbf{3}_R, \quad (2c)$$

$$\mathcal{N}_{yzw}^{\text{LL},\mu\nu} = q_{L,\{y}^\alpha (\overline{q_{L,z}^{\beta C}} \sigma^{\mu\nu} q_{L,w}^\gamma) \epsilon_{\alpha\beta\gamma} \in \mathbf{10}_L \otimes \mathbf{1}_R. \quad (2d)$$

Here α, β, γ denote color and $y, z, w = 1, 2, 3$ denote three light quarks with $q_{1,2,3} = u, d, s$. The curly brackets indicate symmetrization in flavor indices of like-chirality fields. When Ψ s are a single fermion field, i.e., when the $\mathcal{O}$ operators are of dim-6, $\mathcal{O}_{a,b}$ recover the two known operators while $\mathcal{O}_{c,d}$ vanish identically.

Following our previous work [25], we organize in matrix form the $\mathcal{N}$ s in the irreps $\mathbf{8}_L \otimes \mathbf{1}_R$ and $\bar{\mathbf{3}}_L \otimes \mathbf{3}_R$:

$$\mathcal{N}_{8_L \otimes 1_R} = \begin{pmatrix} \mathcal{N}_{uds}^{LL} & \mathcal{N}_{usu}^{LL} & \mathcal{N}_{uud}^{LL} \\ \mathcal{N}_{dds}^{LL} & \mathcal{N}_{dsu}^{LL} & \mathcal{N}_{dud}^{LL} \\ \mathcal{N}_{sds}^{LL} & \mathcal{N}_{ssu}^{LL} & \mathcal{N}_{sud}^{LL} \end{pmatrix}, \quad (3a)$$

$$\mathcal{N}_{\bar{3}_L \otimes \bar{3}_R} = \begin{pmatrix} \mathcal{N}_{uds}^{RL} & \mathcal{N}_{usu}^{RL} & \mathcal{N}_{uud}^{RL} \\ \mathcal{N}_{dds}^{RL} & \mathcal{N}_{dsu}^{RL} & \mathcal{N}_{dud}^{RL} \\ \mathcal{N}_{sds}^{RL} & \mathcal{N}_{ssu}^{RL} & \mathcal{N}_{sud}^{RL} \end{pmatrix}. \quad (3b)$$

Together with $\mathcal{N}_{yzw}^{LR,\mu}$ and $\mathcal{N}_{yzw}^{LL,\mu\nu}$ in Eq. (2), they transform under $(\hat{L}, \hat{R}) \in G_\chi$ as

$$\mathcal{N}_{8_L \otimes 1_R} \rightarrow \hat{L} \mathcal{N}_{8_L \otimes 1_R} \hat{L}^\dagger, \quad (4a)$$

$$\mathcal{N}_{\bar{3}_L \otimes \bar{3}_R} \rightarrow \hat{R} \mathcal{N}_{\bar{3}_L \otimes \bar{3}_R} \hat{R}^\dagger, \quad (4b)$$

$$\mathcal{N}_{yzw}^{LR,\mu} \rightarrow \hat{L}_{yy'} \hat{L}_{zz'} \hat{R}_{ww'} \mathcal{N}_{y'z'w'}^{LR,\mu}, \quad (4c)$$

$$\mathcal{N}_{yzw}^{LL,\mu\nu} \rightarrow \hat{L}_{yy'} \hat{L}_{zz'} \hat{L}_{ww'} \mathcal{N}_{y'z'w'}^{LL,\mu\nu}. \quad (4d)$$

The chirality partners of the $\mathcal{N}$ s along with their chiral transformations are obtained from the above by swapping L with R and $\hat{L}$ with $\hat{R}$.

In the Lagrangian of LEFT each operator $\mathcal{O}_i^{yzw}$ ($i = a, b, c, d$) is multiplied by a Wilson coefficient (WC) C_i^{yzw} which retains the same flavor symmetry as the operator. We denote the product of C_i^{yzw} and $\overline{\Psi}_i$ by the spurion field $\mathcal{P}_i^{yzw} = C_i^{yzw} \overline{\Psi}_i$. To facilitate chiral matching we assign chiral transformations to the spurion fields that are conjugate to those of $\mathcal{O}_i^{yzw}$, so that the effective interactions are apparently chiral invariant:

$$\mathcal{P}_{8_L \otimes 1_R} \rightarrow \hat{L} \mathcal{P}_{8_L \otimes 1_R} \hat{L}^\dagger, \quad (5a)$$

$$\mathcal{P}_{\bar{3}_L \otimes \bar{3}_R} \rightarrow \hat{R} \mathcal{P}_{\bar{3}_L \otimes \bar{3}_R} \hat{R}^\dagger, \quad (5b)$$

$$\mathcal{P}_{yzw}^{LR,\mu} \rightarrow \hat{L}_{yy'}^* \hat{L}_{zz'}^* \hat{R}_{ww'}^* \mathcal{P}_{y'z'w'}^{LR,\mu}, \quad (5c)$$

$$\mathcal{P}_{yzw}^{LL,\mu\nu} \rightarrow \hat{L}_{yy'}^* \hat{L}_{zz'}^* \hat{L}_{ww'}^* \mathcal{P}_{y'z'w'}^{LL,\mu\nu}, \quad (5d)$$

where $\mathcal{P}_{8_L \otimes 1_R}$ and $\mathcal{P}_{\bar{3}_L \otimes \bar{3}_R}$ are organized in matrix form, similar to Eq. (3) but with $\mathcal{N}$ replaced by $\mathcal{P}$ throughout. In Eq. (S3) of the Supplemental Material (SM) [26], we present the explicit forms of these spurion fields related to the dim-6 LEFT BNV interactions. To summarize, the general $\Delta B = 1$ Lagrangian involving triple light quarks can be compactly expressed as [16]

$$\begin{aligned} \mathcal{L}_q^\# = & \text{Tr}[\mathcal{P}_{8_L \otimes 1_R} \mathcal{N}_{8_L \otimes 1_R} + \mathcal{P}_{1_L \otimes 8_R} \mathcal{N}_{1_L \otimes 8_R}] \\ & + \text{Tr}[\mathcal{P}_{\bar{3}_L \otimes \bar{3}_R} \mathcal{N}_{\bar{3}_L \otimes \bar{3}_R} + \mathcal{P}_{\bar{3}_L \otimes \bar{3}_R} \mathcal{N}_{\bar{3}_L \otimes \bar{3}_R}] \\ & + [\mathcal{P}_{yzw}^{LR,\mu} \mathcal{N}_{yzw,\mu}^{LR} + \mathcal{P}_{yzw}^{RL,\mu} \mathcal{N}_{yzw,\mu}^{RL}] \\ & + [\mathcal{P}_{yzw}^{LL,\mu\nu} \mathcal{N}_{yzw,\mu\nu}^{LL} + \mathcal{P}_{yzw}^{RR,\mu\nu} \mathcal{N}_{yzw,\mu\nu}^{RR}] + \text{H.c.}, \quad (6) \end{aligned}$$

where summation over indices y, z, w is implied.

Chiral matching. To obtain the hadronic counterparts of the quark-level interactions in Eq. (6), a consistent approach is to use chiral perturbation theory (ChPT), which systematically accounts for interactions among the octet pseudo-scalars, vectors, and baryons. There are several approaches to incorporating the vector mesons: the matter field approach [27,28], the hidden local symmetry [29,30], the tensor field approach [28,31,32], and the massive Yang-Mills method [33–36]. It has been demonstrated in [28,37] that all these approaches are equivalent. In this work, we adopt the matter field approach in terms of four-vector potentials that transform homogeneously under the non-linear realization of the chiral symmetry.

We arrange the octet meson and baryon fields in the matrix form,

$$\Sigma(x) = \xi^2(x) = \exp\left(\frac{i\sqrt{2}\Pi(x)}{F_0}\right), \quad (7a)$$

$$\Pi(x) = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{6}} & \pi^+ & K^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} + \frac{\eta}{\sqrt{6}} & K^0 \\ K^- & \bar{K}^0 & -\sqrt{\frac{2}{3}}\eta \end{pmatrix}, \quad (7b)$$

$$B(x) = \begin{pmatrix} \frac{\Sigma^0}{\sqrt{2}} + \frac{\Lambda^0}{\sqrt{6}} & \Sigma^+ & p \\ \Sigma^- & -\frac{\Sigma^0}{\sqrt{2}} + \frac{\Lambda^0}{\sqrt{6}} & n \\ \Xi^- & \Xi^0 & -\sqrt{\frac{2}{3}}\Lambda^0 \end{pmatrix}, \quad (7c)$$

$$V_\mu(x) = \begin{pmatrix} \frac{\rho_\mu^0}{\sqrt{2}} + \frac{\phi_\mu^{(8)}}{\sqrt{6}} & \rho_\mu^+ & K_\mu^{*+} \\ \rho_\mu^- & -\frac{\rho_\mu^0}{\sqrt{2}} + \frac{\phi_\mu^{(8)}}{\sqrt{6}} & K_\mu^{*0} \\ K_\mu^{*-} & \bar{K}_\mu^{*0} & -\sqrt{\frac{2}{3}}\phi_\mu^{(8)} \end{pmatrix}, \quad (7d)$$

where $F_0 = (86.2 \pm 0.5)$ MeV is the pion decay constant in the chiral limit. Their chiral transformations are $\Sigma \rightarrow \hat{L}\Sigma\hat{R}^\dagger$, $B \rightarrow \hat{h}B\hat{h}^\dagger$, $V_\mu \rightarrow \hat{h}V_\mu\hat{h}^\dagger$, $\xi \rightarrow \hat{L}\xi\hat{h}^\dagger = \hat{h}\xi\hat{R}^\dagger$, where the matrix $\hat{h}$ is a function of $\hat{L}, \hat{R}$, and ξ . Note that for $X = B, V_\mu$, under the above chiral transformation: $\xi X \xi \rightarrow \hat{L}(\xi X \xi)\hat{R}^\dagger$, $\xi^\dagger X \xi^\dagger \rightarrow \hat{R}(\xi^\dagger X \xi^\dagger)\hat{L}^\dagger$, $\xi X \xi^\dagger \rightarrow \hat{L}(\xi X \xi^\dagger)\hat{L}^\dagger$, and $\xi^\dagger X \xi \rightarrow \hat{R}(\xi^\dagger X \xi)\hat{R}^\dagger$. For the neutral vector meson

$\phi_\mu^{(8)}$, we adopt the ideal mixing in terms of the physical states as $\phi_\mu^{(8)} = \omega_\mu/\sqrt{3} - \phi_\mu\sqrt{2/3}$.

The chiral building blocks consist of the spurion fields, the octet meson and baryon fields, and their derivatives. To identify the chiral leading terms, we adopt the chiral power counting scheme in which $\{\Sigma, \xi, V_\mu, B, D_\mu B\} \sim \mathcal{O}(p^0)$ and $\{D_\mu X_\nu, D_\mu \Sigma\} \sim \mathcal{O}(p^1)$. The covariant derivatives $D_\mu \Sigma$ and $D_\mu X$ ($X = B, V_\nu$) are defined as $D_\mu \Sigma = \partial_\mu \Sigma - i l_\mu \Sigma + i \Sigma r_\mu$ and $D_\mu X = \partial_\mu X + [\Gamma_\mu, X]$, where the external sources l_μ and r_μ are traceless matrices in the flavor space and $\Gamma_\mu = (1/2)[\xi(\partial_\mu - i r_\mu)\xi^\dagger + \xi^\dagger(\partial_\mu - i l_\mu)\xi]$. For the irreps $\mathbf{6}_{L(R)} \otimes \mathbf{3}_{R(L)}$ and $\mathbf{10}_{L(R)} \otimes \mathbf{1}_{R(L)}$, we need to employ appropriate Lorentz projectors to extract their non-trivial Lorentz irreps correctly: $(1, 1/2)$ and $(1/2, 1)$, and $(3/2, 0)$ and $(0, 3/2)$. We found in Ref. [16] that, for the irreps $\mathbf{6}_{L(R)} \otimes \mathbf{3}_{R(L)}$, the $(1, 1/2)$ and $(1/2, 1)$ components of a general vector spinor can be extracted

by $\Gamma_{\mu\nu}^{L,R} \equiv (g_{\mu\nu} - \frac{1}{4}\gamma_\mu\gamma_\nu)P_{L,R}$, and for the irreps $\mathbf{10}_{L(R)} \otimes \mathbf{1}_{R(L)}$, we apply the following Lorentz projectors [38]:

$$\hat{\Gamma}_{\mu\nu\alpha\beta}^{L,R} \equiv \frac{1}{24}(2\{\sigma_{\mu\nu}, \sigma_{\alpha\beta}\} - [\sigma_{\mu\nu}, \sigma_{\alpha\beta}])P_{L,R}, \quad (8)$$

to extract $(\frac{3}{2}, 0)$ and $(0, \frac{3}{2})$ components of a tensor spinor.

Equipped with the above ingredients we studied in Ref. [16] the chiral realizations of the LEFT operators $\mathcal{O}_i$ by the octet baryons and pseudoscalars in the single-baryon sector, which are collected in SM [26]. It turned out that each operator has one chiral realization at its LO which brings in a low-energy constant (LEC) c_i . While two of them have been calculated by lattice QCD [39], the other two have to be estimated by naive dimensional analysis (NDA) [40,41]. Now we do the same thing in the single-baryon and single-vector sector. At the chiral LO of each operator $\mathcal{O}_i$, the result is

$$\begin{aligned} \mathcal{L}_{BV}^\# = & d_1 \text{Tr}[\mathcal{P}_{\mathbf{3}_L \otimes \mathbf{3}_R} \xi \gamma_\mu V^\mu B_R \xi - \mathcal{P}_{\mathbf{3}_L \otimes \mathbf{3}_R} \xi^\dagger \gamma_\mu V^\mu B_L \xi^\dagger] + d'_1 \text{Tr}[\mathcal{P}_{\mathbf{3}_L \otimes \mathbf{3}_R} \xi \gamma_\mu B_R V^\mu \xi - \mathcal{P}_{\mathbf{3}_L \otimes \mathbf{3}_R} \xi^\dagger \gamma_\mu B_L V^\mu \xi^\dagger] \\ & + d_2 \text{Tr}[\mathcal{P}_{\mathbf{8}_L \otimes \mathbf{1}_R} \xi \gamma_\mu V^\mu B_R \xi^\dagger - \mathcal{P}_{\mathbf{8}_L \otimes \mathbf{1}_R} \xi^\dagger \gamma_\mu V^\mu B_L \xi] + d'_2 \text{Tr}[\mathcal{P}_{\mathbf{8}_L \otimes \mathbf{1}_R} \xi \gamma_\mu B_R V^\mu \xi^\dagger - \mathcal{P}_{\mathbf{8}_L \otimes \mathbf{1}_R} \xi^\dagger \gamma_\mu B_L V^\mu \xi] \\ & + d_3 [\mathcal{P}_{yzi}^{LR,\mu} \Gamma_{\mu\nu}^L (\xi B_L \xi)_{yj} (\xi V^\nu \xi)_{zk} \epsilon_{ijk} - \mathcal{P}_{yzi}^{RL,\mu} \Gamma_{\mu\nu}^R (\xi^\dagger B_R \xi^\dagger)_{yj} (\xi^\dagger V^\nu \xi^\dagger)_{kz} \epsilon_{ijk}] \\ & + d'_3 [\mathcal{P}_{yzi}^{LR,\mu} \Gamma_{\mu\nu}^L (\xi V^\nu B_L \xi)_{yj} \Sigma_{zk} \epsilon_{ijk} - \mathcal{P}_{yzi}^{RL,\mu} \Gamma_{\mu\nu}^R (\xi^\dagger V^\nu B_R \xi^\dagger)_{yj} \Sigma_{kz}^* \epsilon_{ijk}] \\ & + d''_3 [\mathcal{P}_{yzi}^{LR,\mu} \Gamma_{\mu\nu}^L (\xi B_L V^\nu \xi)_{yj} \Sigma_{zk} \epsilon_{ijk} - \mathcal{P}_{yzi}^{RL,\mu} \Gamma_{\mu\nu}^R (\xi^\dagger B_R V^\nu \xi^\dagger)_{yj} \Sigma_{kz}^* \epsilon_{ijk}] \\ & + \frac{d_4}{\Lambda_\chi} [\mathcal{P}_{yzw}^{LL,\mu\nu} \hat{\Gamma}_{\mu\alpha\beta}^L (\xi D^\alpha B_L \xi)_{yi} \Sigma_{zj} (\xi V^\beta \xi)_{wk} \epsilon_{ijk} - \mathcal{P}_{yzw}^{RR,\mu\nu} \hat{\Gamma}_{\mu\alpha\beta}^R (\xi^\dagger D^\alpha B_R \xi^\dagger)_{yi} \Sigma_{jz}^* (\xi V^\beta \xi)_{kw}^* \epsilon_{ijk}] + \text{H.c.}, \quad (9) \end{aligned}$$

where d_i s are unknown LECs. They are estimated in the SM [26] by NDA to be $d_{1,2,3,4}^{(r,n)} \sim \Lambda_\chi^2/(4\pi) \approx 0.115 \text{ GeV}^2$. Note that except for the irreps $\mathbf{10}_{L(R)} \otimes \mathbf{1}_{R(L)}$ all other irreps have more than one chiral realizations at each of their LO. Equation (9) is the starting point of all phenomenological analyses on BNV processes in the single-baryon and single-vector meson sector involving any number of pseudoscalars.

We are now in a position to compare our approach with a formalism developed for the vector meson case over 40 years ago [42]. In that work, the authors treated the vector mesons as the massive Yang-Mills fields of the chiral group and extended the results of Ref. [24] by simply replacing the baryon octet field B by $i\gamma_\mu D^\mu B$ to obtain the vector meson Lagrangian for the dim-6 BNV interactions with $\Delta(B-L) = 0$ in $\mathbf{3}_{L(R)} \otimes \mathbf{3}_{R(L)}$ and $\mathbf{8}_{L(R)} \otimes \mathbf{1}_{R(L)}$. The vector mesons enter through the covariant derivative by trading the external sources for them. This replacement leads to double counting for the pseudoscalars since it contributes exactly the same pure pseudoscalar BNV Lagrangian terms. From the EFT perspective, their above

operators can be reduced by using the equation of motion for the octet baryon field or its redefinition, to those appearing already in the pseudoscalar case in the absence of the octet vector [24]. Furthermore, there are additional operators involving covariant derivatives of pseudoscalar mesons that were missed in that work. Finally, we incorporated the chiral realizations of the triple-quark operators in the irreps $\mathbf{6}_{L(R)} \otimes \mathbf{3}_{R(L)}$ and $\mathbf{10}_{L(R)} \otimes \mathbf{1}_{R(L)}$ first pointed out in [16]. Our results are fully based on the consistent representation of vector mesons developed in the seminal works [27,28], thereby providing a consistent framework for matching these BNV interactions onto hadronic states involving vector mesons.

Nucleon decays into a vector meson. To apply the formalism developed in Eq. (9) and Eq. (S2) in the SM [26] to calculate specific decay, it is necessary to identify relevant quark-level interactions first. As an illustration, we focus on the leading dim-6 BNV operators in the LEFT, as summarized in Table S1 in the SM [26]. In addition to these BNV interactions, the standard ChPT interactions are also

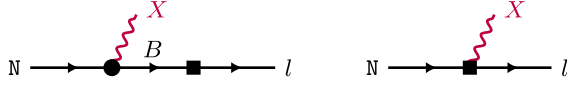


FIG. 2. Diagrams for nucleon BNV two-body decay involving a vector meson in the final state. The black square (blob) represents the insertion of a BNV (usual) vertex.

required. At LO, the relevant baryon-pseudoscalar (vector) interactions include the conventional $D(G_D)$ and $F(G_F)$ terms [43,44],

$$\mathcal{L}_{\text{ChPT}}^{B,0} \supset \frac{D}{2} \text{Tr}(\bar{B} \gamma^\mu \gamma_5 \{u_\mu, B\}) + \frac{F}{2} \text{Tr}(\bar{B} \gamma^\mu \gamma_5 [u_\mu, B]) + G_D \text{Tr}(\bar{B} \gamma_\mu \{V^\mu, B\}) + G_F \text{Tr}(\bar{B} \gamma_\mu [V^\mu, B]), \quad (10)$$

where $u_\mu = i[\xi(\partial_\mu - ir_\mu)\xi^\dagger - \xi^\dagger(\partial_\mu - il_\mu)\xi]$. For the LECs D and F , we adopt the recent lattice results $D = 0.730(11)$ and $F = 0.447_7^6$ [45]. And for the LECs associated with the vector mesons, we use the estimates given in [44], $G_D = 0$ and $G_F = g/\sqrt{2}$ where the factor $1/\sqrt{2}$ accounts for different parametrization of the interactions. g is the self-coupling of vector mesons and is related to the ρ -meson mass and pion decay constant via the well-known Kawarabayashi–Suzuki–Riazuddin–Fayyazuddin relation $g^2 = m_\rho^2/(2F_0^2)$ [46,47].

For the nucleon two-body decay modes, $N \rightarrow lX$, there are two leading-order Feynman diagrams as shown in Fig. 2. The relevant interactions are as follows:

$$\mathcal{L}_{N \rightarrow lX} = C_{N \rightarrow BX} \bar{X}^\mu \bar{B} \gamma_\mu N + C_{Bl}^L \bar{l} B_L + C_{Bl}^R \bar{l} B_R + C_{NlX}^L \bar{X}^\mu \bar{l} \gamma_\mu N_L + C_{NlX}^R \bar{X}^\mu \bar{l} \gamma_\mu N_R, \quad (11)$$

where the coefficients are summarized in Table S2 in the SM [26] for each process, based on the dim-6 LEFT BNV interactions. The decay width is

$$\Gamma_{N \rightarrow lX} = \frac{m_N \sqrt{\lambda_X}}{32\pi} [(\lambda_X x_X^{-1} + 3(1 + x_l - x_X)) \times (|\tilde{C}_{NlX}^L|^2 + |\tilde{C}_{NlX}^R|^2) - 12\sqrt{x_l} \Re(\tilde{C}_{NlX}^L \tilde{C}_{NlX}^{R*})], \quad (12)$$

with $\lambda_X = 1 + x_l^2 + x_X^2 - 2x_l - 2x_X - 2x_l x_X$, $x_l = m_l^2/m_N^2$, and $x_X = m_X^2/m_N^2$. $\tilde{C}_{NlX}^{L,R}$ are related to the parameters in Eq. (11) via

$$\tilde{C}_{NlX}^{L(R)} = C_{NlX}^{L(R)} + \sum_B C_{N \rightarrow BX} \frac{m_B C_{Bl}^{L(R)} + m_l C_{Bl}^{L(R)}}{m_B^2 - m_l^2}, \quad (13)$$

with a summation over all intermediate baryon B .

The decay widths for the pseudoscalar modes generally depend on multiple WCs at leading dim-6 order that are modulated by the LECs $c_{1,2}$. Without assuming single-operator dominance, one can only probe an infinite band or cone in the multiple-dimensional parameter space. For

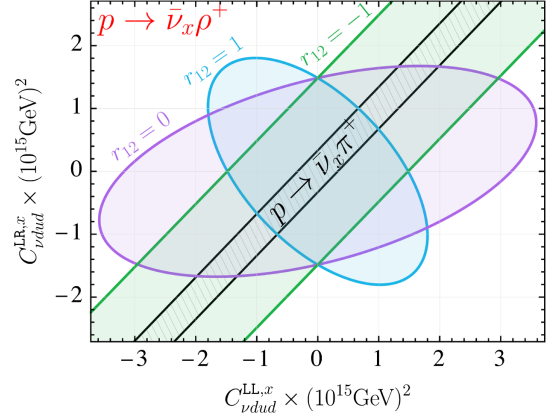


FIG. 3. An example of vector meson mode in breaking the degeneracy of parameter space in the $C_{\nu d u d}^{LR,x} - C_{\nu d u d}^{LL,x}$ plane.

convenience and completeness we summarize our results in Sec. IV and Table S3 in the SM [26]. The decay widths to a vector meson involve the same set of WCs which however are additionally modulated by new LECs $d_{1,2}$. This makes it possible to break the degeneracy with the pseudoscalar modes alone, and constrains the parameter space into a closed region. For the purpose of illustration, we consider the decays $p \rightarrow \bar{\nu}_x \pi^+$ and $p \rightarrow \bar{\nu}_x \rho^+$, both of which depend on the two WCs, $C_{\nu d u d}^{LR,x}$ and $C_{\nu d u d}^{LL,x}$. From Eq. (12), as well as Eq. (S8) and Tables S2 and S3 in the SM [26], we have

$$\Gamma_{p \rightarrow \bar{\nu}_x \rho^+} = \frac{m_p(1 - 3x_\rho^2 + 2x_\rho^3)}{32\pi x_\rho} (|d_1 C_{\nu d u d}^{LR,x} + d_2 C_{\nu d u d}^{LL,x}|^2 + G_F^2 m_n^{-2} |c_1 C_{\nu d u d}^{LR,x} + c_2 C_{\nu d u d}^{LL,x}|^2), \quad (14a)$$

$$\Gamma_{p \rightarrow \bar{\nu}_x \pi^+} = \frac{m_p(1 - x_\pi)^2 [1 + (D + F)m_p m_n^{-1}]^2}{64\pi F_0^2} \times |c_1 C_{\nu d u d}^{LR,x} + c_2 C_{\nu d u d}^{LL,x}|^2. \quad (14b)$$

The vector mode indeed depends on the two WCs in a different combination whenever $d_2/d_1 \neq c_2/c_1$. Once the new LECs $d_{1,2}$ are determined, these two modes can be used jointly to constrain the two WCs. We illustrate this by three benchmark choices for the ratio: $r_{12} \equiv d_1/d_2 = 1, 0, -1$ with $d_1 = 0.115 \text{ GeV}^2$. In Fig. 3, we present the constraints in the $C_{\nu d u d}^{LR,x} - C_{\nu d u d}^{LL,x}$ plane derived from current experimental bounds, assuming both WCs are real. Whereas the $p \rightarrow \bar{\nu}_x \pi^+$ mode restricts the parameter space to a narrow band (hatched gray region), the $p \rightarrow \bar{\nu}_x \rho^+$ mode confines it to an elliptical region when $r_{12} \neq -1$. Combining the two constraints reduces the allowed region to a small area, emphasizing the complementary nature of these decay modes. This example highlights the crucial role played by vector meson modes in probing BNV interactions, and underscores the urgent need for a precise determination of these LECs.

As a second exercise, we neglect the unknown LECs d_i s and derive bounds on the vector meson modes by leveraging the experimental limits on the pseudoscalar modes. Assuming single operator dominance, the branching ratio of a vector meson mode is directly related to that of its corresponding pseudoscalar mode. In Table I, we present the derived bounds along with the associated bridging operators. For comparison, the second column lists the direct experimental lower bounds on their inverse decay widths [48]. As can be seen, for all kinematically allowed two-body vector meson modes (except $p \rightarrow e^+ \bar{K}^{*0}$ and $n \rightarrow e^+ K^{*-}, e^- \rho^+, \mu^- \rho^+$), the obtained indirect bounds are stronger than the direct experimental limits by 2–6 orders of magnitude. In particular, for the processes $p \rightarrow \nu(\bar{\nu}) K^{*+}$ and $n \rightarrow \nu(\bar{\nu}) K^{*0}$, the much stronger bounds associated with the respective operators $\mathcal{O}_{\nu dsu}^{\text{LR/LL},x}(\mathcal{O}_{\nu dsu}^{\text{RL/RR},x})$ and $\mathcal{O}_{\nu uds}^{\text{LR},x}(\mathcal{O}_{\nu uds}^{\text{RL},x})$ are due to cancellations between the Σ^0 - and Λ^0 -mediated contributions. Last but not least, for the experimentally unconstrained mode $n \rightarrow e^- K^{*+}$, we have derived a stringent upper bound based on that of the pseudoscalar mode $n \rightarrow e^- K^+$.

Summary. In this work, we developed a consistent chiral framework for nucleon decays involving vector mesons. Starting from the general triple-quark operator structures in the LEFT [49], we performed their LO chiral realizations, treating both vector-meson and baryon octets as matter fields with well-defined chiral transformation properties. The relevant low-energy constants were estimated using NDA. To highlight the phenomenological significance of vector meson modes, we demonstrated through examples

that including these modes helps resolve degeneracies in the parameter space that cannot be constrained solely by pseudoscalar modes. Moreover, by employing the stringent constraints on the dim-6 WCs derived from the pseudoscalar modes, we established tighter constraints on vector meson modes than those currently available from experiments, when considering only the long-distance contributions.

Should a BNV signal be observed, the presence of these additional vector meson decay channels would provide valuable information for identifying the underlying interaction structures responsible for the signal. For specific new physics scenarios, our formalism serves as a bridge connecting different decay modes and predicts correlations among them that can be tested in experimental searches. In addition, our formalism is not limited to the usual nucleon decays, but also covers scenarios involving light particles in the final state when the spurion fields are accordingly chosen, as well as BNV decays of the tau lepton and heavier hyperons. A comprehensive study of all relevant decay channels thus enables us to probe a broader parameter space and gain deeper insights into fundamental questions such as baryogenesis, neutrino masses, and the nature of dark matter.

Acknowledgments. This work was supported by the Grants No. NSFC-12305110, No. NSFC-12247151, and No. NSFC-12035008.

Data availability. The data are not publicly available. The data are available from the authors upon reasonable request.

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