Neutron-antineutron oscillations from lattice QCD

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Fundamental symmetry tests of baryon number violation in low-energy experiments can probe beyond the Standard Model (BSM) explanations of the matter-antimatter asymmetry of the universe. Neutron-antineutron oscillations are predicted to be a signature of many baryogenesis mechanisms involving low-scale baryon number violation. This work presents a pioneering first-principles calculation of neutron-antineutron matrix elements of the six-quark operators needed to accurately connect measurements of the neutron-antineutron oscillation rate to constraints on $\Delta B = 2$ baryon number violation in BSM theories. Our calculation is performed with state-of-the-art chirally symmetric QCD discretization with physical quark masses and robust control over the majority of lattice systematic uncertainties, including nonperturbative renormalization of the six-quark operators to the MS scheme. We find that our results based on realistic QCD predict at least an order of magnitude more events in neutron-antineutron oscillation experiments compared to previous estimates based on the MIT bag model for the same BSM parameters. Phenomenological implications are highlighted by comparing expected bounds from proposed neutron-antineutron oscillation experiments to predictions of a specific model of post-sphaleron baryogenesis, thus demonstrating the impact future experiments can make in searches for new physics. Lattice artifacts and other systematic uncertainties that are not controlled in this pioneering calculation are not expected to significantly change this conclusion.

Introduction — Beyond the Standard Model (BSM) violation of baryon number conservation is necessary to explain the observed matter-antimatter asymmetry of the universe. Baryogenesis explanations involving physics at high scales, such as leptogenesis, are appealing but difficult to test in low-energy experiments. Alternative explanations such as post-sphaleron baryogenesis involve low-scale baryon number violation that can be directly constrained by current and near-future experiments. Extracting robust BSM theory constraints from these experiments is critical for fundamental symmetry tests of baryon and lepton number violation addressing the long-standing mystery of matter-antimatter asymmetry.

Neutron-antineutron oscillations $(n-\bar{n})$ are predicted as a signature of low-scale baryogenesis in BSM theories including $SO(10)$ grand unified theories (GUTs), left-right symmetric theories, $R$-parity-violating supersymmetry, low-scale quantum gravity, extra-dimensional theories, and string theories with exotic instantons; see Refs. [1–3] for comprehensive reviews and further references. Experimental constraints on $(n-\bar{n})$ come from large underground detection experiments such as Super-K [4] and SNO [5] or from cold neutron time-of-flight experiments [6]. Results are presented as bounds on the neutron-antineutron oscillation time $\tau_{n, \bar{n}}$ governing the time-dependent probability $P_{n, \bar{n}} = \sin^2(t/\tau_{n, \bar{n}})$ for a free neutron in vacuum to turn into an antineutron [3].

The best direct bound on $\tau_{n, \bar{n}}$ is from the cold neutron experiment at Institut Laue-Langevin (ILL), $\tau_{n, \bar{n}} > 0.89 \times 10^8$ s [6], which is essentially background-free and can be improved with larger neutron flux, magnetic shielding, and the latest technologies in neutron optics. There has been a proposal for a new cold neutron experiment at the European Spallation Source (ESS) that would be $\approx 10^3$ times more sensitive than the ILL experiment and improve cold neutron constraints on $\tau_{n, \bar{n}}$ by a factor of $32$ [7, 8], as well as experiments at other reactors [9]. The best bound to date on the neutron-antineutron transition time within Oxygen-16 is from Super-K, $\tau_{O_{16}} > 1.9 \times 10^{12}$ years [4], which constrains the free oscillation time $\tau_{n, \bar{n}} > 2.7 \times 10^8$ s after taking into account nuclear structure effects [10]. In the future, we expect underground neutrino facilities like DUNE to be able to provide competitive bounds thanks to improved background-rejection techniques [11].

BSM theories cannot directly predict $\tau_{n, \bar{n}}$ without including the nonperturbative effects of quantum chromodynamics (QCD) responsible for binding quarks into hadrons. Since BSM and QCD effects are important at different scales, an effective field theory (EFT) description allows BSM and QCD effects to be separated in calculations of $\tau_{n, \bar{n}}$. At low energies, any BSM mechanism mediating $n-\bar{n}$ oscillations is described by a particular linear combination of six-quark operators violating baryon number by two units. Their matrix elements between neutron and antineutron states are determined by QCD at the nonperturbative hadronic scale, and have to be calculated before any BSM predictions for $\tau_{n, \bar{n}}$ can be made. Prior to this work, MIT bag model estimates pro-
vided the state-of-the-art calculation of these six-quark matrix elements [12, 13]. However, MIT bag model estimates contain unknown systematic uncertainties that prevent them from robustly connecting \( \tau_{n,\pi} \) measurements to constraints on the scale of new physics \( \Lambda_{\text{BSM}} \).

This work presents a lattice QCD (LQCD) calculation of the six-quark matrix elements needed to connect measurements of \( \tau_{n,\pi} \) to BSM theory constraints that improves upon previous preliminary LQCD results [14, 15] by accounting for operator renormalization, bias correction, excited-state effects, and quark-mass dependence of the matrix elements. Domain Wall fermions (DWF) with physical pion masses are used to compute lattice-regularized matrix elements that are nonperturbatively renormalized and converted to the \( \overline{\text{MS}} \) scheme at a hadronic scale. In combination with perturbative renormalization for these operators [16, 17], our work connects BSM theory predictions to experimental measurements of \( \tau_{n,\pi} \) with all QCD uncertainties either quantified or, in the case of non-zero lattice spacing and finite-volume effects, parametrically small and improvable. Moreover, our results are already used in generic EFTs for baryogenesis such as the recent work in Ref. [18]. Future lattice calculations with more computational resources can further improve our results by reducing statistical uncertainties and systematic uncertainties associated with discretization and finite-volume effects, but are very unlikely to change the dramatic impact that QCD effects have on these six-quark matrix elements as shown by this work. Our results have already been presented [19] at international workshops [20] dedicated to devise the best physics plan to observe neutron-antineutron oscillations at the ESS, future reactors and future underground laboratories. The impact of our LQCD calculation on possible discovery of new physics at those facilities is the topic of this letter.

Neutron-antineutron operators — The basis for the lowest-dimension operators for \( n-\pi \) transitions that are color singlets and electromagnetically neutral was constructed in Refs. [12, 13, 16, 21, 22]

\[
\begin{align*}
\mathcal{O}^{1}_{\chi_1 \chi_2 \chi_3} &= (u^T C P_{\chi_1} u) (d^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{SSS}, \\
\mathcal{O}^{2}_{\chi_1 \chi_2 \chi_3} &= (u^T C P_{\chi_1} d) (u^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{SSS}, \\
\mathcal{O}^{3}_{\chi_1 \chi_2 \chi_3} &= (u^T C P_{\chi_1} d) (u^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{AAS},
\end{align*}
\]

where quark spin indices are implicitly contracted in the parentheses and quark color indices are implicitly contracted with the tensors

\[
\begin{align*}
\mathcal{T}^{SSS}_{\{ij\}\{kl\}\{mn\}} &= \varepsilon_{ikm} \varepsilon_{jln} + \varepsilon_{jkm} \varepsilon_{ijn} + \varepsilon_{ilm} \varepsilon_{jkn} + \varepsilon_{ikn} \varepsilon_{jlm}, \\
\mathcal{T}^{AAS}_{\{ij\}\{kl\}\{mn\}} &= \varepsilon_{ijm} \varepsilon_{kln} + \varepsilon_{ijm} \varepsilon_{klm},
\end{align*}
\]

\( P_{L,R} = \frac{1}{2} (1 \mp i \gamma_5) \) are chiral projectors, and the Euclidean charge conjugation matrix \( C \) satisfies \( C \gamma_\mu C^\dagger = -\gamma_\mu \).

Symmetries together with the Fierz relation \( \mathcal{O}^{1}_{\chi_1 \chi_2 \chi_3} = 3 \mathcal{O}^{3}_{\chi_1 \chi_2 \chi_3} \) reduce the number of independent operators to 14 (7+7 related by parity). In this work, we use basis operators from chiral isospin multiplets that renormalize multiplicatively [15, 17] and have particular implications for phenomenology. An extended discussion of the operators and their chiral properties is reported in Ref. [23]. There are three SM gauge-singlet operators\(^1\) that provide the dominant contributions to \( n-\pi \) transitions in SM EFT,

\[
\begin{align*}
Q_1 &= -4(u^T C P_{\chi_1} u) (d^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{AAS}, \\
Q_2 &= -4(u^T C P_{\chi_1} d) (u^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{AAS}, \\
Q_3 &= -4(u^T C P_{\chi_1} d) (u^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{AAS}.
\end{align*}
\]

The fourth \( SU(2)_L \times U(1) \) singlet operator

\[
\begin{align*}
Q_4 &= -\frac{4}{5} (u^T C P_{\chi_1} u) (d^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{SSS} - \frac{16}{5} (u^T C P_{\chi_1} d) (u^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{SSS}
\end{align*}
\]

has vanishing matrix elements in the \( SU(2) \) isospin limit and is not studied in this work. Isospin-violating corrections provide a subdominant systematic uncertainty and are also neglected. \( SU(2)_L \)-non-singlet operators can also lead to \( n-\pi \) oscillations,

\[
\begin{align*}
Q_5 &= (u^T C P_{\chi_1} u) (d^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{SSS}, \\
Q_6 &= -4(u^T C P_{\chi_1} d) (u^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{SSS}, \\
Q_7 &= -\frac{4}{3} (u^T C P_{\chi_1} u) (d^T C P_{\chi_2} d) (d^T C P_{\chi_3} d) \mathcal{T}^{SSS},
\end{align*}
\]

These operators transform in the same chiral irreducible representation but describe different chiral multiplet components and do not mix under renormalization [17, 23]. Isospin \( SU(2) \) symmetry leads to the following relation between matrix elements

\[
\langle \pi | Q_5 | n \rangle = \langle \pi | Q_6 | n \rangle = -\frac{3}{2} \langle \pi | Q_7 | n \rangle.
\]

The complete chiral basis of QCD and QED singlet operators is given by \( Q_I, Q_I^F, I = 1 \ldots 7 \) where parity-transformed operators are \( Q_I^F = (-Q_I)|_{L \leftrightarrow R} \). The operators (2-4) are related to the ones in Eq. (1) as

\[
\begin{align*}
Q_1 &= -4O^{3}_{RRR}, \\
Q_2 &= -4O^{1}_{LLR}, \\
Q_3 &= -4O^{1}_{LRR}, \\
Q_4 &= -\frac{4}{5} O^{1}_{LRR} - \frac{16}{5} O^{2}_{RRR}, \\
Q_5 &= O^{1}_{RLL}, \\
Q_6 &= -4O^{2}_{RLL}, \\
Q_7 &= -\frac{4}{3} O^{1}_{LLR} - \frac{8}{3} O^{1}_{LRR}.
\end{align*}
\]

\(^1\) Two additional operators arise in dimensional regularization at two-loop order [17] that are exactly equal to \( Q_{1,4} \) in lattice regularization by Fierz identities. These Fierz identities are preserved in \( \overline{\text{MS}} \) if one-loop matching is consistently included with two-loop running. At high scales \( \Lambda_{\text{BSM}} \gg \Lambda_{\text{QCD}} \) Fierz identity violations can be neglected even if BSM matching is performed at tree level.
Because of symmetries and Eq. (5), only four separate nucleon matrix elements \langle \bar{n} Q_1, 2, 3, 5 | n \rangle need to be determined using lattice QCD methods.

Lattice QCD results — For this calculation, we use an ensemble of QCD gauge field configurations on a 48^3 \times 96 lattice generated with Iwasaki gauge action and \( N_f = 2 + 1 \) flavors of dynamical M"obius Domain Wall fermions with masses almost exactly at the physical point [24]. The pion mass is \( m_{\pi} = 139.2(4) \text{ MeV} \) and the lattice spacing is \( a = 0.1141(3) \text{ fm} \). With the physical lattice size \( L \approx 5.45 \text{ fm} \), and \( m_{\pi} L = 3.86 \), finite volume effects on the matrix elements are estimated from chiral perturbation theory to be \( \lesssim 1\% \) [25].

We calculate lattice correlation functions on 30 independent gauge field configurations separated by 40 MD steps using point-source quark propagators aided by AMA sampling [26] to reduce stochastic uncertainty.

On each configuration, we compute 1 exact and 81 low-precision samples evenly distributed over the 4D volume. For the latter, quark propagators are computed with low-mode deflation and 250 iterations of the conjugate gradient algorithm. The propagators are contracted into intermediate baryon blocks [27, 28] representing (anti)neutron source or sink operators made of point and Gaussian-smeared quarks and denoted by \( n^{P,S} \), respectively. These blocks are finally contracted into (anti)neutron two-point correlation functions of \( P \) source and \( J = P, S \) sink operators,

\[
G_{2pt}^{P,S}(t) = \sum_{x} \langle \langle n_t \rangle_{\alpha}^{P,S}(x, t) (\pi_{\tau})_{\alpha}^{P,S}(0) \rangle,
\]

as well as three-point correlation functions involving \( J = S, P \) antineutron sources, neutron sinks, and six-antiquark operators \( Q_1 \) that are obtained from \( Q_1 \) by charge conjugation and have identical matrix elements,

\[
G_{3pt}^{J,J'}(\tau, t; Q_1) = \sum_{x, y} \langle \langle n_t \rangle_{\alpha}^{J,J'}(x, t - \tau) (\bar{Q}_1)(0) (n_\tau)_{\alpha}^{J,J'}(y, -\tau) \rangle
\]

where \( |m, \uparrow\rangle \rangle \) denote the spin-up (anti)neutron states, \( (\pi, t) \rangle \) are the Euclidean time intervals from the source to the operator and the sink, respectively.

To extract the \( n-\pi \) matrix elements, we perform correlated \( \chi^2 \)-minimization 2-state fits. First, we fit the \( G_{2pt}^{P,P} \) and \( G_{2pt}^{P,S} \) data to Eq. (7) to find optimal values for energies \( E_{0,1} \) and obtain the corresponding optimal \( \sqrt{G_{0,1}^{P,S}} \) using VarPro [29]. This is done for a variety of fit ranges \( t^f \in [t^f_{\min}, t^f_{\max}] \), where \( J = P, S \) corresponds to the sink smearing, to assess the contamination from excited states. The energy parameters \( E_{0,1} \) entering Eq. (8) are then fixed and the \( n-\pi \) ground-state \( (n, m = 0) \) matrix elements are obtained from correlated linear fits of \( G_{3pt}^{P,S,S} \) data to Eq. (8). An extensive and detailed description of the analysis strategy for two and three-point functions is reported in our companion paper [23] which includes results and quality of fits for different \( t^f_{\min}, t^f_{\max} \) ranges, together with a quantification of the systematic uncertainties. The data and fit results for the six-quark operator \( Q_2 \) are shown in Fig. 2, where ratios of \( G_{2pt}^{J,J'} \) and \( G_{3pt}^{J,J'} \) are plotted. Those ratios are expected to reach a plateau when the ground state saturates the correlation functions and excited states contaminations have disappeared. The gray bands in the plots represent the matrix element for the operator \( Q_2 \) from a 2-state fit, including the statistical uncertainty.

The operators \( Q_1^{P,S} (2-4) \) do not mix in perturbation theory because of chiral symmetry constraints [17]. Non-
perturbative effects, quark mass effects, and residual chiral symmetry breaking could lead to operator mixing and must be investigated in nonperturbative renormalization studies. We compute the renormalization of these operators on a lattice nonperturbatively using the RI-MOM scheme [30] in the Landau gauge. RI-MOM renormalization factors $Z_{IJ}^{\text{RI-MOM}}$, $Q_{IJ}^{\text{RI-MOM}} = Z_{IJ}^{\text{RI-MOM}} Q_{IJ}^{\text{bare}}$ at momentum $p$ are defined as

$$
(Z_q^{\text{RI-MOM}}(p))^{-3} Z_{IJ}^{\text{RI-MOM}}(p) \Lambda_{JK}(p) = \delta_{IK},
$$

(9)

where $Z_q^{\text{RI-MOM}}$ is the lattice quark field renormalization and $\Lambda_{JK}(p)$ are amputated Green's functions of the lattice operators $Q_J$ and quark fields carrying momenta $\pm p$ projected onto the spin-color-flavor structure of $Q_K$. A more detailed discussion of operator mixing and the numerical procedure used to calculate the renormalization factors can be found in the companion paper [23]. We find that the renormalization matrices $\Lambda_{JK}$ and therefore $Z_{IJ}^{\text{RI-MOM}}$ are diagonal in our chiral basis (2–4) up to $O(10^{-3})$ [23], as expected from the good chiral symmetry properties of the lattice implementation we adopt. We neglect this residual mixing and identify $Z_I = Z_{II}$ below. Chiral symmetry is crucial to obtaining precise

results for these matrix elements and we demonstrate this for the first time.

The "scale-independent" combination of perturbative and nonperturbative factors

$$
Z_{I}^{\text{SI}}(\mu_0, p) = Z_I^{\text{RI-MOM}}(p) \left[ \frac{Z_I^{\text{RI-MOM}}(\mu_0)}{Z_I^{\text{RI-MOM}}(|p|)} \right]^{\text{pert}},
$$

(10)

where the factor in brackets is calculated perturbatively, has residual dependence on the lattice momentum $p$, due to discretization effects, rotational symmetry breaking, and higher orders in perturbation theory. We analyze the lattice artifacts following Ref. [31], with a representative fit for operator $Q_2$ and $\mu_0 = 2$ GeV shown in Fig. 3. Similar plots for all operators and a summary of all fits and their corresponding uncertainties are included in our companion paper [23]. We have found no substantial difference between fits using 1-loop and 2-loop perturbative factors in Eq. (10). Different fits including several intervals of $p \in [1.6, 4.5]$ GeV are used to define a central value for $Z_{I}^{\text{RI-MOM}}$ with statistical and systematic uncertainties using a "weighted" average as described in the analysis of Ref. [23]. These statistical and systematic uncertainties are added in quadrature to those of the bare matrix elements to estimate the total RI-MOM matrix element uncertainties. RI-MOM matrix element results are then converted to the $\overline{\text{MS}}(2$ GeV$)$ and $\overline{\text{MS}}(700$ TeV$)$ scheme using 1-loop matching [17],

$$
\langle \pi | Q_I^{\overline{\text{MS}}} | n \rangle = \left[ \frac{Z_{IJ}^{\overline{\text{MS}}}}{Z_{IJ}^{\text{RI-MOM}}} \right]_{N_f = 3}^{N_f = 4} Z_{I|N_f = 3}^{\text{RI-MOM}} (\pi|Q_I|n),
$$

(11)

where perturbative running to/from the charm quark threshold $\mu = M_c$ is used to convert $N_f = 3$ lattice results to $\overline{\text{MS}}$ with $N_f = 4$ and $\alpha_s$ is taken from the four-loop determination of Ref. [32]. Higher-order matching uncertainties are estimated to be $\lesssim 7\%$ based on the size of
1-loop matching effects [17] and are neglected. Final results for the ≈ matrix elements are shown in Tab. I with their total uncertainties and comparisons to old MIT bag model estimates.

Phenomenological implications — In BSM theories where \( \Delta B = 2 \) transitions are permitted, experimentally observable \( n - \pi \) oscillations are low-energy phenomena that can be described in an EFT containing only SM fields. The low-energy EFT will include \( \Delta B = 2 \) terms involving the operators \( Q_I \) above,

\[
\mathcal{L}_{n-\pi} = \sum_{I=1}^{7} \left( C_I(\mu)Q_I(\mu) + C_I^{(P)}(\mu)Q_I^{(P)}(\mu) \right),
\]

where the \( C_I \) are numerical coefficients with mass dimension \( -5 \) that are predicted to be non-zero in some BSM theories. The \( SU(2)_L \)-singlet operators are EW-symmetric and their coefficients should scale as \( C_{1,2,3,4} \sim \Lambda_{\text{BSM}}^{-5} \) in naive dimensional analysis. In contrast, the \( SU(2)_L \)-non-singlet operators \( Q_{6,7}^{(P)} \) and \( Q_{5,6,7} \) can only appear in an SM gauge-invariant Lagrangian in products with additional SM Higgs (or BSM) fields to make them \( SU(2)_L \)-singlets. Assuming the former, their coefficients should scale as \( C_{6,7}^{(P)}(\Lambda_{\text{BSM}}) \sim v^2 \Lambda_{\text{BSM}}^{-7} \) and \( C_{5,6,7} \sim v^4 \Lambda_{\text{BSM}}^{-9} \), where \( v \) is the vacuum expectation value of the Higgs field. For \( v \ll \Lambda_{\text{BSM}} \), this provides a significant additional suppression on \( n - \pi \) oscillation rate contributions from \( SU(2)_L \)-non-singlet operators.

The \( n - \pi \) oscillation rate is given by the matrix element of the associated Hamiltonian between neutron and antineutron states, which in the isospin limit of QCD simplifies to

\[
\tau_{n-\pi}^{-1} = \left| \sum_{I=1,2,3,5} \tilde{C}_I(\mu) \langle \bar{n} | Q_I(\mu) | n \rangle \right|.
\]

where \( \tilde{C}_I = C_I - C_I^{(P)} \) for \( I = 1, \ldots, 4 \) and \( \tilde{C}_5 = (C_5 - C_5^{(P)} + (C_6 - C_6^{(P)})/3) \). Contributions involving \( C_4^{(P)} \) vanish exactly in the isospin limit considered here, although in principle isospin-violating \( C_4^{(P)} \) contributions could play a role in particular BSM models. Using the LQCD results above, the \( n - \pi \) oscillation rate is given by

\[
\tau_{n-\pi}^{-1} = \left( \frac{10^{-9} \text{ s}^{-1}}{(700 \text{ TeV})^{-5}} \right) \left| 4.2(1.1)\tilde{C}_1(\mu) - 8.6(1.5)\tilde{C}_2(\mu) + 4.5(1.1)\tilde{C}_3(\mu) + 0.096(43)\tilde{C}_5(\mu) \right|_{\mu=2 \text{ GeV}}.
\]

Predictions for non-zero \( \tau_{n-\pi}^{-1} \) arise in some BSM theories explaining the matter-antimatter asymmetry of the universe and other outstanding problems of the SM and cosmology. An example is provided by \( n - \pi \) oscillations in left-right symmetric gauge theories where the SM gauge group is embedded in \( SU(2)_L \times SU(2)_R \times SU(4)_C \) with \( (B - L) \) acting as a fourth color [33-35]. Post-sphaleron baryogenesis occurs after a colored BSM scalar field develops a \( (B + L) \)-breaking vacuum expectation value that leads to Majorana neutrino masses and \( n - \pi \) oscillations. The \( n - \pi \) oscillation rate in this model only involves \( Q_1 \) at tree-level and \( \tau_{n-\pi}^{-1} \) is given by the first term in Eq. (14).

Conclusions — An LQCD calculation of six-quark matrix elements is presented that provides, for the first time, renormalized \( n - \pi \) transition matrix elements in the \( \overline{MS} \) scheme at 2 GeV. These renormalized results with well-defined scale dependence are required to reliably connect experimental measurements of \( \tau_{n-\pi} \) to the baryon number-violating new physics scale \( \Lambda_{\text{BSM}} \).

This calculation is performed with physical pion masses, a chirally-symmetric fermion discretization, and a large spacetime volume. Ground-state matrix elements are extracted using 2-state fits and systematic uncertainties associated with excited state effects are estimated through variation of the fitting region. Finite-volume effects are predicted by chiral EFT to be \( \lesssim 1\% \) for the \( L \approx 5.45 \text{ fm} \) volume used in this study. Even though we only use one lattice spacing \( a \approx 0.114 \text{ fm} \), discretization effects are expected to be small because of automatic \( O(a) \) improvement due to the chiral symmetry of the DWF action. Additional systematic effects due to the uncertainty in the lattice scale \( \delta a/a \approx 1.3\% \) are negligible. The error budget of our final results is dominated by limited statistics. Statistical uncertainties and systematic uncertainties associated with discretization and finite-volume effects will be improved in future calculations using additional gauge field ensembles, e.g., \( a \approx 0.084 \text{ fm} \) ensembles of Ref. [24].

To conclude, there has been recent phenomenological interest in \( n - \pi \) oscillations and the possibility of new
searches for $n$-$\pi$ vacuum oscillations and transitions in nuclei at ESS, DUNE, and other experiments [7–9, 11]. The magnitudes of electroweak-singlet $n$-$\pi$ transition matrix elements are 4-8 times larger than those computed in the MIT bag model [12]. Experiments should consequently observe 16-64 times more neutron-antineutron oscillation events for fixed BSM physics parameters than was previously estimated using the bag model results. Future searches for $n$-$\pi$ oscillations at large-mass detectors (HyperK, DUNE) and neutron sources (WWR-M, ESS) will be able to probe the parameter space of several viable baryogenesis scenarios [18]. Our results, despite being pioneering, indicate that experimental searches of $n$-$\pi$ transitions are about 1 to 2 orders of magnitude more sensitive to baryon number violating interactions in BSM physics than previously expected and will be able to put more stringent constraints on various baryogenesis mechanisms.

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