

Displaced vertex signals of low temperature baryogenesis

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We explore the connection of baryogenesis at temperatures below the electroweak scale and signals for long-lived particles at the LHC. The model features new SM singlets, with a long-lived fermion decaying to quarks to generate the baryon asymmetry. The model avoids strong flavor physics bounds while predicting a rich diquark phenomenology, monojet signals, and displaced vertices. We show how the displaced vertex signals can be probed at the high luminosity LHC. The large transverse production makes a strong physics case for constructing far detector experiments such as MATHUSLA, ANUBIS, and CODEX-b, complementary to the central and forward long-lived particle program.

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I. INTRODUCTION

The observed matter antimatter asymmetry remains one of the biggest motivations for physics beyond the standard model (BSM). Some of the traditional BSM baryogenesis mechanisms such as electroweak baryogenesis [1,2], leptogenesis [3], and Affleck-Dine models [4] generate the baryon asymmetry at high temperatures, above the electroweak (EW) scale, and sometimes much above, e.g., $\sim 10^{15}$ GeV. While these are well-motivated models, they can be elusive for testing given the high new physics scales involved. Alternatively, low-temperature baryogenesis models propose to generate the baryon asymmetry after the EW transition [5–11]. Baryogenesis occurs through a new particle decaying after the sphaleron decoupling and before big bang nucleosynthesis (BBN) times. These decays violate baryon number and generate a CP asymmetry, meeting all three Sakharov conditions [12].

If a particle N decays to generate the asymmetry, its decay length has to satisfy the out-of-equilibrium condition, $\tau_N > H^{-1}$ at temperatures $T \simeq m_N$. For masses around the EW scale, this relation implies

$$c\tau_N > 20 \text{ mm} (m_N/100 \text{ GeV})^{-2}, \quad (1)$$

resulting in a decay length of macroscopic size. Furthermore, if the lifetime of N is longer than a picosecond, its decays occur after the sphaleron transition. If it is shorter than a second the decay occurs before BBN. This connection between long-lived particles (LLP) and low-temperature baryogenesis represents an excellent opportunity for displaced vertex searches precisely because they select the decay length window of $0.1 \text{ mm} \lesssim c\tau_N \lesssim 200 \text{ m}$. Such signals can be most effectively searched at the Large Hadron Collider (LHC) and dedicated far detector experiments such as MATHUSLA [13,14], CODEX-b [15,16], ANUBIS [17], AL3X [18], and the Forward Physics Facilities [19] with FASER/FASER2. The relationship between LLP and baryogenesis models is surprisingly underexplored in the literature, with only a few proposed models. As reviewed in [14], some of these models are weakly interacting massive particles baryogenesis [20–23] and baryogenesis via exotic baryon oscillations [24–26]. In these, charged BSM states enter on shell in the decay loop amplitude so as to generate CP violation. Because extra SM-charged particles should be heavy, such models prefer masses above the EW scale. Instead, we focus on a lower mass range, $10 \text{ GeV} \lesssim m_N \lesssim 200 \text{ GeV}$, resulting in generally longer lifetimes. To reach these masses, we propose a model with only SM singlets below the TeV scale, one of which can successfully decay to create a baryon asymmetry. We combine several features of different models proposed in the literature [27–35]. The resulting simplified framework naturally avoids proton decay, neutron oscillations, neutron electric dipole moment (EDM), and strong flavor physics bounds while providing a diverse phenomenology for colliders and cosmology. In a companion paper [36], we expand the model to include spontaneous symmetry breaking of baryon number. Here we focus on the model's aspects that are important for displaced vertex signals.

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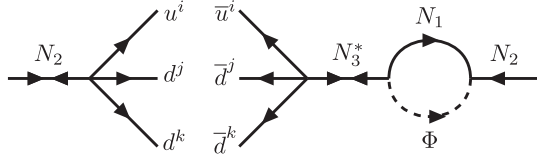


FIG. 1. Diagrams for CP violation in the interference between tree-level and loop decays of N_2 with and without mixing with the heavier state N_3 . The lightest singlet fermion N_1 and scalar Φ are on shell in the loop to generate a CP asymmetry.

II. MODEL

We consider three flavors of singlet Majorana fermions, $N_\alpha = N_{1,2,3}$, one neutral scalar, Φ , along with new dynamics at a high UV scale. Baryon number violation fixes the dynamics at the UV scale, which generates an effective coupling between N and a neutral combination of quarks. The simplest operator that couples N to the SM baryon number happens at dimension six. Since N is neutral, the only quark flavor structure allowed for pairing is udd' . A minimal possibility that meets these requirements is

$$\mathcal{L}_{\text{eff}} = \frac{\kappa_\alpha^{ijk}}{M_X^2} (\bar{N}_\alpha^c u_R^i) (\bar{d}_R^c d_R^k) + \xi_{\alpha\beta} \bar{N}_\alpha^c \Phi N_\beta + \text{H.c.} \quad (2)$$

Here, M_X is the UV scale; i, j, k are quark generation indices; and α, β are the neutral fermion flavors. We assume a mass hierarchy of $m_{N_1} \lesssim m_\Phi < m_{N_2} < m_{N_3}$. Then, CP violation occurs by the interference between decays with and without mixing of one of the neutral fermions and another heavier state as shown in Fig. 1. A nonzero CP phase requires the mixing term to have on shell intermediate states in the loop. Therefore, the intermediate states must be lighter than the decaying particle. These requirements show that CP violation occurs exclusively in the decay of N_2 as it mixes at loop level with the heavier N_3 and has an on shell loop contribution from the lighter N_1 .

The baryon asymmetry parameter is the product of the yield of N_2 , $Y_2 = \frac{n_{N_2}}{s}$, the CP asymmetry ϵ_{CP} and the branching ratio of N_2 to quarks, $\text{Br}(udd')$,

$$Y_{\Delta B} = Y_2 \epsilon_{CP} \text{Br}(udd'). \quad (3)$$

Using the diagrams of Fig. 1, the CP asymmetry is computed to be

$$\epsilon_{CP} \simeq \frac{3}{8\pi} \frac{m_{N_1}}{m_{N_3}} \frac{|\kappa_2 \xi_{12}^* \xi_{13} \kappa_3^*| \sin \delta}{|\kappa_2|^2} \sqrt{1 - \frac{m_{N_1}^2 + m_\delta^2}{m_{N_2}^2}}, \quad (4)$$

where δ is the phase resulting from the coupling product $\kappa_2^{ijk} \xi_{12}^* \xi_{13} \kappa_3^{ijk*}$. We summed over the quark final states c, d, s, b assuming the flavor hierarchies described in the Appendix, which allows us to write $\kappa_\alpha^{cjk} \approx \kappa_\alpha$ in (4).

To compute N_2 's yield, we consider the thermal history of the neutral sector. The relevant processes are the annihilation of N_2 to quarks, $N_2 u \leftrightarrow dd'$, and the decay of $N_2 \rightarrow udd'$. For $m_{N_2} \sim 100$ GeV with small nonrenormalizable interactions, the freeze-out of the annihilation processes can occur when N_2 is still relativistic, with freeze-out temperature $m_{N_2} < T_{\text{FO}} < M_X$. We can estimate the relativistic freeze-out temperature by setting $n\langle\sigma_{\text{ann}}v\rangle \simeq H(T_{\text{FO}})$ during radiation domination to get

$$T_{\text{FO}} \simeq 280 \text{ GeV} \left(\frac{M_X}{2 \text{ TeV}} \right)^{4/3} \left(\frac{10^{-6}}{\kappa_2} \right)^{2/3}. \quad (5)$$

The relativistic freeze-out of N_2 is welcome since it maximizes the baryon asymmetry by not having a Boltzmann suppressed population of N_2 . With a sufficiently large lifetime, N_2 can decay after freeze-out and also after the sphaleron decoupling. We can obtain the N_2 yield using the relativistic equilibrium expression

$$Y_2 = \frac{45\zeta(3)}{2\pi^4} \frac{g_{N_2}}{g_{*,s}(T_{\text{FO}})}, \quad (6)$$

where g_{N_2} is the number of N_2 degrees of freedom, and $g_*(T_{\text{FO}})$ is the total number of degrees of freedom in the bath at the freeze-out temperature.

Lastly, to compute the branching ratio, we must consider the relative contributions of the tree-body and two-body decays, $N_2 \rightarrow udd'$ and $N_2 \rightarrow N_1 \Phi$. For $m_{N_1} \sim m_\Phi$ the partial decay widths are given by

$$\Gamma_{N_2 \rightarrow udd'} \simeq \frac{3|\kappa_2|^2}{192\pi^3} \frac{m_{N_2}^5}{M_X^4}, \quad \Gamma_{N_2 \rightarrow N_1 \Phi} \simeq \frac{m_{N_2} |\xi_{12}|^2}{\pi}. \quad (7)$$

The decay rate of $N_2 \rightarrow udd'$ should be comparable to $N_2 \rightarrow N_1 \Phi$ in order to generate the baryon asymmetry efficiently. In turn, this requirement translates into a relationship between the couplings κ_2 and ξ_{12} ,

$$\frac{|\xi_{12}|}{|\kappa_2|} \approx \frac{1}{8\sqrt{3}\pi} \left(\frac{m_{N_2}}{M_X} \right)^2, \quad (8)$$

where the equality corresponds to a 50% branching ratio to the udd' final state. Putting together (4) and (6) and requiring (8) to get a $\mathcal{O}(1)$ branching ratio to the baryon number violating channel, the baryon asymmetry for $M_X = 2 \text{ TeV}$ is given by

$$\frac{Y_{\Delta B}}{Y_{\Delta B}^{\text{exp}}} \simeq \left(\frac{m_{N_2}}{100 \text{ GeV}} \right)^2 \frac{\kappa_3 \xi_{13} \sin \delta m_{N_1}}{3 \times 10^{-2} m_{N_3}} \sqrt{1 - \frac{m_{N_1}^2 + m_\delta^2}{m_{N_2}^2}}, \quad (9)$$

where we used the central value $Y_{\Delta B}^{\text{exp}} = 8.7 \times 10^{-11}$ measured by Planck [37]. In Fig. 2, we impose the measured baryon asymmetry to fix the masses of the $N_{1,2,3}$.

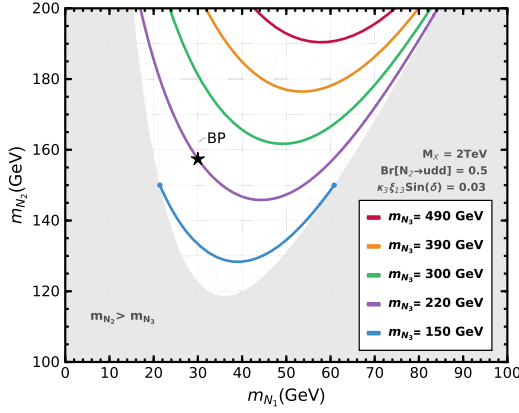


FIG. 2. Masses of the particles in the model that reproduce the observed baryon asymmetry. We plot the curves respecting relation (9) in the $m_{N_1} \times m_{N_2}$ plane. Each curve represents a choice of mass m_{N_3} for the scale $M_X = 2$ TeV. We choose the values $\kappa_3 \xi_{13} \sin \delta = 0.03$. We assume that N_2 decays to u, d, s, c, b quarks, with a branching ratio of 0.5. BP indicates the benchmark point chosen throughout the paper.

It is useful to define a simple tree-level UV completion for (2) by integrating in a $(\mathbf{3}, \mathbf{1})_{2/3}$ SM charged scalar diquark, X . The allowed gauge invariant terms are

$$\mathcal{L}_X = \lambda_{ai} X^\dagger \bar{N}_a^c u_R^i + \lambda'_{jk} X \bar{d}_R^c d_R^k + \text{H.c.} \quad (10)$$

Such diquark has some interesting low-energy properties. First, if $m_{N_a} > 1$ GeV and N does not mix with neutrinos, the proton does not decay [29]. Second, QCD gauge invariance implies color antisymmetry, which means the flavor antisymmetry of the $\bar{d}_R^c d_R^k$ quark couplings in (10). Therefore, there are only three independent λ' couplings, $\lambda'_{jk} = \epsilon_{jkl} \lambda'_l$. Because of this, there is no tree-level neutron oscillation. The lightest baryon that oscillates is Λ_0 , which imposes weak constraints on the diquark mass [36]. Additionally, there is no tree-level $K-\bar{K}$ and $B-\bar{B}$ mixing. At one loop, neutral Kaon mixing must involve the b quark. In the case of B-meson mixing, the loop must contain an s quark. Then, if one of the couplings is small, e.g., $\lambda'_{bs} < \lambda'_{ds}, \lambda'_{db}$, the bounds from meson oscillations can be negligible while still allowing for order 1 diquark couplings. Neutron-antineutron oscillations are suppressed by two-loop and Cabibbo-Kobayashi-Maskawa contributions rendering negligible bounds for this model. Regarding neutron EDM bounds, the new CP -violating phases only contribute at three loops or higher. A detailed discussion of low-energy bounds is given in the Appendix.

III. LLPS AT THE LHC

Having defined the UV model, we now discuss the LHC sensitivity for the predicted LLPs. We start by considering the direct production bounds on the diquark that govern the necessary UV dynamics of the model. CMS [38,39] and

ATLAS [40] conducted searches for nonresonant pair production of dijet resonances. Their benchmark model is the R-parity-violating supersymmetric top squark that decays to ds quarks and can be directly associated with the diquark X above. The leading limits from [38] excludes the $Y = 2/3$ scalar diquark at 95% confidence level between 0.50 TeV and 0.77 TeV, with a 3.6(2.5) local(global) excess occurring at 0.95 TeV. There are also resonant searches for X to a pair of jets conducted by CMS [41,42] and ATLAS [43] through their diquark couplings. Reference [44] imposed diquark bounds from resonant production for different values of the coupling pair $(\lambda'_{ds}, \lambda'_{sb})$. For a diquark of mass $M_X = 2$ TeV, the values below $(\lambda'_{ds}, \lambda'_{sb}) = (0.3, 0.12)$ are allowed by both direct searches and the flavor constraints of meson oscillations.

The production of N_a leads to monojet signals when there is a single X or multijet plus missing p_T for diquark pair production. However, pair production is subdominant for $M_X = 2$ TeV and couplings of $\lambda_{ai} \lesssim 0.3$. Therefore we consider only the monojet channels in our analysis. Moreover, for the monojet channels, N_a can be singly produced in the $ds \rightarrow Nu_i$ channel or doubly produced by the gluon initiated $gu \rightarrow X^{(*)}N \rightarrow NN'u_i$ and diquark t -channel $u\bar{u}(g) \rightarrow NN'(g)$ processes. In the case of the double production t channel, the monojet comes from the initial state emission of a gluon, but these are limited by the high p_T cut required by the analysis. For the other two channels, the resulting p_T distribution of the jets is a Jacobian peak at $M_X/2$. Because of this feature, the cross section can be sizable as it does not suffer from the high p_T cuts. CMS [45] and ATLAS [46] have searched for monojets in the context of fermion portal dark matter [47] and light nonthermal dark matter models [48]. In our model, these bounds are most relevant for N_3 , which is produced and decays promptly to the LLPs. The monojet N_3 channel is the largest since κ_3 has to be maximized in (3) to get a sizable $Y_{\Delta B}$. Because of this, the monojet N_3 channel is the one driving our LLP phenomenology, as we discuss next.

First we focus on the benchmark point shown in the plots with a star labeled BP. We hold the diquark couplings fixed to $(\lambda'_{ds}, \lambda'_{db}, \lambda'_{sb}) = (0.25, 0.28, 0.1)$. We choose $\lambda_{au} \ll \lambda_{ac} = \lambda_{at}$ to respect the bounds described in the Appendix and set $\lambda_{1c} = \lambda_{2c} = 5 \times 10^{-4}$ and $\lambda_{3c} = 0.1$, which gives decay lengths $c\tau_1 = 10$ m and $c\tau_2 = 1$ cm for N_1 and N_2 , while N_3 decays promptly. The production channel considered is $pp \rightarrow jN_3$, with N_3 decaying to either three N_1 s or one N_2 and two N_1 s. This happens through the $N_3 \rightarrow N_{1,2}\Phi$ and $\Phi \rightarrow N_1 N_1$ cascade decay, where we assume the Φ Yukawa couplings to be $(\xi_{13}, \xi_{23}) = (1, 0.3)$. For the signal simulation, we use MadGraph5_aMC@NLO version 2.9.19 with the parton distribution function NNPDF2.3QED and generate the UFO model using FeynRules [49].

We extract the angular and velocity distribution of N_1 and N_2 from the simulated data. The majority of events are

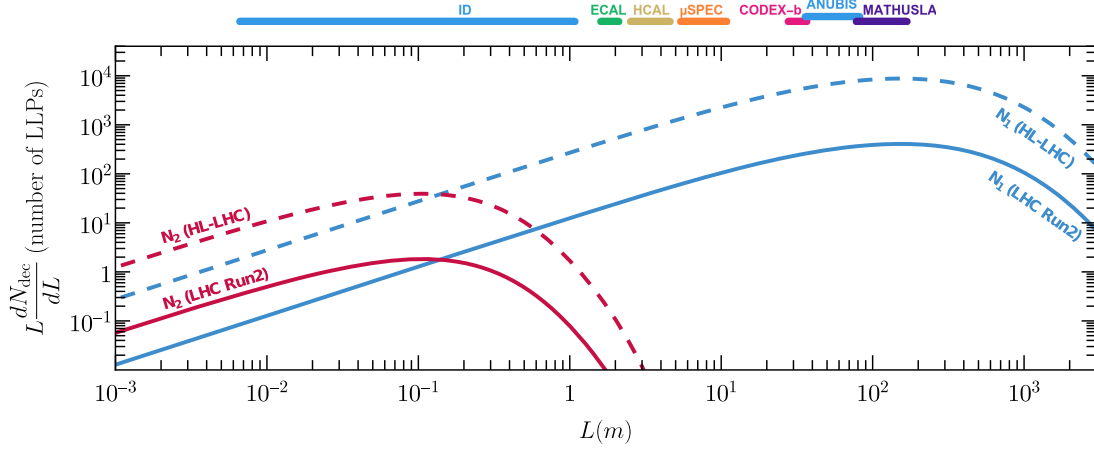


FIG. 3. Differential distribution for the number of LLP decays as a function of the distance from the interaction point L . The red curve is the number of N_2 with proper decay length $c\tau_2 \simeq 1$ cm, and the blue is for N_1 with $c\tau_1 \simeq 10$ m. The solid lines are for LHC run 2 with $\mathcal{L} = 139 \text{ fb}^{-1}$, and the dashed are for the high-luminosity phase (HL-LHC) with $\mathcal{L} = 3000 \text{ fb}^{-1}$. The color lines above the plot indicate the ATLAS inner detector (ID); the electromagnetic (ECAL) and hadronic calorimeters (HCAL); the muon spectrometer (μ SPEC); and the transverse far experiments CODEX-b, ANUBIS, and MATHUSLA.

transverse to the interaction plane, forming an approximately uniform angular distribution for $|\eta| < 5$. From the velocity distribution, we extract the boosted lifetimes $\beta\gamma c\tau$, and construct the differential probability distribution for $N_{1,2}$ to decay at position L ,

$$\frac{d\mathcal{P}_{\text{dec}}(L)}{dL} = \frac{e^{-\frac{L}{\beta\gamma c\tau_i}}}{\beta\gamma c\tau_i}. \quad (11)$$

Then, the differential number of events observed as a function of the distance L is given by

$$\frac{dN_{\text{obs}}}{dL} = \overbrace{(n\sigma\mathcal{L}) \otimes \text{Br}(udd')}^{\equiv dN_{\text{dec}}/dL} \otimes \frac{d\mathcal{P}_{\text{dec}}}{dL} \otimes \epsilon_{\text{geom}}^{\text{LLP}} \otimes \epsilon_{\text{recon}}^{\text{LLP}}, \quad (12)$$

where n is the multiplicity of $N_{1,2}$ of the process, $\mathcal{L}$ is the luminosity, $\epsilon_{\text{geom}}^{\text{LLP}}$ is the geometric acceptance of the experiment, and $\epsilon_{\text{recon}}^{\text{LLP}}$ is the efficiency of reconstruction of the displaced vertices. In Fig. 3, we obtain the differential number of LLP decays defined in the overbrace of (12) as a

function of the distance L . Notice that the value obtained is for the whole solid angle without specifying the geometric acceptance and detector efficiencies. The number of displaced vertices is sizable in several regions of the future LHC experiments, as well as for proposed far detectors with sensitivity to transverse events.

To get an estimate of the number of LLPs in each experiment, we can integrate the length L in (11) to obtain the decay probability inside a detector that starts at L_{in} and ends at L_{fin} . In Table I, we estimate the number of events for the benchmark point used throughout the paper for various LHC experiments during run 2 and the HL-LHC. For decays inside the inner detector of the ATLAS experiment, we assume the large-radius tracking reconstruction algorithm used in a similar search for RPV SUSY (supersymmetry) long-lived neutralinos [50]. Since the ATLAS RPV SUSY search targets heavy neutralinos, we only count N_2 decays inside the inner detector. A dedicated study of the three displaced vertices' signal from N_3 decays into light N_1 plus a monojet could improve sensitivity but is left for future work. There are also displaced decays inside the calorimeters and muon spectrometers of the ATLAS/CMS experiments. However, we are not aware of a reconstruction

TABLE I. Number of observable displaced vertices from the decay of $N_{1,2} \rightarrow udd'$ inside the ID, ECAL and HCAL, and μ SPEC. We assume the large-tracking radius algorithm of [50] to estimate the efficiency of LLP reconstruction for decays inside the ID and optimistically choose $\epsilon_{\text{recon}}^{\text{LLP}} = 1$ for far detectors. The number of events inside the ECAL/HCAL and μ SPEC is left as a function of reconstruction efficiency.

	ATLAS/CMS				Transverse far detectors			Forward far detectors
	ID (N_2)	ECAL	HCAL	μ SPEC	CODEX-b	ANUBIS	MATHUSLA	FASER/FASER2, ...
N_{obs} (Run2)	<3	$18 \times \epsilon_{\text{recon}}^{\text{LLP}}$	$116 \times \epsilon_{\text{recon}}^{\text{LLP}}$	$163 \times \epsilon_{\text{recon}}^{\text{LLP}}$	...	...	...	<1
N_{obs} (HL)	60	$397 \times \epsilon_{\text{recon}}^{\text{LLP}}$	$2509 \times \epsilon_{\text{recon}}^{\text{LLP}}$	$3537 \times \epsilon_{\text{recon}}^{\text{LLP}}$	26	72	370	<1

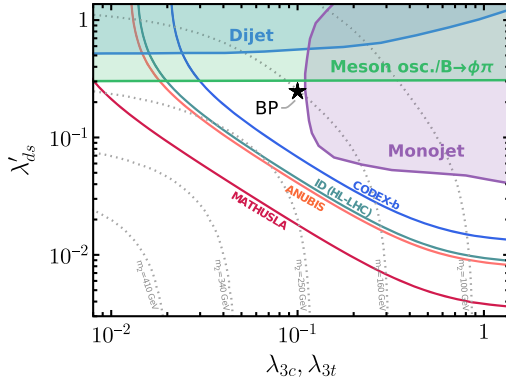


FIG. 4. Exclusion bounds on $(\lambda_{3c} = \lambda_{3t}, \lambda'_{ds})$. The purple region is excluded by CMS monojet searches [45], blue by ATLAS + CMS dijet searches, and the green by meson mixing constraints [44,51]. Dotted lines show the m_{N_2} values required for the asymmetry (9), assuming $m_{N_3} = m_{N_2} + 70$ GeV.

strategy for the three-jet signals of the LLPs of our model beyond the ID. Because of that, we show the resulting number of events without specifying the reconstruction efficiencies of each part of the detector. For the far detectors, we consider the transverse and forward experiments CODEX-b, ANUBIS, MATHUSLA, and FASER/FASER2. To obtain the geometric acceptances, we integrate the number of events distribution over the geometric coverage of each experiment integrating from L_{in} to L_{fin} for each detector. We choose the optimistic scenario by assuming that the reconstruction efficiencies of the far detectors are $\epsilon_{recon}^{LLP} = 1$. While the forward far detectors have negligible event rates due to the limited solid angle coverage, the transverse detectors present a significant sensitivity to LLP decays. This enhanced sensitivity is primarily due to their optimal positioning relative to the interaction point, which allows for a larger geometric acceptance and increased event visibility. As shown in Fig. 3, the experiments considered are complementary in their range for probing $c\tau$. CODEX-b, ANUBIS, and MATHUSLA are strategically positioned at successive distances from the interaction point, allowing for a continuous probing of the region ranging from 26 to 170 meters. Lastly, in Fig. 4, we show the projected sensitivity of various HL-LHC experiments in the $(\lambda_{c3}, \lambda'_{ds})$ plane, assuming $\lambda_{c3} = \lambda_{t3}$ and $\lambda_{u3} \ll \lambda_{c3}$. In some regions, the displaced vertex searches extend the reach for this baryogenesis model by 1 order of magnitude compared to current flavor and monojet bounds. This highlights the potential of the HL-LHC to probe the low scale baryogenesis and motivates the construction of the proposed far detectors.

IV. CONCLUSIONS

In this paper, we explore the connection of baryogenesis through out-of-equilibrium decays at low temperatures and the detection of long-lived particles at the LHC. The proposed model can be probed through the detection of

LLPs and associated phenomenology using current and future LHC experiments. The model requires the presence of three flavors of majorana fermions $N_{1,2,3}$, a neutral scalar Φ , and some TeV scale dynamics which we assume to be related to a scalar diquark with $2/3$ hypercharge. The model avoids proton decay because N does not mix with neutrinos, has no sizable neutron oscillations, and has small loop-suppressed $\Delta F = 2$ neutral meson oscillations. At the same time, there are plenty of interesting collider signals ranging from the detection of the diquarks to monojet signals with missing energy and displaced vertices, which are the main focus of the paper. The HL-LHC provides a promising scenario for probing the model as a significant number of LLPs can be produced. Transverse far detectors like CODEX-b, ANUBIS, and MATHUSLA have a significantly higher sensitivity to LLP decays compared to forward detectors due to their optimal positioning and larger geometric acceptance. The increased sensitivity of transverse far detectors highlights the importance of their construction and complementary ability to probe new physics at the HL-LHC.

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APPENDIX: LOW ENERGY BOUNDS

This appendix is dedicated to detailing the low energy bounds from the UV completed model with a scalar diquark $X = (3, 1)_{2/3}$. In the end, we summarize the required couplings for the model to be compatible with current constraints. The Lagrangian of the model is

$$\begin{aligned} \mathcal{L} = & |\partial_\mu \Phi|^2 - V(\Phi) + \frac{1}{2} \bar{N}_\alpha (\not{\partial} - m_{N_\alpha}) N_\alpha - \xi_{\alpha\beta} \bar{N}_\alpha^c \Phi N_\beta \\ & + |D_\mu X|^2 - M_X^2 |X|^2 - \lambda_{ai} X^\dagger \bar{N}_a^c u_R^i - \lambda'_{jk} X \bar{d}_R^j d_R^k + \text{H.c.} \end{aligned} \quad (\text{A1})$$

As discussed in the main text, the QCD structure, together with the fermionic nature of the spinors, imply that the couplings λ'_{jk} are antisymmetric in flavor. This means that there are three independent flavor combinations, $\lambda'_{ds}, \lambda'_{db}$ and λ'_{sb} . When rotating to the mass basis, the coupling transforms as $\lambda' \rightarrow \tilde{\lambda}' = U_{d_R}^T \lambda' U_{d_R}$. This transformation preserves the antisymmetry:

$$\tilde{\lambda}'^T = (U_{d_R}^T \lambda' U_{d_R})^T = U_{d_R}^T \lambda'^T U_{d_R} = -U_{d_R}^T \lambda' U_{d_R} = -\tilde{\lambda}'. \quad (\text{A2})$$

Therefore, the coupling remains antisymmetric in the mass basis. This property holds as long as the coupling involves only down-type (or only up-type) quarks since the corresponding fields are rotated independently.

After integrating out the heavy diquark, we can write the following effective interactions:

$$\begin{aligned} \mathcal{L}_{\text{eff}}^{d=6} = & \frac{\lambda_{ai}\lambda'_{jk}}{M_X^2} (\bar{N}_\alpha^c u_R^i) (\bar{d}_R^j d_R^k) + \frac{\lambda'_{jk}\lambda'_{lm}}{M_X^2} (\bar{d}_R^j d_R^k) (\bar{d}_R^l d_R^m) \\ & + \frac{\lambda_{ai}\lambda_{\beta n}}{M_X^2} (\bar{N}_\alpha^c u_R^i) (\bar{u}_R^n N_\beta) + \text{H.c.} \end{aligned} \quad (\text{A3})$$

Apart from generating the baryogenesis mechanism described in the main text, these effective interactions lead to a distinct low-energy phenomenology. As we discussed, provided that $m_N \geq m_p + m_K \approx 1.4$ GeV and the operator $\bar{L}HN$ is forbidden in the theory, the proton does not decay. Next, because the diquarks do not directly couple to two quarks of the same generation, tree-level $\Delta F = 2$ flavor-changing neutral current processes are avoided. For the same reason, tree-level neutron-antineutron oscillations are also forbidden in the model. At tree level, two main observables put bounds on the couplings of the model. First, tree-level flavor-changing decays can mimic SM decays induced by penguin diagrams. The most constraining ones are B meson decays like $B \rightarrow \phi\pi$ and $B \rightarrow \phi\phi$. Following the calculation done in [51], current limits on the branching ratio of these decays put the following constraints on the diquark couplings:

$$\begin{aligned} \frac{\text{Br}(B^\pm \rightarrow \phi\pi^\pm)}{1.5 \times 10^{-7}} = & \left(\frac{|\lambda'_{ds}\lambda'_{sb}|}{0.035} \right)^2 \left(\frac{2 \text{ TeV}}{M_X} \right)^4 < 1, \\ & \text{at 90\% C.L. [52, 53]} \end{aligned} \quad (\text{A4})$$

$$\begin{aligned} \frac{\text{Br}(B^0 \rightarrow \phi\pi^0)}{1.5 \times 10^{-7}} = & \left(\frac{|\lambda'_{ds}\lambda'_{sb}|}{0.049} \right)^2 \left(\frac{2 \text{ TeV}}{M_X} \right)^4 < 1, \\ & \text{at 90\% C.L. [53, 54].} \end{aligned} \quad (\text{A5})$$

Another tree-level process is the $\Delta B = 2$ and $\Delta F = 2$ dinucleon decay $pp \rightarrow K^+ K^+$ induced by the first operator of (A3). These are constrained by the SuperKamiokande experiment through reactions of the type $O^{16} \rightarrow C^{14} K^+ K^+$ [52,53]. Following [54], this process put the following limit:

$$\begin{aligned} |\lambda_{au}\lambda'_{ds}| \left(\frac{2 \text{ TeV}}{M_X} \right)^2 \left(\frac{200 \text{ GeV}}{m_{N_a}} \right)^2 & < 1.5 \times 10^{-6}, \\ \alpha = 1, 2, 3. \end{aligned} \quad (\text{A6})$$

Since we assume the couplings λ'_{ds} to be sizable for the production of the new states at the LHC, the dinucleon decay put strong bounds on the up-quark flavor couplings of the diquarks and N_α . We therefore assume a hierarchy in which $\lambda_{au} \ll \lambda_{ac}, \lambda_{at}$.

At the one-loop level, the most important bounds come from meson oscillations. References [44,51,55–57] imposed the following constraints from Kaon and B meson oscillations:

$$\sqrt{|\text{Re}(\lambda_{db}^* \lambda'_{sb})^2|} \left(\frac{2 \text{ TeV}}{M_X} \right) \leq 9.2 \times 10^{-2}, \quad (K^0 - \bar{K}^0) \quad (\text{A7})$$

$$\sqrt{|\text{Re}(\lambda_{ds}^* \lambda'_{sb})^2|} \left(\frac{2 \text{ TeV}}{M_X} \right) \leq 7.2 \times 10^{-2}, \quad (B_d^0 - \bar{B}_d^0). \quad (\text{A8})$$

Similarly, owing to the λ_{ai} couplings, there can also be D^0 meson oscillations. The effective Hamiltonian for the transition is given by

$$\mathcal{H}_{\text{eff}} = \sum_{\alpha\beta} \frac{\lambda_{au}\lambda_{\beta u}\lambda_{ac}^*\lambda_{\beta c}}{16\pi^2 M_X^2} f\left(\frac{m_{N_{\alpha\beta}}^2}{m_X^2}\right) (\bar{u}P_R c)(\bar{u}P_R c). \quad (\text{A9})$$

Since the diquark couplings with $N_{1,2}$ must be small so that they are long lived, the dominant contribution comes from the λ_{3u} and λ_{3c} couplings. With only one singlet in the loop the function $f(x)$ is given by

$$\begin{aligned} f(x) = & \frac{2(1+x^2-2x+x\log x-x^2\log x)}{(1-x)^3} \Big|_{x \ll 1} 2 \\ & + \mathcal{O}(x[1+\ln(x)]). \end{aligned} \quad (\text{A10})$$

Then, following [58], the off-diagonal D^0 mass term is

$$M_{12} = \frac{1}{2m_D} \langle D^0 | \mathcal{H}_{\text{eff}} | \bar{D}^0 \rangle = \frac{5}{12} \frac{m_D f_D^2 B_D}{2} \frac{|\text{Re}(\lambda_{3u}^* \lambda_{3c})^2|}{48\pi^2 M_X^2}, \quad (\text{A11})$$

where $m_D \simeq 1.9$ GeV, $f_D \simeq 0.21$ GeV are the D -meson mass and decay constant and $B_D \approx 0.8$ is the bag factor. Requiring that the mass difference $\Delta m_D = 2M_{12}$ satisfies the experimental bound from [59], $\Delta m_D^{\text{exp}} \lesssim 0.94 \times 10^{10} \hbar s^{-1} = 6.2 \times 10^{-15}$ GeV, leads to

$$\sqrt{|\text{Re}(\lambda_{3u}^* \lambda_{3c})^2|} \left(\frac{2 \text{ TeV}}{M_X} \right) \leq 2.9 \times 10^{-2}, \quad (D^0 - \bar{D}^0). \quad (\text{A12})$$

Other one-loop bounds from $\Delta F = 1$ processes like $b \rightarrow d\gamma, s\gamma$ and contributions to the chromomagnetic sd operator are weaker than the meson oscillation ones we quote above. For a detailed discussion, we refer to [51].

There are no one-loop contributions for neutron-antineutron oscillations. For $n - \bar{n}$, the absence of one-loop effects can be seen by factorizing the loop diagram

into two tree-level ones ($udd \rightarrow \psi_1 \psi_2$) $\times$ ($\bar{\psi}_1 \bar{\psi}_2 \rightarrow \bar{u} \bar{d}$) using the optical theorem. The intermediate two-particle state $\psi_1 \psi_2$ must be neutral and carry $\Delta B = 2$. This means we need a single Majorana mass insertion, which forces $\psi_1 = N_\alpha$. However, the remaining state ψ_2 must also be a fermionic singlet, implying that $\psi_2 = N_\beta$. This leaves no room for a nonvanishing spinor contraction, as one would need two mass insertions to obtain the correct fermionic lines. For the neutron EDM, the one-loop diagrams always involve the modulus squared of a single coupling, producing no imaginary part. Therefore, there is no one-loop neutron EDM.

At two loops, neutron oscillation is possible but very suppressed. The operator that generates $n - \bar{n}$ is

$$\mathcal{O}_{n-\bar{n}} = \frac{1}{\Lambda_{n\bar{n}}^5} (u_R d_R d_R)^2. \quad (\text{A13})$$

At two loops, this operator is generated dominantly by the diquark sector as EW flavor-changing processes cannot couple to right-handed currents without light quark mass insertions. One can estimate the effective scale $\Lambda_{n\bar{n}}^5$ by power counting the loop contributions of the type shown in Fig. 5:

$$\frac{1}{\Lambda_{n\bar{n}}^5} \sim \sum_\alpha \sum_{i,j,k,l} \frac{|\lambda_{au}|^2 \lambda'_{dk} \lambda'_{ik} \lambda'_{dl} \lambda'_{jk} m_{N_\alpha}}{256\pi^4 M_X^6} \sim \frac{|\lambda_{3u}|^2 |\lambda'_{ds}|^4 m_{N_3}}{256\pi^4 M_X^6}. \quad (\text{A14})$$

Then the bound for $n - \bar{n}$ oscillations can be obtained as

$$\tau_{n-\bar{n}} = \frac{\Lambda_{n\bar{n}}^5}{\Lambda_{\text{QCD}}^6} \sim 8.6 \times 10^7 \text{ s} \times \left(\frac{(0.09)^6}{|\lambda_{3u}|^2 |\lambda'_{ds}|^4} \right) \left(\frac{200 \text{ GeV}}{m_{N_3}} \right) \left(\frac{M_X}{2 \text{ TeV}} \right)^6. \quad (\text{A15})$$

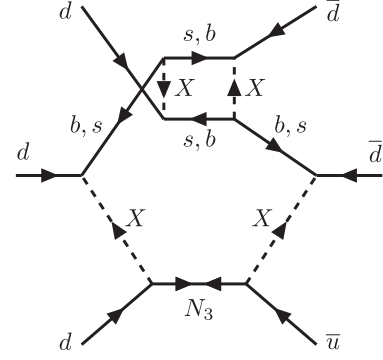


FIG. 5. Two-loop contribution to neutron-antineutron oscillations induced by the diquark couplings of the model.

Notice that because the coupling λ_{3u} is suppressed due to the $pp \rightarrow K^+ K^+$ bound, neutron oscillations are well within the allowed values of $\tau_{n-\bar{n}}^{\text{exp}} \approx 8.6 \times 10^7 \text{ s}$ [59].

Finally, the need for at least two different diquark couplings and a chirality flip suggests that a nonzero neutron EDM contribution only appears at higher loops. In fact, since the only chirality flip in the model occurs through Higgs insertions, there are also no anomalous neutron EDM corrections at two loops. In Ref. [51], it is suggested that the diquark couplings can generate a neutron EDM at three loops, but an estimate of the bounds must be subleading given the other constraints on the couplings and is beyond the scope of this work.

In conclusion, to avoid all low energy bounds it is sufficient to adopt the following parameters:

$$\lambda'_{ds} < 0.3, \quad \lambda'_{db} < 0.3, \quad \lambda'_{sb} < 0.11, \quad (\text{A16})$$

$$\lambda_{au} < 1.5 \times 10^{-6}, \quad \lambda_{ac} < 1, \quad \lambda_{at} < 1, \quad \alpha = 1, 2, 3. \quad (\text{A17})$$

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