Testing leptogenesis at the LHC and future muon colliders: a $Z'$ scenario

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If the masses of at least two generations of right-handed neutrinos (RHNs) are near-degenerate, the scale of leptogenesis can be as low as $\sim 100$ GeV. In this work, we study probing such resonant leptogenesis in the $B - L$ model at the LHC and future multi-TeV muon colliders via the process $Z' \rightarrow NN \rightarrow \ell^±\ell^± +$jets, with $Z'$ the $U(1)_{B-L}$ gauge boson and $N$ the RHN. The same-sign dilepton feature of the signal makes it almost background-free, while the event number difference between positive and negative leptons is a hint for $CP$ violation, which is a key ingredient of leptogenesis. We found that resonant leptogenesis can be tested at the HL-LHC for $M_{Z'}$ up to 12 TeV, while at a 10 (30) TeV muon collider the reach can be up to $M_{Z'} \sim 30$ (100) TeV via the off-shell production of $Z'$.

I. INTRODUCTION

The baryon asymmetry of the Universe (BAU) is one of the most mysterious unsolved problems in the Standard Model (SM) of particle physics. Leptogenesis is a very attractive explanation [1–3], as it links the BAU to the origin of neutrino masses. In that mechanism, the $CP$ violating decay of the heavy right-handed neutrinos (RHN) to the SM leptons generates a lepton asymmetry, which is then converted to the BAU via the electroweak (EW) sphaleron process. The same interaction also accounts for the tiny neutrino masses [4–6] via the Type-I seesaw mechanism [7]. In the conventional thermal leptogenesis formalism, the $CP$ violating effects are related to the RHN mass $M_N$, and the explanation of the BAU requires $M_N \gtrsim 10^9$ GeV [8], making it very challenging to test the mechanism experimentally. However, the constraints on $M_N$ can be relaxed if at least two of the RHNs are highly degenerate, so that the $CP$ asymmetry is resonantly enhanced, and hence even $O$(TeV) RHNs can generate the observed BAU [9–13]. In particular, if such a resonant leptogenesis mechanism is embedded into a gauge theory in which the RHNs exist naturally for anomaly cancellation, then the $O$(TeV) leptogenesis is testable at the colliders via searches for RHNs or new gauge/scalar bosons as well as the lepton charge asymmetries [14–21].

In this article, we perform a comprehensive study of the collider phenomenology of the resonant leptogenesis mechanism in the gauged $U(1)_{B-L}$ extended SM (i.e. the so-called $B - L$ model [23–26]), focusing on the interplay between leptogenesis and the $Z'$ gauge boson of the $U(1)_{B-L}$ group, since the $Z'$-mediated scattering process greatly impacts the generated BAU. In particular, we use the $Z' \rightarrow NN \rightarrow \ell^±\ell^± +$jets channel to probe the $Z'$, $N$ particles as well as the $CP$ violation. Compared with previous studies on similar topics [14–21], our work includes not only the newest constraints from the LHC and the corresponding projected reach at the future HL-LHC, but also the first projections at the future multi-TeV muon colliders. The study on the physics potential of muon colliders started around three decades ago [27, 28], and receives a renewed interest recently [29–60]. Due to the high energy and precision measurement environment, the future multi-TeV muon colliders offer us the opportunity to probe both SM and beyond the SM physics very accurately.

This paper is organized as follows. The $B - L$ model and the corresponding (resonant) leptogenesis is described in Section II, where the relation between BAU and the size of $CP$ asymmetry is derived. Then we study the collider phenomenology in Section III, including the reach for $Z'$ boson and the $CP$ asymmetry (via RHNs) at the HL-LHC and muon colliders. Finally, the conclusion is given in Section IV.

II. RESONANT LEPTOGENESIS IN THE $B - L$ MODEL

A. The model

In the $B - L$ model, three generations of RHNs (with $B - L = -1$) are needed naturally for gauge anomaly cancellation. Besides, the model also contains an extra gauge boson $Z'$ and a complex scalar $\Phi = (\phi + i\eta)/\sqrt{2}$ with $B - L = 2$ that breaks the $U(1)_{B-L}$ spontaneously. The relevant Lagrangian reads

$$\mathcal{L}_{B-L} = \sum_i \bar{\nu}_i D_\mu \nu_i^\dagger \Phi - \frac{1}{2} \sum_{i,j} \left( \lambda_{ij}^D \bar{\nu}_i^\dagger \Phi \nu_j^R + \text{h.c.} \right) - \sum_{i,j} \left( \lambda_{ij}^D \bar{\nu}_i^\dagger \tilde{H} \nu_j^R + \text{h.c.} \right)$$

$$+ D_\mu \Phi^\dagger D^\mu \Phi - \lambda_\Phi \left( |\Phi|^2 - \frac{v_\Phi^2}{2} \right)^2 - \frac{1}{4} Z'_{\mu\nu} Z'^{\mu\nu},$$

with $i$, $j$ being the family indices, $D_\mu = \partial_\mu - ig_{B-L} X Z'_\mu$ the covariant derivative ($X$ is the $B - L$ quantum

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* Also at Ref. [22] for a recent review.
number, $\ell_L$ the SM left-handed lepton doublet and $H$ the SM Higgs doublet. Without loss of generality, we assume $\lambda^N_{ij} = \text{diag}\{\lambda_{N_1}, \lambda_{N_2}, \lambda_{N_3}\}$, and define the four-component Majorana RHNs as $N_i = \nu^c_R + (\nu^c_R)^c$. The minimum of the scalar potential is at $\langle \Phi \rangle = v_\phi/\sqrt{2}$, breaking the $U(1)_{B-L}$ symmetry and providing masses for the particles

$$M_{Z'} = 2g_{B-L}v_\phi, \quad M_{N_i} = \lambda_{N_i}v_\phi/\sqrt{2}, \quad M_\phi = \sqrt{2\lambda_{N_i}v_\phi},$$

(2)

and the imaginary part $\eta$ of $\Phi$ is absorbed to be the longitudinal mode of $Z'$. We are interested in the parameter region $v_\phi \sim \mathcal{O}(\text{TeV})$. The usage of Yukawa interactions between $\ell$ and $N$ is two-fold. On one hand, they account for the tiny left-handed neutrino mass via the Type-I seesaw

$$m_{\nu_L} \sim \frac{\lambda_{N_0}^2 v^2}{M_N} \sim 0.06 \text{ eV} \times \left(\frac{\lambda_D}{10^{-6}}\right)^2 \left(\frac{1 \text{ TeV}}{M_N}\right),$$

(3)

while on the other hand they trigger the RHN decay $N \rightarrow \ell H/\ell H^*$, which is the crucial process in thermal leptogenesis. Due to the CP violation phase in the Yukawa couplings, the widths of $N \rightarrow \ell H$ and $N \rightarrow \ell H^*$ are different, and hence a CP asymmetry can be defined as [9–12]

$$\epsilon_i = \frac{\sum_j \Gamma_{N_i \rightarrow \ell_j H} - \Gamma_{N_i \rightarrow \ell_j H^*}}{\sum_j \Gamma_{N_i \rightarrow \ell_j H} + \Gamma_{N_i \rightarrow \ell_j H^*}} \approx -\sum_{j \neq i} \frac{M_{N_i} \Gamma_{N_i \rightarrow \ell_j H}}{M_{N_j}^2} \left(\frac{V_{ij}}{2} + S_{ij}\right) \frac{\text{Im}(\lambda_D^0 \lambda_D^1)_{ij}}{(\lambda_D^0 \lambda_D^1)_{ii} (\lambda_D^0 \lambda_D^1)_{jj}},$$

(4)

where

$$V_{ij} = 2 \frac{M_{N_i}^2}{M_{N_i}^2} \left[1 + \frac{M_{N_j}^2}{M_{N_j}^2}\right] \ln \left(1 + \frac{M_{N_j}^2}{M_{N_i}^2}\right) - 1,$$

$$S_{ij} = \frac{M_{N_i}^2 (M_{N_j}^2 - M_{N_i}^2)}{(M_{N_i}^2 - M_{N_j}^2)^2 + M_{N_i}^2 \Gamma_{N_i}},$$

(5)

are respectively the vertex correction and RHN self-energy correction to the decay process, and the tree level width is

$$\Gamma_{N_i} = \frac{M_{N_i}}{8\pi} (\lambda_D^0 \lambda_D^1)_{jj}.$$  

(6)

If there is a mass hierarchy among the three RHNs, i.e. $M_{N_1} \ll M_{N_2} \ll M_{N_3}$, $V_{ij}$ and $S_{ij}$ are comparable, and $\epsilon_i$ is proportional to the lightest RHN mass but typically $\gtrsim 10^{-6}$. Therefore, a sizable BAU requires a RHN with mass $\gtrsim 10^9 \text{ GeV}$ [8]. However, if at least two RHNs are highly degenerate that $|M_{N_2}^2 - M_{N_3}^2| \sim M_{N_2} \Gamma_{N_2}$, then $S_{ij} \sim M_{N_2}/\Gamma_{N_2} \gg 1$, and $\epsilon_i$ can reach $\mathcal{O}(1)$ [9–13]. In this case, even $\mathcal{O}(\text{TeV})$ RHNs can generate a successful BAU. This is the scenario under consideration in this article.

### B. Thermal leptogenesis

We assume first two flavors of RHNs are near-degenerate, i.e. $(M_{N_2}^2 - M_{N_3}^2) \sim M_{N_2} \Gamma_{N_2}$, while the third flavor RHN is much heavier, i.e. $M_{N_3} \gg M_{N_2}$. The BAU receives contributions from the CP violating decays of both $M_{N_1}$ and $M_{N_2}$. If the decay widths $\Gamma_{N_1}$ and $\Gamma_{N_2}$ are comparable, then the generated BAU should be twice of that from mere $N_1$ decay. However, if the decay widths have a hierarchy, e.g. $\Gamma_{N_2} \ll \Gamma_{N_1}$, then so do the CP asymmetries, as $\epsilon_2 \sim 2\epsilon_1 \Gamma_{N_1}/\Gamma_{N_2} \ll \epsilon_1$, and in the meantime the washout from $N_2$ is negligible [12]. In that case, the generated BAU is dominated by $N_1$ decay, and a one-flavor discussion on $N_1$ is sufficient. For simplicity, we will consider such a scenario throughout this article.

We then denote $\epsilon_1$ as $\epsilon$, and $M_{N_1} \approx M_{N_2} = M_N$ from now on. In the radiation dominated era, the energy and entropy densities of the Universe are respectively

$$\rho = \frac{\pi^2}{30} g_* T^4, \quad s = \frac{2\pi^2}{45} g_* T^3,$$

(7)

where $g_* = 106.75$ is the number of relativistic degrees of freedom. The Hubble constant is then derived by the first Friedman equation $H^2 = (8\pi/3M_P^2)\rho$, with $M_{P1} = 1.22 \times 10^{19} \text{ GeV}$ the Planck scale. Defining the dimensionless parameter $z \equiv M_N/T$, the Boltzmann equations for the RHN and net $B - L$ number yield in the thermal bath read [12, 61, 62]

$$\frac{s_N H_N}{z^4} \frac{dY_N}{dz} = -\frac{Y_N^{\text{eq}}}{Y_N^{\text{eq}}} (\gamma_D + 2\gamma_{h,s} + 4\gamma_{h,t}) - \frac{Y_N^2}{(Y_N^{\text{eq}})^2} - 1 \frac{2\gamma_{Z'}}{\gamma_D - \frac{Y_{B-L}^{\text{eq}}}{Y_L}} \left[2(\gamma_{N,s} + \gamma_{N,t} + \gamma_{h,t}) + \frac{Y_N}{Y_N^{\text{eq}}} \gamma_{h,s}\right],$$

(8)

where $s_N$ and $H_N$ are the entropy density and Hubble constant at $z = 1$, respectively, and the abundances are defined as number density to entropy density ratios (e.g. $Y_N = n_N/s$), with

$$Y_N^{\text{eq}} = \frac{45 z^2}{2\pi^2} K_2(z), \quad Y_L^{\text{eq}} = \frac{3\zeta(3) M_N^3}{2\pi^2 s_N},$$

(9)

the equilibrium abundances, and $K_i(z)$ is the modified
Bessel function of the $i$-th kind. When writing Eq. (8), we have neglected the charge lepton flavor effects, which can affect the production and washout of the lepton asymmetry and the low energy experiments such as $\mu \rightarrow e\gamma$. See the review [63] for a fully flavor-covariant treatment on the resonant leptogenesis. In this article, since we are interested in the $\ell^+\ell^- + \text{jets}$ final state with $\ell = e, \mu$, the flavor effects from the first two generation of charged leptons are expected to be subdominant.

The reaction rates in Eq. (8) are defined as [61],

$$\gamma_D = \frac{M_N^3}{\pi^2} \Gamma_N \frac{K_1(z)}{z},$$

for the $N \rightarrow \ell H/\bar{\ell}H^*$ decay, and

$$\gamma_{ab\rightarrow cd} \equiv \langle \sigma_{ab\rightarrow cd} \rangle n_{eq} n_{b}^0,$$

for the $2 \rightarrow 2$ scattering, where $s_{\text{min}} = \max\{M_a + M_b, M_c + M_d\}$, and the dimensionless reduced scattering cross section is

$$\hat{\sigma}_{ab\rightarrow cd}(s) \equiv 2 \sigma_{ab\rightarrow cd}(s) \cdot s \cdot \left[ 1 - 2 \left( \frac{M_a^2}{s} + \frac{M_b^3}{s} \right) + \left( \frac{M_a^2}{s} - \frac{M_b^2}{s} \right)^2 \right].$$

The correspondence between the reaction rates in Eq. (8) and the $2 \rightarrow 2$ processes are listed in Table I, except the case of $\gamma_{Z'\gamma}$, which corresponds to the scattering $N N \rightarrow Z' \rightarrow f \bar{f}$ ($f$ denotes the SM fermions) and is highlighted below

$$\hat{\sigma}_{Z'}(s) = \frac{13g_{B-L}^4}{6\pi} \sqrt{x(x-4)} \left( \frac{x}{x - M_{Z'}^2/M_N^2} \right)^3, \Gamma_{Z'}/M_{Z'},$$

where

$$\Gamma_{Z'} = \frac{g_{B-L}^2}{24\pi} \left[ 13 + 2 \left( 1 - \frac{M_{Z'}^2}{M_N^2} \right)^{3/2} \theta(M_{Z'} - 2M_N) \right],$$

with $x = s/M_{Z'}^2$ and $\theta$ the Heaviside step function. We also assume $M_{Z'} > M_N$ so that scatterings $NN \rightarrow Z'Z'$ or $\phi \phi$ are suppressed.

For a $\mathcal{O}(\text{TeV})$ leptogenesis, $\lambda_D \sim 10^{-6}$, making the $\gamma$'s in Table I rather small, as they are proportional to $\lambda_D^2$. Only the reaction rate $\gamma_{Z'}$ can be sizable since it is $\propto g_{B-L}^2$. Therefore we can omit the reaction rates in Eq. (8) except $\gamma_D$ and $\gamma_{Z'}$, and simplify Eq. (8) into a single equation,

$$\frac{dY_{B-L}}{dz} + \left( \frac{z^4}{s_N H_N} \frac{\gamma_D}{2Y_{\ell}^{eq}} \right) \frac{Y_{B-L}}{Y_{B-L}^{eq}} = \epsilon \left( \frac{\gamma_D}{\gamma_D + 4\gamma_{Z'}} \right),$$

where we have approximated $Y_N/Y_N^{eq} + 1 \approx 2$ and $dY_N/dz \approx dY_{B-L}^{eq}/dz$, since $N$ is not far away from equilibrium. Eq. (15) can be solved analytically as

$$Y_{B-L}(z) \approx \epsilon \int_{z_{\text{in}}}^z dz' dY_{B-L}^{eq} (z) \left( \frac{\gamma_D}{\gamma_D + 4\gamma_{Z'}} \right) \times$$

$$\exp \left\{ - \int_{z_{\text{in}}}^z dz'' \frac{z''^4}{s_N H_N} \frac{\gamma_D}{2Y_{B-L}^{eq}} \right\},$$

where we adopt $z_{\text{in}} = 1$ as the lower limit of the integral. It is very clear in above equation that $\epsilon$ and $\gamma_D$ generate the lepton asymmetry, while $\gamma_{Z'}$ tends to washout this asymmetry because it tends to push $N$ back to the equilibrium. The final generated baryon asymmetry is

$$Y_B = \frac{28}{79} Y_{B-L}(z_{\text{ph}},)$$

where $z_{\text{ph}} = M_N/T_{\text{ph}}$, with $T_{\text{ph}} \approx 130 \text{ GeV}$ the decoupled temperature of the EW sphaleron [64]. We have checked that Eq. (16) matches the numerical solution of the complete equation set Eq. (8) very well.

Given the value of $g_{B-L}$, one is able to derive the $CP$ asymmetry $\epsilon$ as a function of $(M_{Z'}/M_N)$ via Eq. (16) by the observed BAU $Y_{B-L}^{\text{obs.}} \approx 10^{-10}$ [68, 69]. This is shown in the left panel of Fig. 1, where $g_{B-L} = 0.8$ is fixed. Near the line $M_{Z'} = 2M_N$, the washout process $NN \rightarrow Z' \rightarrow f \bar{f}$ is resonantly enhanced and hence a large $\epsilon$ is needed to realize $Y_{B-L}^{\text{obs.}}$. The region with $\epsilon > 1$ is

\begin{table}[h]
| Reaction rates | 2 $\rightarrow$ 2 scattering | $\hat{\sigma}(s)$ |
|----------------|-------------------------------|-----------------|
| $\gamma_{\ell,\bar{\ell}}$ | $N_{\ell,\bar{\ell}} \rightarrow \ell q$ | $\frac{3g^2 \lambda^2}{4\pi} \left( \frac{x}{x-1} \right)^2$ |
| $\gamma_{\ell,t}$ | $N_{\ell,t} \rightarrow t\ell$ or $N_{\ell,t} \rightarrow q\ell t$ | $\frac{3g^2 \lambda^2}{4\pi} \left( \frac{x}{x-1} \right)^2 + \frac{1}{2} \ln \frac{x-1+M_1^2/M_N^2}{M_1^2/M_N^2}$ |
| $\gamma_{N,s}$ | $tH \rightarrow \ell H^*$ | $\frac{\lambda^2}{2\pi} \left[ 1 + \frac{2}{2(\pi^2 s)} + \frac{s}{2D_1(s)} \right] - \frac{1}{2} \frac{\gamma_D(1+s)}{D_1(s)} \ln(1+x)$ |
| $\gamma_{N,t}$ | $t\ell \rightarrow H^* H^*$ | $\frac{\lambda^2}{2\pi} \left( \frac{x}{x+1} + \frac{1}{2} \right)$ |

TABLE I. The reaction rates and the reduced cross sections taken from Refs. [12, 61], where $x = s/M_N^2$ and $D_1(x) = x^{-1} + (\Gamma_N^2/M_N^2)/(x-1)$.

\footnote{The impact of $NN \rightarrow Z'Z'$ and $NN \rightarrow \phi \phi$ can be found in Ref. [16].}
forbidden in the leptogenesis mechanism. We can see that there is plenty of parameter space allowed by leptogenesis for $\mathcal{O}(\text{TeV}) > Z'$ and $N$. Since this mass region is accessible at current or near future colliders [70–74], some of the parameter space is already excluded by the LEP [65] and LHC [66, 67], see also [75–77] searches for $Z' \rightarrow \ell^+\ell^−/jj$, as plotted in the shaded region in the right panel of Fig. 1. Also plotted in the figure is the projected reach for $M_{Z'} \approx \sqrt{s}$, for a $g_{B-L} = 0.8$ and focus on $Z'$ with 6 TeV < $M_{Z'}$ < 30 TeV to perform the collider phenomenology study.

III. COLLIDER PHENOMENOLOGY

A. Same-sign dilepton final state

Since the SM fermions are charged under the $U(1)_{B-L}$ group, $Z'$ can be produced at the LHC via quark fusion $q\bar{q} \rightarrow Z'$ or at the muon colliders via $\mu^+\mu^- \rightarrow Z'\gamma/Z'Z$ (the so-called radiative return) and $\mu^+\mu^- \rightarrow Z'(\rightarrow \ell^+\ell^-/\nu\bar{\nu})\gamma$ at the multi-TeV muon colliders taken/rescaled from Ref. [46]. As shown in the right panel of Fig. 1, the 10 (30) TeV muon collider can reach very low $g_{B-L}$ due to the resonant enhancement where $M_{Z'} \approx \sqrt{s}$. We are mostly interested in the parameter space that is allowed by current data but can be probed at the HL-LHC and future muon colliders; for this sake we fix $g_{B-L} = 0.8$ and focus on $Z'$ with 6 TeV < $M_{Z'}$ < 30 TeV to perform the collider phenomenology study.

channel as it is directly related to the essential ingredients of the leptogenesis mechanism: the new particles $Z'$, $N$ and the $CP$ asymmetry. By reconstructing the invariant masses of the decay products, we can find clues for the $Z'$ and $N$ resonances; while by counting the event number difference between the $\ell^+\ell^+$ and $\ell^−\ell^−$ final states, we can probe the $CP$ violation.

The cross sections of $Z' \rightarrow NN$ for various production channels are plotted in Fig. 2 at the 13 TeV LHC, 10 and 30 TeV muon colliders where $M_N = M_{Z'}/3$ is fixed. We can see the LHC cross section drops rapidly as $M_{Z'}$ increases, and it is only $\sigma \sim 10^{-4}$ fb for $M_{Z'} = 12$ TeV, yielding less than 1 event even at the HL-LHC. In

4 We adopt the FeynRules [81] model file from Refs. [82, 83], which is also publicly available in the FeynRules model database, [84] package for a parton-level simulation.
Table II. The cross sections of the signal and main backgrounds at the 13 TeV LHC and 10 TeV muon collider after the trigger cuts, the same-sign dilepton and W-jet requirements. (*): The decay products of $t\bar{t}$ are further required to have invariant masses $>6$ TeV. The signal process is generated at $M_{Z'} = 8$ TeV, $M_N = 500$ GeV, and $g_{B-L} = 0.8$.

\begin{table}[h]
\centering
\begin{tabular}{|c|c|c|c|}
\hline
LHC & Trigger cut [fb] & Same-sign lepton [fb] & $W$-jet [fb] \\
\hline
Signal & $\sim 10^{-3}$ & $\sim 10^{-3}$ & $\sim 10^{-4}$ \\
$t\bar{t}$ & $\sim 10^{-4}$ (*)& $\lesssim 10^{-7}$ & $\lesssim 10^{-10}$ \\
$W^\pm W^\mp jj$ & $\lesssim 10^{-2}$ & $\lesssim 10^{-4}$ & $\lesssim 10^{-7}$ \\
\hline
10 TeV muon collider & Trigger cut [fb] & Same-sign lepton [fb] & $W$-jet [fb] \\
\hline
Signal & $\sim 1$ & $\sim 1$ & $\sim 10^{-1}$ \\
$\mu^+\mu^- \rightarrow e^+e^- W^+W^-$ & $\sim 10^{-2}$ & $\sim 10^{-5}$ & $\sim 10^{-6}$ \\
$\mu^+\mu^- \rightarrow e^+e^- W^+W^-\gamma/Z$ & $\sim 10^{-2}$ & $\sim 10^{-5}$ & $\sim 10^{-6}$ \\
$\mu^+\mu^- \rightarrow W^+W^-jj$ & $\sim 10^{-1}$ & $\sim 10^{-6}$ & $\sim 10^{-9}$ \\
\hline
\end{tabular}
\end{table}

In contrast, the $\mu^+\mu^- \rightarrow Z'\gamma/Z'Z$ cross section increases when $M_{Z'}$ is close to the collision energy of the muon collider: for $M_{Z'} > 6$ TeV, $\sigma \gtrsim 10$ (1) fb at a 10 (30) TeV muon collider, providing $\gtrsim 10^3$ (9 $\times 10^4$) events at an integrated luminosity of 10 (90) ab$^{-1}$. The 10 TeV muon collider can probe $M_{Z'} > 10$ TeV via the off-shell production of $\mu^+\mu^- \rightarrow Z'^* \rightarrow NN$. In this case, the cross section drops when $M_{Z'}$ increases, showing a tendency similar to the LHC but with much higher rates (for example $\sigma \sim 10$ fb for $M_{Z'} = 12$ TeV).

The RHN decays to a lepton and a vector/Higgs boson, and due to the Goldstone equivalence theorem the branching ratios are approximately

$$\frac{1}{2}\text{Br}(\ell^+W^-) \approx \text{Br}(\ell\nu Z) \approx \text{Br}(\ell\nu h),$$

$$\frac{1}{2}\text{Br}(\ell^-W^+) \approx \text{Br}(\nu\ell Z) \approx \text{Br}(\nu\ell h).$$

The difference between $\text{Br}(\ell^+W^-)$ and $\text{Br}(\ell^-W^+)$ connects to the $CP$ violation $\epsilon$, as defined in Eq. (4). For simplicity, in the following we assume the $N$ decays exclusively to the first generation of leptons, and focus on the $NN \rightarrow e^+e^-Z^\pmZ^\mp$ final state, which is clean due to its same-sign dilepton feature. As we are considering $M_{Z'} \gtrsim 6$ TeV, the $W^\pm$ bosons from cascade decay are typically boosted, and hence can be treated as boosted fat jets.

To trigger the signal events, we require the final state to have exactly two electrons within

$$p_T^e > 100 \text{ GeV}, \quad |\eta_e| < 2.5, \quad \text{(for LHC)}$$

and two jets within

$$p_T^j > 500 \text{ GeV}, \quad |\eta_j| < 2, \quad \text{(for LHC)}$$

and

$$p_T^W > 500 \text{ GeV}, \quad |\eta_W| < 2.43, \quad \text{(for muon colliders)}$$

The electrons are further required to be same-sign, and the jets are required to be $W$-tagged. At the LHC, the main backgrounds come from the dilepton decay of $t\bar{t}$ (with one lepton’s charge misidentified and two QCD jets mistagged as $W$) and $W^\pm W^\mp jj$ (with jets mistagged) [14]. While at the muon colliders, the $t\bar{t}$ background is subdominant; instead, the main backgrounds are $\mu^+\mu^- \rightarrow e^+e^- W^+W^-$ and $\mu^+\mu^- \rightarrow e^+e^- W^+W^-\gamma/Z$ (with lepton charge misidentified), $\mu^+\mu^- \rightarrow W^+W^-jj$ (with jets mistagged). The backgrounds from charge misidentification are significantly suppressed by the misidentification rate, which we adopt as 0.1% based on the LHC detector performance [85]. Note that this is a conservative estimate as the charge identification efficiency is expected to be improved at the muon colliders.

In our parton-level study, the $W$-jet is simulated by a parton-level $W$ boson multiplying by the hadronic decay branching ratio 67.4% [86] and the tagging efficiency 60%. The backgrounds from mistagged QCD jets are then suppressed by the mistag rate, which we adopt as 5%. The $t\bar{t}$ background is further suppressed by requiring the final state objects’ invariant mass to be above 6 TeV, while the same cut almost does not affect the signal rate. The background rates are listed in Table II, where the signal cross sections for $M_{Z'} = 8$ TeV, $M_N = 500$ GeV, and $g_{B-L} = 0.8$ are also given. We can see that at the LHC the backgrounds are under control, while at the muon colliders the backgrounds are negligible after the selection cuts.

At the muon colliders, we can further remove the backgrounds also by putting cuts on the invariant mass of the final state particles. This is shown in Fig. 3, where $p_T^W$ and the invariant masses of the signal and $\mu^+\mu^- \rightarrow e^+e^- W^+W^-\gamma$ at the 10 TeV muon collider are illustrated for one benchmark point with $M_{Z'} = 8$ TeV and $M_N = 500$ GeV. The jet four-momentum is smeared according to a jet energy resolution of $\Delta E/E = 10\%$. For the reconstruction of $N$, we pair the two leptons and $W$-jets by minimizing $\chi^2 = (M_{\ell\ell}W_1 - M_N)^2 + (M_{\ell\ell}W_2 - M_N)^2$. From the figure, it is clear to see that the signal and background have very different kinematical distributions especially when reconstructing the $N$; a cut on the $\ell\ell W$ invariant mass can significantly remove the backgrounds, even in case of $M_{Z'} > \sqrt{s}$. 
dileptons at colliders, upper limits can be put on the CP final state. If we observe no asymmetry of the same-sign N where ϵ > ϵ
"no leptogenesis" is for ab
leptogenesis can be tested at the HL-LHC and the 10 (30) ϵ > ϵ
FIG. 4. The regions where ϵ > ϵ

B. Testing Leptogenesis

We have shown in the previous subsection that the backgrounds for the same-sign dilepton channel are negligible after imposing a set of selection cuts, which however keep the signal rates almost unchanged. There are also other channels available for the Z' → NN process (e.g. N → νh/νZ), which can further increase the signal significance of the model. Based on this consideration, we make the zero background assumption in this subsection to derive the projected sensitivities for the thermal leptogenesis parameter space of the B – L model. The CP asymmetry ϵ defined in Eq. (4) is related to the asymmetry between the positive and negative same-sign dileptons as [14]

\[ \epsilon = \frac{1}{2} \left| \frac{N_+ - N_-}{N_+ + N_-} \right|, \]

where N± denotes the event number of the e±e±W±W± final state. If we observe no asymmetry of the same-sign dileptons at colliders, upper limits can be put on the CP asymmetry. This is done by assuming the number of the signal events to follow a Poisson distribution, and hence at 1σ confidence level [86]

\[ N_+ = \langle N_+ \rangle \pm \Delta N_+ = \langle N_+ \rangle \pm \sqrt{\langle N_+ \rangle}, \]

yielding a sensitivities of

\[ \epsilon_{\text{min}} = \frac{1}{2 \sqrt{\langle N_+ \rangle}}. \]

The CP asymmetry is detectable if ϵ > ϵ_{\text{min}}.

With Eq. (24) in hand, one is able to test resonant leptogenesis at a specific collider environment. Given a set of (M_{Z'}, M_N, g_{B-L}), the ϵ required by leptogenesis is derived by Eq. (16). On the other hand, the corresponding N_+ and hence ϵ_{\text{min}} is also available by collider simulation of the Z' → NN process. If ϵ > ϵ_{\text{min}}, then this parameter setup is detectable at the collider. The projected detectable parameter space with a fixed g_{B-L} = 0.8 is plotted in Fig. 4 for the HL-LHC (13 TeV, 3 ab^{-1}) and two setups of muon colliders (10 TeV, 10 ab^{-1} and 30 TeV, 90 ab^{-1}) with different colors. As the signal significance is proportional to the CP asymmetry ϵ, the reachable regions have similar shapes to the ϵ contours in the left panel of Fig. 1, except that there is a vertical band-like region in the 10 TeV muon collider case near M_{Z'} ≈ 10 TeV due to the enhancement from resonant production of Z'. The region shaded in dark blue is for ϵ > 1 or M_N < T_{sph}.
the parameter space for $M_{Z'} \sim \mathcal{O}(10 \text{ TeV})$ and $M_N \sim 2 \text{ TeV}$ is not reachable for all collider setups, because the $\epsilon$ needed for BAU is too small (see the contours in Fig. 1). Although larger $\epsilon$ can be generated from tuning the $\Delta M_N$ which is a free parameter, the resulting BAU will be greater than the observation value. However, a band in this region can be filled up by the 10 TeV muon collider for $M_{Z'} \sim 10 \text{ TeV}$, due to the resonant enhancement of the $Z'Z/Z\gamma$ process. In summary, a considerable fraction of parameter space can be probed at the future muon colliders.

IV. CONCLUSION

Both the BAU and tiny neutrino mass can be explained simultaneously within the framework of the thermal leptogenesis via the RHNs. If two of the RHNs are nearly degenerate, the scale of leptogenesis can be as low as $\sim 100 \text{ GeV}$ and hence accessible at current and future colliders. In this article, we take the $B-L$ model as an example to perform a study on the collider probe of the leptogenesis mechanism. In this model, the new gauge boson $Z'$ mediates a scattering process that tends to washout the BAU. Therefore, if both $Z'$ and $N$ are at $\mathcal{O}(\text{TeV})$ scale, a sizable CP asymmetry $\epsilon$ is needed to explain the observed BAU. In other words, successful (resonant) leptogenesis can be realized via TeV scale $Z'$, $N$ and a sizable $\epsilon$, which are testable at the colliders.

We choose the channel $Z' \rightarrow NN \rightarrow \ell^+\ell^- + \text{jets}$ to probe leptogenesis, as this process involves the three key ingredients of the scenario: a new gauge boson $Z'$, the RHN $N$, and the CP violation $\epsilon$ from the asymmetry between $N(\ell^+\ell^+)$ and $N(\ell^-\ell^-)$. Three collider setups are considered: 13 TeV LHC with 3 ab$^{-1}$ and 10 (30) TeV muon colliders with 10 (90) ab$^{-1}$. At the LHC, the $Z'$ is produced via the Drell-Yan process; while at the muon colliders, $Z'$ can be produced in association with a $Z'$ or $\gamma$, or via the off-shell $s$-channel fusion. We have shown that the backgrounds are safely negligible after a set of selection cuts, and hence the sensitivities for $\epsilon$ can be derived under the zero background assumption. Our quantitative study shows that leptogenesis in the $B-L$ model can be probed at the LHC for $6 \text{ TeV} \lesssim M_{Z'} \lesssim 12 \text{ TeV}$ and $M_N \sim \text{TeV}$. At the muon colliders, due to the higher signal rates, the projected probe limits can cover $M_{Z'}$ up to $30 \text{ TeV}$ or even higher. For the 30 TeV muon collider, the leptogenesis can be probed up to $M_{Z'} \sim 100 \text{ TeV}$ via the off-shell production of $Z'$. Our work demonstrates that the muon colliders can serve as a machine to efficiently probe the early Universe dynamics.

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