Enhanced Z boson decays as a new probe of first-order electroweak phase transition at future lepton colliders

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We study phenomenological consequences of the strong first-order electroweak phase transition in an extension of the standard model with an inert doublet and vectorlike leptons motivated by the muon $g-2$ anomaly and dark matter. We find that a condition for the strong first-order electroweak phase transition inevitably induces a large logarithmic enhancement in Z boson decays, which relegates the explanation of the anomalous muon $g-2$ at below 2σ level. Our analysis shows that future lepton collider experiments, especially the Giga-Z at the International Linear Collider and Tera-Z at the Circular Electron Positron Collider as well as Future Circular Collider have great capability to explore the nature of the electroweak phase transition, which is complementary to conventional approaches via measurements of the triple Higgs boson coupling and gravitational waves.

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I. INTRODUCTION

From cosmological observations, the baryon asymmetry of the Universe (BAU) is found to be $n_B/n_\gamma = (6.09 \pm 0.06) \times 10^{-10}$ [1], where n_B denotes the baryon number density and n_γ represents the photon number density. To obtain the observed BAU from a baryon symmetric Universe, the following Sakharov conditions [2] must be satisfied: (i) B violation, (ii) C and CP violation and (iii) departure from thermal equilibrium. The last condition could be exempted if CPT is violated.

While a plethora of baryogenesis scenarios are present in the literature [3], the discovery of the Higgs boson at the LHC attracts more people's attention to a scenario of electroweak baryogenesis (EWBG) [4] in which the Higgs physics plays an essential role. One of the necessary ingredients for the successful EWBG is a strong first-order electroweak phase transition (SFOEWPT) that can achieve departure from the thermal equilibrium and prevent the generated BAU from washing out. It is shown by lattice simulations that the 125 GeV Higgs boson is too heavy to realize a SFOEWPT in the standard model (SM) [5], and therefore the minimal Higgs sector has to be extended by introducing, for instance, an additional Higgs doublet. Besides the baryogenesis issue, the scalar extensions of the SM are also motivated by other fundamental problems, such as dark matter (DM), inflation and neutrino masses and mixings, and some of the extended models could provide a solution of the muon magnetic dipole moment $[(g-2)_\mu]$ anomaly as well. Among such extensions, the SM augmented by an inert Higgs doublet [6], right-handed neutrinos and vectorlike leptons (denoted as VLIDM for short) has been studied from the viewpoints of DM and neutrino physics [7], or DM and the $(g-2)_\mu$ anomaly [8], or DM and inflation [9]. One can expect that a SFOEWPT would still be possible in the VLIDM, as is the ordinary inert Higgs doublet model. However, so far there has been no explicit demonstration and compatibility of a SFOEWPT with other observables, $(g-2)_\mu$ in particular is not clear.

In this paper, we study a SFOEWPT and its compatibility with the $(g-2)_\mu$ explanation in the VLIDM, and also discuss phenomenological consequences, focusing particularly on the correlation between a SFOEWPT and the Z boson decays. Here, we consider a case in which the vectorlike leptons preferentially couple to muons, which is motivated by the $(g-2)_\mu$ explanation.

We point out that a condition for SFOEWPT inevitably leads to sizable radiative corrections to $Z \to \mu^+\mu^-$ due to a logarithmic enhancement factor, whereas $(g-2)_\mu$, by contrast, is suppressed, preventing one from explaining the $(g-2)_\mu$ anomaly within 2σ level. Since the essential point in this correlation is a mass splitting between the neutral scalars in the same multiplet, our findings would hold in other models as long as the mass splitting is crucial for realizing a SFOEWPT and EWBG.

We also show that the regions of the SFOEWPT and DM in our scenario can be thoroughly probed by the future lepton collider experiments, especially the precise measurements

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of the Z boson, such as a Giga-Z option at the International Linear Collider (ILC) [10] as well as the Tera-Z phase at the Circular Electron Positron Collider (CEPC) [11] and Future Circular Collider (FCC-ee) [12], which plan to produce around 10^9 and 10^{12} Z bosons, respectively.[1] This would be a new approach to explore the nature of EWPT along with conventional probes by the triple Higgs boson coupling [17,18] and gravitational waves [19].

The paper is organized as follows. In Sec. II, we describe the VLIDM and give a brief overview of DM physics in this model. In Sec. III, we derive the condition of the SFOEWPT and show its impacts on the Z decays and δa_μ analytically. The quantitative studies of the correlations are presented in Secs. IV and V is devoted to the conclusion and discussions. One-loop functions appearing in the Z boson decays are listed in the Appendix.

II. THE MODEL

We study the model in which the inert Higgs doublet (η), right-handed neutrinos ($N_{Ri=1\text{–}3}$) [20] and vectorlike leptons ($E_{i=1\text{–}3}$) are added to the SM [7,8,21]. The SM quantum numbers and the Z_2 parity of each field are assigned as $(\mathbf{1},\mathbf{2},1/2,-)$ for η, $(\mathbf{1},\mathbf{1},0,-)$ for N_{Ri} and $(\mathbf{1},\mathbf{1},-1,-)$ for E_i, respectively. This model offers the possible explanation for the DM, neutrino masses/mixings and the $(g-2)_\mu$ anomaly.[2] Moreover, if we further assume a nonminimal coupling of η to gravity, a successful scenario of inflation can also be accommodated [9]. Owing to the Z_2 parity, the lightest Z_2-odd particle can be stable and becomes the DM candidate. In this work, we focus on the case of the scalar DM by assuming the right-handed neutrinos are heavy enough, and thus they are omitted hereafter.

The new vectorlike lepton interactions can be written as

$$-\mathcal{L} \supset m_{E_i}\bar{E}_{iL}E_{iR} + y_{ij}\bar{\ell}_{iL}\eta E_{jR} + \text{H.c.}, \tag{1}$$

where y_{ij} are the general 3-by-3 complex matrices that may provide the necessary CP-violating sources for EWBG [22].[3] Since we focus on the $(g-2)_\mu$-favored region, we assume that $y_{\mu E_i} \neq 0$ and that other elements are negligibly small. Moreover, we set $y_{\mu E_1} = y_{\mu E_2} = y_{\mu E_3} \equiv y_{\mu E}$ and $m_{E_1} = m_{E_2} = m_{E_3} \equiv m_E$ for the sake of simplicity.

[1]For earlier studies on DM and/or $(g-2)_\mu$-favored regions at the Z factories, see, e.g., Refs. [13–16].

[2]In this model, the neutrino masses are generated by the same mechanism as the Ma model [20], and the vectorlike leptons do not play a significant role.

[3]In principle, there is a possibility of leptogenesis [23] (for recent studies, see, e.g., Refs. [9,24]). The amount of BAU strongly depends on parameters in the lepton sector. In our scenario, the BAU is not generated via leptogenesis by assumption.

The tree-level scalar potential is given by

$$\begin{aligned} V_0(\Phi,\eta) &= \mu_1^2|\Phi|^2 + \mu_2^2|\eta|^2 + \frac{\lambda_1}{2}|\Phi|^4 + \frac{\lambda_2}{2}|\eta|^4 + \lambda_3|\Phi|^2|\eta|^2 \\ &\quad + \lambda_4|\Phi^\dagger\eta|^2 + \left[\frac{\lambda_5}{2}(\Phi^\dagger\eta)^2 + \text{H.c.}\right], \end{aligned}$$

where Φ is the SM Higgs doublet that develops a vacuum expectation value (VEV) with $v = 246$ GeV. The masses of the physical scalar particles at tree level can be written as

$$\begin{aligned} m_h^2 &= \lambda_1 v^2, \\ m_{H^\pm}^2 &= \mu_2^2 + \frac{1}{2}\lambda_3 v^2, \\ m_H^2 &= \mu_2^2 + \frac{1}{2}(\lambda_3+\lambda_4+\lambda_5)v^2 = m_{H^\pm}^2 + \frac{1}{2}(\lambda_4+\lambda_5)v^2, \\ m_A^2 &= \mu_2^2 + \frac{1}{2}(\lambda_3+\lambda_4-\lambda_5)v^2 = m_{H^\pm}^2 + \frac{1}{2}(\lambda_4-\lambda_5)v^2, \end{aligned} \tag{2}$$

where m_h is the SM Higgs boson mass, m_H and m_A are the masses of the CP-even and -odd scalar particles from the inert Higgs doublet η, respectively, and $m_{H^\pm}$ is the charged Higgs boson mass. Without loss of generality, we consider that $\lambda_4+\lambda_5<0$ and $\lambda_5<0$, which renders the CP-even scalar H the lightest Z_2-odd particle and hence the stable DM candidate. As is known, the contributions from the extra scalars to the T parameter becomes zero if $m_{H^\pm}=m_H$ or $m_{H^\pm}=m_A$ [6]. Although the small deviations from those mass relations are still experimentally allowed, we choose $m_{H^\pm}=m_A$, corresponding to $\lambda_4=\lambda_5<0$, in order to make our discussion simpler. We also note that in the parameter space of our interest, S parameter [6] is small.

From the Planck 2018 data [25], the observed DM abundance is determined as

$$\Omega_{\text{DM}}h^2 = 0.11933 \pm 0.00091. \tag{3}$$

If the DM mass is less than the W boson mass, the main DM annihilation processes are the h-mediated s channel and the vectorlike lepton-mediated t channel. Its cross section is approximately given by

$$\sigma v_{\text{rel}} \simeq \frac{N_C^f}{16\pi}\frac{\lambda_L^2 m_f^2\beta_f^3}{(s-m_h^2)^2+m_h^2\Gamma_h^2} + \frac{|y_{\mu E}|^4}{60\pi m_H^2}\frac{v_{\text{rel}}^4}{(1+r_E)^4}, \tag{4}$$

where v_{rel} denotes the relative velocity of the DM, f denote the SM fermions, N_C^f are the color degrees of freedom of f, $\lambda_L = \lambda_3+\lambda_4+\lambda_5$, Γ_h is the width of the Higgs boson, $r_E = m_E^2/m_H^2$, $s = 4m_H^2/(1-v_{\text{rel}}^2/4)$, $\beta_f = \sqrt{1-4m_f^2/s}$. Although the t-channel process is d-wave suppressed, it would come into play if $|y_{\mu E}| \gtrsim 1$ [8].

As for the DM direct detection, the recent XENON1T data put a strong constraint on the DM-nucleon spin-independent elastic scatter cross section $\sigma_{\rm SI}$ [26]. For instance, the most excluded region at 90% confidence level reaches $\sigma_{\rm SI} = 4.1 \times 10^{-47}$ cm^2 with the DM mass of 30 GeV. Therefore, for a light DM, the direct detection data favor the so-called Higgs funnel region where the DM mass is close to half of the Higgs mass, namely, $m_H \simeq m_h/2 \simeq 63$ GeV. In this model, the cross section $\sigma_{\rm SI}$ is approximated as

$$\sigma_{\rm SI} \simeq \frac{\lambda_L^2 f_N^2}{4\pi}\left(\frac{m_N^2}{m_H m_h^2}\right)^2, \tag{5}$$

where $f_N \simeq 0.3$. To evade the current DM direct detection constraints in this Higgs funnel region, $\lambda_L \lesssim 0.003$ is required, which implies that

$$m_H^2 \simeq \mu_2^2, \qquad m_A^2 \simeq m_H^2 + \frac{\lambda_3}{2}v^2. \tag{6}$$

As discussed below, λ_3 has to be $\mathcal{O}(1)$ in magnitude to achieve a SFOEWPT [27–36].

Although Eqs. (4) and (5) make it easy to see the model parameter dependences, we use MICROMEGAs [37] in order to get more precise values of $\Omega_{\rm DM}h^2$ and $\sigma_{\rm SI}$. We have confirmed that our numerical results are consistent with those in Ref. [8] when we take their input parameters.

III. STRONG FIRST-ORDER PHASE TRANSITION AND ITS IMPLICATIONS FOR Z BOSON DECAYS AND $(g-2)_\mu$

In EWBG, the BAU arises via B-violating processes (sphaleron processes) in the symmetric phase. To maintain the generated BAU in the broken phase, the sphaleron processes must be sufficiently suppressed. This is realized if the sphaleron energy, which is proportional to the Higgs VEV at finite temperature, becomes large enough. Conventionally, the condition of the SFOEWPT is approximately described by

$$v_C/T_C \gtrsim 1, \tag{7}$$

where T_C denotes the critical temperature at which two degenerate minima coexist in the finite-temperature effective scalar potential and v_C is the corresponding VEV in the broken phase.

Let us consider how the condition of Eq. (7) is satisfied in this model. The effective potential with high-temperature expansion takes the form

$$V_{\rm eff}(\varphi; T_C) = \frac{\lambda_{T_C}}{4}\varphi^2(\varphi - v_C)^2, \qquad v_C = \frac{2ET_C}{\lambda_{T_C}}, \tag{8}$$

where φ is the classical background field of the SM-like Higgs, E denotes the coefficient of the φ^3 term and λ_{T_C} is the quartic coupling at T_C, which is more or less fixed by the Higgs boson mass. Therefore, E has to be enhanced in order to satisfy Eq. (7). As is well known, the bosonic particles can contribute to E. In this model, E would be enhanced if $\mu_2^2 \ll \lambda_3 v^2/2$, which enforces the large mass splitting among the neutral scalars of the inert doublet, i.e.,

$$m_H \ll m_A. \tag{9}$$

To what extent the mass splitting is needed depends on the condition (7). Even though Eq. (8) is useful to understand the essence of the SFOEWPT, the high-temperature expansion is not always valid. We therefore numerically evaluate v_C/T_C in Sec. IV. Note that since the vectorlike leptons do not generate the thermal cubic term, the SFOEWPT in the parameter space of our interest is essentially the same as that in the ordinary inert doublet model. Note also that too large of a mass splitting breaks the perturbativity of λ_3 as seen from Eq. (6). In what follows, we discuss the implications of the condition (9) for the Z boson decays and $(g-2)_\mu$.

Let us parametrize the Z boson couplings to fermions as

$$\mathcal{L} = -g_Z Z_\mu \bar{f}\gamma^\mu[g^L_{Z\bar{f}f}P_L + g^R_{Z\bar{f}f}P_R]f, \tag{10}$$

where $g_Z = g_2/c_W$ with g_2 being the SU(2) gauge coupling and c_W the cosine of the weak mixing angle. With those Z boson couplings, the partial decay width of $Z \to \ell^+\ell^-$ can be written as

$$\Gamma(Z \to \ell^+\ell^-) = \frac{m_Z}{24\pi}g_Z^2[|g^L_{Z\bar{\ell}\ell}|^2 + |g^R_{Z\bar{\ell}\ell}|^2], \tag{11}$$

where the SM lepton masses are ignored. We parametrize the new physics effects as $g^{L,R}_{Z\bar{\ell}\ell} = g^{L,R,{\rm SM}}_{Z\bar{\ell}\ell} + \Delta g^{L,R}_{Z\bar{\ell}\ell}$. It is straightforward to calculate the new physics effects by evaluating the Feynman diagrams as depicted in Fig. 1 (the leg correction to μ^- also exists.). In the current model, $\Delta g^R_{Z\bar{\mu}\mu} = 0$ and

$$\begin{aligned}\Delta g^L_{Z\bar{\mu}\mu} = \frac{3|y_{\mu E}|^2}{32\pi^2}\Bigg[&\tilde{F}_3(m_E, m_H, m_A) \\ &+ \sum_{\phi=H,A}\left\{\left(-\frac{1}{2}+s_W^2\right)F_2(m_E, m_\phi)\right. \\ &\left.+ s_W^2 F_3(m_E, m_\phi)\right\}\Bigg],\end{aligned} \tag{12}$$

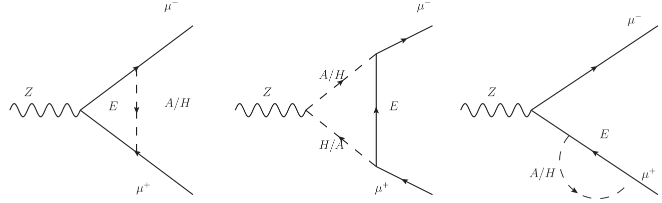

FIG. 1. Representative Feynman diagrams for the one-loop corrections to $Z \to \mu^+\mu^-$.

where the loop functions F_2, F_3 and $\tilde{F}_3$ are listed in the Appendix. One can show that $\Delta g^L_{Z\bar{\mu}\mu} \to 0$ for $m_\phi(\simeq m_Z) \ll m_E$, $m_E(\simeq m_Z) \ll m_\phi$ and $m_Z \ll m_\phi = m_E$ as long as $m_\phi = m_H = m_A$. For m_H, $m_E \ll m_A$, however, $\Delta g^L_{Z\bar{\mu}\mu}$ would be enhanced by $\ln(m_A^2/m_H^2)$, which arises from the correction of the middle triangle diagram and the right leg corrections of $\mu^\pm$ in Fig. 1. As a result, the corrections to $\Gamma(Z \to \mu^+\mu^-)$ in this limit is cast into the form

$$\Delta\Gamma(Z \to \mu^+\mu^-) \simeq \frac{m_Z^2 g_Z^2 |y_{\mu E}|^2}{128\pi^3}\left(-\frac{1}{2} + s_W^2\right)\left[C + \frac{1}{4}\ln\frac{m_A^2}{m_H^2}\right], \quad (13)$$

where C denotes nonlogarithmic contributions. We are aware that such a nondecoupling behavior by the mass splitting is already noticed in the calculation of $Z \to b\bar{b}$ in the two-Higgs doublet model [38], and more recently in the study of $(g-2)_\mu$ in the lepton-specific two-Higgs doublet model [39].[4] Nevertheless, to our best knowledge, the importance of its correlation with a SFOEWPT has not been well recognized in the literature and therefore detailed numerical studies will be conducted below.

Since the vectorlike leptons couple only to $\mu^\pm$ in this model, the lepton flavor universality of Z boson decays is violated. We thus utilize

$$R_{\mu/e} = \frac{\Gamma(Z \to \mu^+\mu^-)}{\Gamma(Z \to e^+e^-)} \quad (14)$$

to test this model precisely. Its current experimental value is $R^{\rm EXP}_{\mu/e} = 1.0009 \pm 0.0028$ [1]. Let us define the deviation of $R_{\mu/e}$ from the SM value as

$$\Delta R_{\mu/e} \equiv \frac{R_{\mu/e} - R^{\rm SM}_{\mu/e}}{R^{\rm SM}_{\mu/e}} \simeq \frac{2g^{L,{\rm SM}}_{Z\bar{\mu}\mu}{\rm Re}(\Delta g^L_{Z\bar{\mu}\mu}) + |\Delta g^L_{Z\bar{\mu}\mu}|^2}{|g^{L,{\rm SM}}_{Z\bar{\mu}\mu}|^2 + |g^{R,{\rm SM}}_{Z\bar{\mu}\mu}|^2}, \quad (15)$$

where $g^{L,{\rm SM}}_{Z\bar{\mu}\mu} \simeq -0.27$ and $g^{R,{\rm SM}}_{Z\bar{\mu}\mu} \simeq 0.23$. As the experimental constraints, we require that $\Delta R_{\mu/e} < 2.8 \times 10^{-3}$.

As is the case of $\Gamma(Z \to \mu^+\mu^-)$, $\Gamma(Z \to \nu\bar{\nu})$ is also modified by the new particles as

$$\Delta g^L_{Z\bar{\nu}\nu} = \frac{3|y_{\mu E}|^2}{16\pi^2}\left[s_W^2\{F_3(m_E, m_{H^\pm}) + \tilde{F}_3(m_E, m_{H^\pm}, m_{H^\pm})\} + \frac{1}{2}\{F_2(m_E, m_{H^\pm}) - \tilde{F}_3(m_E, m_{H^\pm}, m_{H^\pm})\}\right], \quad (16)$$

[4]In the lepton-specific model, leptonic τ decays are also enhanced by the similar logarithmic contribution, which is as important as the Z boson decays [39–41].

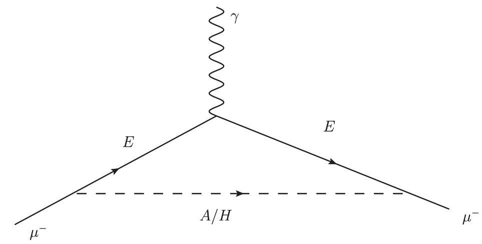

FIG. 2. Schematic Feynman diagrams for the one-loop contributions to $(g-2)_\mu$.

and $\Delta g^R_{Z\bar{\nu}\nu} = 0$. Unlike the $Z \to \mu^+\mu^-$ case, $\Delta g^L_{Z\bar{\nu}\nu}$ does not have logarithmic enhancement due to the absence of the mass splitting. Since this quantity is always numerically unimportant, we do not discuss it henceforward.

Here, we also make a comment on other experimental constraint, especially the $W-\mu-\nu_\mu$ coupling ($\Delta g^L_{W\mu\nu}$) whose deviation from the SM induces the lepton flavor nonuniversality in the muon decay. As with $\Delta g^L_{Z\bar{\mu}\mu}$, $\Delta g^L_{W\mu\nu}$ would receive the logarithmic enhancement of $\ln(m_A^2/m_H^2)$ in the limit of $m_H \ll m_A$. In our parameter space, however, $R_{\mu/e}$ gives a stronger bound than this constraint so that we do not discuss it in the following.

Now we move to considering the new physics effects on $(g-2)_\mu$ in the case of $m_H \ll m_A$.

The discrepancy of $(g-2)_\mu$ between the experimental value and the SM prediction is estimated as [42]

$$\delta a_\mu = a_\mu^{\rm EXP} - a_\mu^{\rm SM} = (26.1 \pm 8.0) \times 10^{-10}. \quad (17)$$

In this model, the one-loop contributions to $(g-2)_\mu$ arise from the diagram shown in Fig. 2, from which it is found to be [43]

$$\delta a_\mu = \sum_{\phi=H,A} \frac{3|y_{\mu E}|^2}{32\pi^2} S_1(r_{E\mu}, r_{\phi\mu}), \quad (18)$$

where $r_{E\mu} = m_E^2/m_\mu^2$, $r_{\phi\mu} = m_\phi^2/m_\mu^2$ with m_μ being the muon mass and

$$S_1(r_1, r_2) = \int_0^1 dx \frac{x^2(1-x)}{x(x-1) + xr_1 + (1-x)r_2}. \quad (19)$$

For $m_E \simeq m_H \ll m_A$, δa_μ is approximated as

$$\delta a_\mu \simeq \frac{3|y_{\mu E}|^2}{32\pi^2}\left[\frac{m_\mu^2}{12m_H^2} + \frac{m_\mu^2}{m_A^2}\left\{\frac{1}{3} + \frac{m_E^2}{m_A^2}\left(\frac{11}{6} + \ln\frac{m_E^2}{m_A^2}\right)\right\}\right]. \quad (20)$$

Therefore, the heavy particle simply decouples, which is in stark contrast to the case of $Z \to \mu^+\mu^-$. This is understandable since the logarithmic enhancement originates from the vertex correction of the middle diagram and the leg corrections of $\mu^\pm$ in Fig. 1 while their counterparts are absent in the diagram of δa_μ.

IV. NUMERICAL ANALYSIS

Before showing our numerical results, we outline the current experimental constraints. As mentioned above, the DM data enforce that $m_H \simeq m_h/2$ and $\lambda_L \lesssim 0.003$ for the light DM case. As for the LHC constraints, there are two important processes (for recent studies, see, e.g., Refs. [44,45]). One is the dimuon plus missing energy (MET) process, $q\bar{q} \to E^+E^- \to \mu^+\mu^- + \mathrm{MET}$, from which the vectorlike lepton mass is bounded as $105\ \mathrm{GeV} \lesssim m_E \lesssim 125\ \mathrm{GeV}$ [8]. The other is the monojet plus MET, $gg \to gh \to gHH$. However, this cross section would be suppressed by λ_L^2 if the aforementioned DM constraint is taken into account. Other than those, the potentially relevant constraint is the signal strength of the Higgs boson decays to two photons ($\mu_{\gamma\gamma}$), which can be affected by the charged Higgs bosons. In parameter space of the SFOEWPT, it is known that $\mu_{\gamma\gamma} \simeq 0.9$ [32], which is consistent with the current LHC data $\mu_{\gamma\gamma}^{\mathrm{ATLAS}} = 0.99^{+0.15}_{-0.14}$ [46] and $\mu_{\gamma\gamma}^{\mathrm{CMS}} = 1.18^{+0.17}_{-0.14}$ [47] within 2σ level.

As a benchmark scenario, we consider

$$m_E = (105\text{–}125)\ \mathrm{GeV}, \qquad |y_{\mu E}| = 1.0, 0.5 \text{ and } 0.3. \tag{21}$$

The DM relic abundance is always satisfied by judiciously choosing m_H and λ_L. For instance, for $m_E = 110$ GeV and $|y_{\mu E}| = 0.5$, the choice of $m_H = 62.55$ GeV and $\lambda_L = 0.001$ gives $\Omega_{\mathrm{DM}}h^2 = 0.12$ and $\sigma_{\mathrm{SI}} = 8.7 \times 10^{-48}\ \mathrm{cm}^2$. Here, we set $m_A = m_{H^\pm} = 300$ GeV and $\lambda_2 = 0.3$, though they are not sensitive to the results.

It would be nice if the above scenario can be tested at the LHC Run 3 or high luminosity (HL)-LHC. However, it might be difficult since such light vectorlike leptons could escape from the searches via the soft lepton plus MET as well as the dilepton plus MET [8]. Furthermore, the monojet plus MET process can constrain λ_L only down to $\mathcal{O}(0.1)$ at the HL-LHC [44], which is still way above the range of our interest. It is definitely worth conducting dedicated studies taking all the detailed information into account [21,45]. In this work, however, we consider detectability at future lepton colliders instead, which can offer more robust tests of the scenario.

We firstly consider the case in which the vectorlike leptons cannot be pair produced at the future lepton collider with the center of mass energy of 240 GeV (CEPC, FCC-ee) or 250 GeV (ILC), respectively, i.e., $m_E = 120$ or 125 GeV. Since physics discussion would not differ between the two cases, we only present the $m_E = 120$ GeV case below. After that, we also consider the case in which the vectorlike leptons can be directly produced at those colliders.

In Fig. 3, $\Delta R_{\mu/e}$ is plotted as a function of the heavy scalar mass m_A. The current upper bound of $\Delta R_{\mu/e}$ is represented by the upper horizontal dashed line. The solid curves in red, blue and black correspond to the deviation of

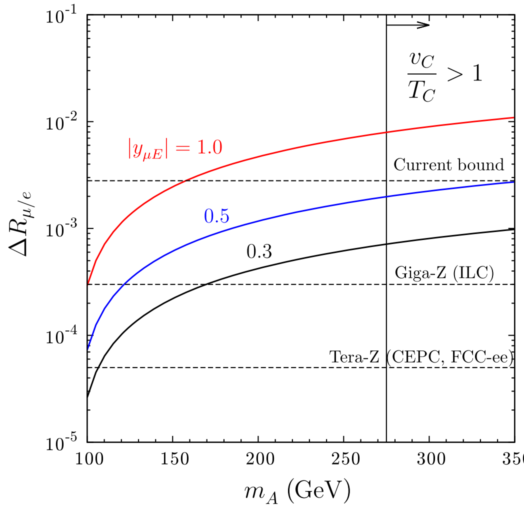

FIG. 3. Deviation of the lepton flavor universality in the Z boson decays $\Delta R_{\mu/e}$ as a function of m_A with $m_E = 120$ GeV. The solid curves in red, blue and black correspond to the deviation $\Delta R_{\mu/e}$ in the Z boson decay for $|y_{\mu E}| = 1.0$, 0.5 and 0.3, respectively. The upper, middle and lower horizontal dashed lines represent the current bound, future Giga-Z sensitivity and Tera-Z sensitivity of $\Delta R_{\mu/e}$, respectively. The region of the SFOEWPT is the right-hand side of the vertical line. The greater v_C/T_C corresponds to the larger enhancement of $Z \to \mu^+\mu^-$.

$\Delta R_{\mu/e}$ in the Z boson decay for $|y_{\mu E}| = 1.0$, 0.5 and 0.3, respectively. In all cases, $\Delta R_{\mu/e}$ gets enhancement as m_A increases, which is attributed to the logarithmic enhancement of $\ln(m_A^2/m_H^2)$ as discussed above. As a result, m_A cannot exceed around 160 GeV for $|y_{\mu E}| = 1.0$ while m_A is allowed up to 350 GeV for $|y_{\mu E}| = 0.5$ and 0.3.[5] We also display the region of $v_C/T_C > 1$ as the right-hand side of the vertical line [27–36].[6] It can been seen that the SFOEWPT requires the large mass splitting of m_H and m_A as explained in the previous section, which leads to the strong correlation between the Z boson decays and the SFOEWPT, i.e., the stronger the SFOEWPT becomes, the more $Z \to \mu^+\mu^-$ is enhanced. One can see that the case of $|y_{\mu E}| = 1.0$ is not consistent with the SFOEWPT. It is an interesting and important question how large $\mathrm{Im}(y_{\mu E})$ is needed for the sufficient BAU and whether it is compatible with the results obtained here for successful baryogenesis [22].

[5] There exists an upper bound on m_A coming from the perturbativity of λ_3 or stability of the DM. In the IDM and its extensions, the latter gives the stronger bound. In the IDM, $m_A^{\max} \simeq (300\text{–}350)$ GeV modulated by the input parameters [32,35]. In the VLIDM, on the other hand, the presence of the vectorlike leptons can push it upward [22].

[6] Note that there still exist a lot of theoretical uncertainties in an ordinary perturbative treatment of an EWPT [36] that we adopt here based on Ref. [35]. Thus, when we interpret the $v_C/T_C = 1$ line, such uncertainties are kept in mind. Nevertheless, the correlation between the enhancement of $Z \to \mu^+\mu^-$ and the SFOEWPT is still robust.

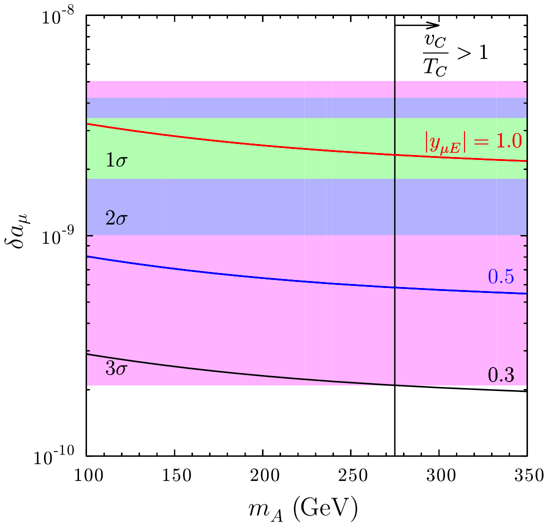

FIG. 4. Corrections to the muon magnetic dipole moment anomaly δa_μ as a function of m_A. The shaded regions in green, blue and magenta correspond to the $(g-2)_\mu$ regions within 1σ, 2σ and 3σ, respectively. The solid curves in red, blue and black correspond to the deviation δa_μ for $|y_{\mu E}| = 1.0$, 0.5 and 0.3, respectively.

Future sensitivities of $\Delta R_{\mu/e}$ are also shown by the middle and lower horizontal dashed lines, where the former is the Giga-Z at the ILC [10] and the latter is Tera-Z such as the CEPC [11] and FCC-ee [12].[7] From Fig. 3, it is found that the future precise measurements of the Z boson can provide new and thorough tests of the SFOEWPT in our benchmark scenario. This new probe can give the cross-check of the EWPT together with the conventional approaches through the measurements of the triple Higgs boson coupling [17,18] and gravitational waves [19]. It should be emphasized that depending on the parameter regions, the precise measurements of the Z boson decays are more powerful than the conventional methods [48].

Now we investigate the compatibility with $(g-2)_\mu$ in our benchmark scenario. In Fig. 4, δa_μ is shown as a function of m_A. The color coordinates of the solid curves are the same as those in Fig. 3. The shaded regions in green, blue and magenta correspond to the $(g-2)_\mu$ regions within 1σ, 2σ and 3σ, respectively. One can see that the case of $|y_{\mu E}| = 1.0$ can explain $(g-2)_\mu$ at 1σ level. However, as shown in Fig. 3, the region of $m_A \gtrsim 160$ GeV is excluded by the measurement of $R_{\mu/e}$, thereby the SFOEWPT and $(g-2)_\mu$ at the 1σ explanation are not compatible. We find that the explanation of $(g-2)_\mu$ is impossible within 2σ level in the regions of the SFOEWPT.

Now let us consider the case in which the pair production of the vectorlike leptons are kinematically

[7] Specifically, the sensitivities at the CEPC and FCC-ee are not necessarily the same due to different machine properties, etc. Here, we just ignore such a difference for simplicity.

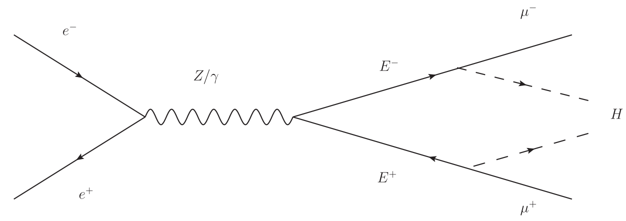

FIG. 5. Direct search for the vectorlike leptons via the process involving dimuon plus MET at the lepton colliders.

allowed. Since the vectorlike leptons exclusively decay into the DM and muons, the most relevant process is $e^+e^- \to E^+E^- \to \mu^+\mu^- + 2H$, as shown in Fig. 5. The dominant contribution to the production cross section of the vectorlike leptons is

$$\sigma(e^+e^- \to \gamma \to E^+E^-) = \frac{4\pi\alpha^2}{3s^2}(s + 2m_E^2)\sqrt{1 - \frac{4m_E^2}{s}}, \tag{22}$$

where α denotes the fine structure constant at the scale of Z boson mass, s is the square of the center of mass energy. Here, the masses of the electron and positrons are neglected.

Figure 6 shows the cross section $\sigma(e^+e^- \to \gamma/Z \to E^+E^-)$ as a function of m_E, which depends only on the collider energy and the vectorlike lepton mass. Here, we also include $\sigma(e^+e^- \to Z \to E^+E^-)$ and the interference effects between the Z boson and photon me diators. It is found that the cross section gets enhanced with decreasing m_E and reaches about $\mathcal{O}(1)$ pb, which is large enough to be measured at future lepton colliders.

We can use the direct production channel to fix the vectorlike lepton mass, and the mass splitting between the neutral scalars and the Yukawa coupling $|y_{\mu E}|$ can be extracted from the precise measurements of the Z boson decays, and if available, together with the triple Higgs

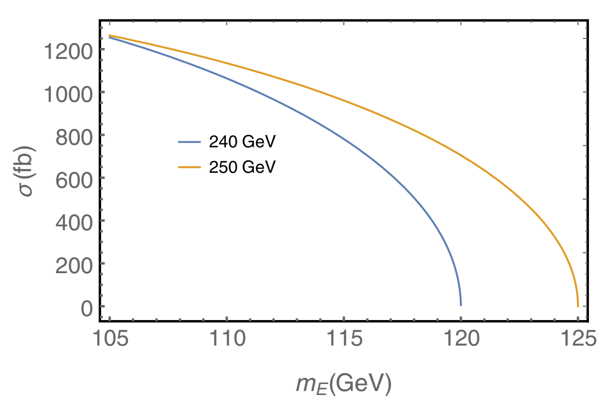

FIG. 6. The cross section $\sigma(e^+e^- \to \gamma/Z \to E^+E^-)$ as a function of m_E. We take $\sqrt{s} = 240$ GeV (blue) and 250 GeV (orange), respectively.

boson coupling [17,18] and/or gravitational waves [19]. In this way, our scenario can be fully tested.

V. CONCLUSION AND DISCUSSIONS

We have studied the possibility of a SFOEWPT and its phenomenological consequences in the VLIDM. It is found that the significant mass splitting between the DM H and the CP-odd Higgs boson A is required to induce a SFOEWPT. Such a condition inevitably leads to the enhanced $\Gamma(Z \to \mu^+\mu^-)$ owing to the logarithmic enhancement of $\ln(m_A^2/m_H^2)$. As a result, the couplings of the vectorlike leptons to the muons have to be small in order to evade the current experimental constraints of the Z boson decays, which limits the size of the corrections to $(g-2)_\mu$. Our numerical studies show that $(g-2)_\mu$ cannot be explained within 2σ level in the region of a SFOEWPT. In other words, the EWPT would be weak first order if $(g-2)_\mu$ is confirmed in this model.

We also showed that the precise measurement of $Z \to \mu^+\mu^-$ at future lepton colliders, such as the ILC Giga-Z and CEPC/FCC-ee Tera-Z as well as the direct search of the vectorlike leptons via the process $e^+e^- \to \gamma/Z \to E^+E^-$ can provide new exquisite probes of the SFOEWPT.

It should be emphasized that the deep correlation between the SFOEWPT and the enhancement of the Z boson decays would be generic as long as the SFOEWPT requires the large mass splitting of the scalar mass spectrum in the same multiplet, which opens a novel and promising avenue to probe thermal history of EWPT in the early Universe and EWBG in addition to the well-studied approaches using the triple Higgs coupling and gravitational waves.

As a by-product, our study clarified that the future lepton colliders, especially the Z factories, can provide a new alternative approach to explore the DM blind spots, where the DM-Higgs coupling λ_L is too small to be detected by the DM direct detection.

Lastly, we make a remark about other DM scenarios such as the heavy scalar DM ($m_H \gtrsim 500$ GeV) [49] and the right-handed neutrino DM [50]. In the former case, he SFOEWPT induced by the thermal loop effects would not be realized since the conditions of $\mu_2^2 \ll \lambda_3 v^2/2$ cannot be satisfied. In the latter case, on the other hand, λ_5 has to be much smaller than unity to be consistent with neutrino and DM physics, which implies that $m_{H^\pm} = m_A \simeq m_H$. It is still possible to have a SFOEWPT as long as $\mu_2^2 \ll \lambda_3 v^2/2$. However, its correlation with the Z decays would be lost due to the lack of a large mass splitting between H and A.

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APPENDIX: LOOP FUNCTIONS

The one-loop functions appearing in $\Delta g^L_{Z\bar{\ell}\ell}$ and $\Delta g^L_{Z\bar{\nu}\nu}$ are defined as

$$F_2(m_E, m_\phi) = \int_x x \ln[(1-x)m_E^2 + x m_\phi^2], \quad \text{(A1)}$$

$$F_3(m_E, m_\phi) = \int_{xy} \left[\frac{xym_Z^2 + m_E^2}{\Delta_3} - 1 - \ln \Delta_3\right], \quad \text{(A2)}$$

$$\tilde{F}_3(m_E, m_H, m_A) = \int_{xy} \ln \tilde{\Delta}_3, \quad \text{(A3)}$$

where $\int_x = \int_0^1 dx$, $\int_{xy} = \int_0^1 dx \int_0^{1-x} dy$ and

$$\Delta_3 = -xym_Z^2 + (x+y)m_E^2 + m_\phi^2(1-x-y), \quad \text{(A4)}$$

$$\tilde{\Delta}_3 = -xym_Z^2 + xm_H^2 + ym_A^2 + m_E^2(1-x-y). \quad \text{(A5)}$$

Incidentally, for $m_H = m_A = m_\phi$, our loop functions are reduced to those in Ref. [8]:

$$I_a(m_E, m_\phi) = F_3(m_E, m_\phi) + \tilde{F}_3(m_E, m_\phi), \quad \text{(A6)}$$

$$I_c(m_E, m_\phi) = F_2(m_E, m_\phi) - F_3(m_E, m_\phi). \quad \text{(A7)}$$

[1] M. Tanabashi et al. (Particle Data Group), Phys. Rev. D **98**, 030001 (2018).

[2] A. D. Sakharov, Pis'ma Zh. Eksp. Teor. Fiz. **5**, 32 (1967) [JETP Lett. **5**, 24 (1967)]; Usp. Fiz. Nauk **161**, 61 (1991) [Sov. Phys. Usp. **34**, 392 (1991)].

[3] M. Dine and A. Kusenko, Rev. Mod. Phys. **76**, 1 (2003).

[4] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, Phys. Lett. **155B**, 36 (1985); see, e.g., A. G. Cohen, D. B. Kaplan, and A. E. Nelson, Annu. Rev. Nucl. Part. Sci. **43**, 27 (1993); M. Quiros, Helv. Phys. Acta **67**, 451 (1994); V. A. Rubakov and M. E. Shaposhnikov, Usp. Fiz. Nauk **166**, 493 (1996) [Sov. Phys. Usp. **39**, 461 (1996)]; K. Funakubo, Prog. Theor. Phys. **96**, 475 (1996); A. Riotto,

arXiv:hep-ph/9807454; M. Trodden, Rev. Mod. Phys. **71**, 1463 (1999); W. Bernreuther, Lect. Notes Phys. **591**, 237 (2002); J. M. Cline, arXiv:hep-ph/0609145; D. E. Morrissey and M. J. Ramsey-Musolf, New J. Phys. **14**, 125003 (2012); T. Konstandin, Usp. Fiz. Nauk **183**, 785 (2013) [Sov. Phys. Usp. **56**, 747 (2013)].
[5] K. Kajantie, M. Laine, K. Rummukainen, and M. E. Shaposhnikov, Phys. Rev. Lett. **77**, 2887 (1996); K. Rummukainen, M. Tsypin, K. Kajantie, M. Laine, and M. E. Shaposhnikov, Nucl. Phys. **B532**, 283 (1998); F. Csikor, Z. Fodor, and J. Heitger, Phys. Rev. Lett. **82**, 21 (1999); Y. Aoki, F. Csikor, Z. Fodor, and A. Ukawa, Phys. Rev. D **60**, 013001 (1999).
[6] R. Barbieri, L. J. Hall, and V. S. Rychkov, Phys. Rev. D **74**, 015007 (2006).
[7] D. Borah, S. Sadhukhan, and S. Sahoo, Phys. Lett. B **771**, 624 (2017).
[8] L. Calibbi, R. Ziegler, and J. Zupan, J. High Energy Phys. 07 (2018) 046.
[9] D. Borah, P. S. B. Dev, and A. Kumar, Phys. Rev. D **99**, 055012 (2019).
[10] A. Irles, R. Pschl, F. Richard, and H. Yamamoto, arXiv: 1905.00220.
[11] J. Guimaraes da Costa *et al.* (CEPC Study Group), arXiv: 1811.10545.
[12] A. Blondel *et al.*, arXiv:1809.01830.
[13] Q. H. Cao, Y. Li, B. Yan, Y. Zhang, and Z. Zhang, Nucl. Phys. **B909**, 197 (2016).
[14] E. Coluccio Leskow, G. D'Ambrosio, A. Crivellin, and D. Müller, Phys. Rev. D **95**, 055018 (2017).
[15] K. Kowalska and E. M. Sessolo, J. High Energy Phys. 09 (2017) 112.
[16] J. Liu, L. T. Wang, X. P. Wang, and W. Xue, Phys. Rev. D **97**, 095044 (2018).
[17] C. Grojean, G. Servant, and J. D. Wells, Phys. Rev. D **71**, 036001 (2005).
[18] S. Kanemura, Y. Okada, and E. Senaha, Phys. Lett. B **606**, 361 (2005).
[19] F. P. Huang and J. H. Yu, Phys. Rev. D **98**, 095022 (2018).
[20] E. Ma, Phys. Rev. D **73**, 077301 (2006).
[21] B. Barman, D. Borah, L. Mukherjee, and S. Nandi, arXiv: 1808.06639.
[22] Fa Peng Huang and Eibun Senaha (to be published).
[23] M. Fukugita and T. Yanagida, Phys. Lett. B **174**, 45 (1986).
[24] T. Hugle, M. Platscher, and K. Schmitz, Phys. Rev. D **98**, 023020 (2018).
[25] N. Aghanim *et al.* (*Planck* Collaboration), arXiv:1807.06209.
[26] E. Aprile *et al.* (XENON Collaboration), Phys. Rev. Lett. **121**, 111302 (2018).
[27] T. A. Chowdhury, M. Nemevsek, G. Senjanovic, and Y. Zhang, J. Cosmol. Astropart. Phys. 02 (2012) 029.
[28] D. Borah and J. M. Cline, Phys. Rev. D **86**, 055001 (2012).
[29] G. Gil, P. Chankowski, and M. Krawczyk, Phys. Lett. B **717**, 396 (2012).
[30] J. M. Cline and K. Kainulainen, Phys. Rev. D **87**, 071701 (2013).
[31] S. S. AbdusSalam and T. A. Chowdhury, J. Cosmol. Astropart. Phys. 05 (2014) 026.
[32] N. Blinov, S. Profumo, and T. Stefaniak, J. Cosmol. Astropart. Phys. 07 (2015) 028.
[33] Q. H. Cao, F. P. Huang, K. P. Xie, and X. Zhang, Chin. Phys. C **42**, 023103 (2018).
[34] M. Laine, M. Meyer, and G. Nardini, Nucl. Phys. **B920**, 565 (2017).
[35] E. Senaha, arXiv:1811.00336.
[36] K. Kainulainen, V. Keus, L. Niemi, K. Rummukainen, T. V. I. Tenkanen, and V. Vaskonen, J. High Energy Phys. 06 (2019) 075.
[37] D. Barducci, G. Belanger, J. Bernon, F. Boudjema, J. Da Silva, S. Kraml, U. Laa, and A. Pukhov, Comput. Phys. Commun. **222**, 327 (2018).
[38] H. E. Haber and H. E. Logan, Phys. Rev. D **62**, 015011 (2000).
[39] E. J. Chun and J. Kim, J. High Energy Phys. 07 (2016) 110.
[40] T. Abe, R. Sato, and K. Yagyu, J. High Energy Phys. 07 (2015) 064.
[41] L. Wang, J. M. Yang, M. Zhang, and Y. Zhang, Phys. Lett. B **788**, 519 (2019).
[42] K. Hagiwara, R. Liao, A. D. Martin, D. Nomura, and T. Teubner, J. Phys. G **38**, 085003 (2011).
[43] J. P. Leveille, Nucl. Phys. **B137**, 63 (1978).
[44] A. Belyaev, G. Cacciapaglia, I. P. Ivanov, F. Rojas-Abatte, and M. Thomas, Phys. Rev. D **97**, 035011 (2018).
[45] A. Bhardwaj, P. Konar, T. Mandal, and S. Sadhukhan, arXiv:1905.04195.
[46] M. Aaboud *et al.* (ATLAS Collaboration), Phys. Rev. D **98**, 052005 (2018).
[47] A. M. Sirunyan *et al.* (CMS Collaboration), J. High Energy Phys. 11 (2018) 185.
[48] F. P. Huang, Y. Wan, D. G. Wang, Y. F. Cai, and X. Zhang, Phys. Rev. D **94**, 041702 (2016).
[49] S. Banerjee, F. Boudjema, N. Chakrabarty, G. Chalons, and H. Sun, arXiv:1906.11269.
[50] A. Vicente and C. E. Yaguna, J. High Energy Phys. 02 (2015) 144.