Anomaly-free leptophilic axionlike particle and its flavor violating tests

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Motivated by the recent Xenon1T result, we study a leptophilic flavor-dependent anomaly-free axionlike particle (ALP) and its effects on charged-lepton flavor violation. We present two representative models. The first one considers that the ALP origins from the *flavon* that generates the charged-lepton masses. The second model assumes a larger flavor symmetry such that more general mixings in the charged-lepton are possible, while maintaining flavor-dependent ALP couplings. We find that a keV ALP explaining the Xenon1T result is still viable for lepton flavor violation and stellar cooling astrophysical limits. On the other hand, if the Xenon1T result is confirmed, future charged-lepton flavor violation measurements can be complementary to probe such a possibility.

I. INTRODUCTION

Recently, the Xenon collaboration reported the observation of an excess in the electron recoiling energy around the keV scale in the Xenon1T detector [1]. Shortly after its announcement, a lot of theoretical work has been done to interpret the results in the context of different extensions of the Standard Model, like axionlike particles (ALPs) [2–10], dark matter [11–35], neutrinos [36–45] and solar axions [46–48], which, however, are subject to stringent constraints from stellar cooling [49–51].

In this work we focus on the ALP framework to explain this excess. This scenario assumes the existence of an ALP as a sizeable part of the observed dark matter (DM) abundance [2], with a mass of a few keV and a relatively weak coupling to the electron. However, constraints from x-ray observations forbid the existence of an anomalous coupling of the ALP to photons for $m_a \gtrsim 0.1$ keV.

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Therefore, we are forced to consider an anomaly-free ALP with respect to $U(1)_{em}$, to avoid these bounds. With the SM fermionic particle content (plus right-handed neutrinos), only hypercharge and $B - L$ are completely anomaly free with family universal charges. Hypercharge can be immediately discarded as it cannot be broken above the electroweak scale. The breaking of $B - L$ gives rise to a pseudo-Goldstone boson, the Majoron, coupling only to neutrinos at tree level. Although it couples to charged leptons at one-loop level, the $B - L$ breaking scale is related to the right-handed Majorana masses that is necessarily too high to explain Xenon1T result with our minimal particle content. Lepton number, also anomaly free with respect to electromagnetism, faces the same problem with family universal charges.

Thus, we have to consider an anomaly-free $U(1)$ symmetry with family-dependent charges which, as we will see, necessarily implies flavor-changing couplings between the ALP and the SM-lepton sector. If this excess is confirmed in the future, it will be necessary to investigate the lepton flavor violating signatures of this particle in low-energy experiments. In this paper, we consider the flavor violation effects induced by such anomaly-free ALP, and we show that lepton flavor violating (LFV) measurements are essential to probe this possibility.

Our paper is organized as follows: in Sec. II, we present two different models, both flavor dependent, but with distinct mixing patterns; the most important constraints to our models are collected in Sec. III; in Sec. IV, we

present our results and discuss how these flavored models can be tested by LFV data; finally, Sec. V is dedicated to our conclusions.

II. MODELS

We consider a $U(1)_\phi$ global symmetry spontaneously broken by the vev of a complex scalar field, ϕ, whose angular component is identified with an ALP. We propose two models with flavor dependence on the lepton sector and evaluate the importance of present and future experiments on LFV decays. In the first model, the presence of the ALP is directly connected to the SM flavor puzzle and the breaking of the $U(1)_\phi$ is the only responsible of the observed hierarchy among the lepton generations. Instead, model II generalizes the previous structure assuming the existence of a larger symmetry, which includes $U(1)_\phi$ as a subgroup, whose breaking produces the Yukawa structures at high energies. In this way, we can partly decouple the nonanomalous flavor-dependent $U(1)_\phi$ charges from the observed leptonic masses and mixings. In both models, below the $U(1)_\phi$-breaking scale, the ALP has flavor-dependent couplings.

A. Model I: hierarchical mixing

Flavor symmetries à la Froggatt-Nielsen [52] offer an attractive solution to the origin of the observed hierarchy among the charged-fermion families. In its simplest version, the spontaneous breaking of a U(1) flavor symmetry by the vev of a scalar field, usually called *flavon*, generates it as powers of the ratio between its vev, v_ϕ, and Λ, the scale at which the heavy fields mediating the processes live, $\epsilon = |v_\phi/\Lambda|$.

In model I, we identify this symmetry with the global $U(1)_\phi$ so that the angular component of the *flavon* corresponds to the ALP. The case of the anomalous QCD axion has been previously explored in [53,54], with the scalar receiving the name of *flaxion* or *axiflavon*.

As usual, in flavor models, distinct mixing patterns can be derived for different charge assignments. Here we focus on the leptonic sector, hence quarks are assumed to be uncharged under the symmetry. Besides, a sufficient condition to obtain an electromagnetic anomaly-free ALP is

$$\sum_i Q_{L_i} = 0, \qquad \sum_i Q_{e_i} = 0, \qquad (1)$$

where Q_{L_i} and Q_{e_i} are the charges of the left-handed and right-handed charged leptons under the $U(1)_\phi$. Choosing them as $L(1, 0, -1)$ and $e(-1, 0, 1)$, the desired Froggatt-Nielsen structure can be generated.

Two Higgs doublets are introduced with charge 0 and -2, and an additional Z_2 symmetry is imposed as in the type-X 2HDM so that the only odd fields are $H_2 \rightarrow -H_2$, $e \rightarrow -e$ and $N_R \rightarrow -N_R$ [55,56]. The Higgses, H_1 and H_2, couple to quarks and leptons, respectively. To summarize, the following particles and charges under $U(1)_\phi \times Z_2$ are considered for model I:

$$H_1(0; 1), H_2(2; -1), \phi(1; 1),$$
$$L(1, 0, -1; 1), e(-1, 0, 1; -1), N_R(0, 0, 0; -1). \quad (2)$$

From (2), it can be seen that the anomalies cancel for both the left- and right-handed sector. The most general scalar potential is

$$\begin{aligned}
V(H_1, H_2, \phi) &= m_1^2 H_1^\dagger H_1 + m_2^2 H_2^\dagger H_2 + \lambda_1 (H_1^\dagger H_1)^2 \\
&+ \lambda_2 (H_2^\dagger H_2)^2 + \lambda_3 (H_1^\dagger H_1)(H_2^\dagger H_2) \\
&+ \lambda_4 |H_1 \cdot H_2|^2 + m^2 H_1 \cdot H_2 \\
&+ \lambda(\phi^\dagger \phi - v_\phi^2)^2, \quad (3)
\end{aligned}$$

where the soft term $m^2 H_1 H_2$, with $m \sim \mathcal{O}(\text{EW})$, breaks explicitly the $U(1)_\phi \times Z_2$ symmetry and contributes to the ALP mass after the EWSB. The corresponding Yukawa terms are

$$\begin{aligned}
\mathcal{L}_Y &\supset Y_u \bar{Q} \tilde{H}_1 u + Y_d \bar{Q} H_1 d + c_{ij}^e \epsilon^{n_{ij}^e} \bar{L}_i \tilde{H}_2 e_j \\
&+ c_{ij}^\nu \epsilon^{n_{ij}^\nu} \bar{L}_i H_2 N_j + (M_R)_{ij} N_{R_i} N_{R_j}^c, \quad (4)
\end{aligned}$$

with c_{ij}^e and c_{ij}^ν $\mathcal{O}(1)$ coefficients and $n_{ij}^e = q_{L_i} - q_{e_j} + q_{H_2}$, $n_{ij}^\nu = q_{L_i} - q_{N_{R,j}} - q_{H_2}$. In model I, we have

$$n_{ij}^e = \begin{pmatrix} 4 & 3 & 2 \\ 3 & 2 & 1 \\ 2 & 1 & 0 \end{pmatrix}, \qquad n_{ij}^\nu = \begin{pmatrix} 1 & 1 & 1 \\ 2 & 2 & 2 \\ 3 & 3 & 3 \end{pmatrix}. \quad (5)$$

Once the EW symmetry is broken by the Higgs vevs, v_{H_1} and v_{H_2} with $v_{EW}^2 = v_{H_1}^2 + v_{H_2}^2 = 246$ GeV, the Dirac mass matrices are simply given by

$$M_{ij}^e = \frac{v_{H_2}}{\sqrt{2}} c_{ij}^e \epsilon^{n_{ij}^e}, \qquad M_{ij}^\nu = \frac{v_{H_2}}{\sqrt{2}} c_{ij}^\nu \epsilon^{n_{ij}^\nu}. \quad (6)$$

At leading order, the charged lepton masses are

$$\frac{m_e}{m_\tau} = \frac{(c_{12}^e - c_{23}^e)^2}{c_{23}^{e2} - 1} \epsilon^4, \qquad \frac{m_\mu}{m_\tau} = (1 - c_{23}^{e2}) \epsilon^2. \quad (7)$$

Taking $\epsilon = 0.1$, the following matrix of c_{ij}^e coefficients reproduce the correct hierarchy between generations:

$$c_{ij}^e = \begin{pmatrix} 1.0 & 1.6 & 1.0 \\ 1.6 & 1.0 & -2.7 \\ 1.0 & -2.7 & 1.0 \end{pmatrix}. \quad (8)$$

Since the tau mass is not suppressed by any additional factor, we expect $v_{H_2} \simeq \epsilon^2 v_{EW}$. For this hierarchical scenario, the mixing pattern is

$$|U_L^e|_{ij} = |U_R^e|_{ij} \approx \delta_{ij} + \epsilon^{n_{ij}^e}/\epsilon^{n_{jj}^e} \quad \text{with } i \leq j. \quad (9)$$

Then, the $e - \mu$ mixing is $\mathcal{O}(\epsilon) \sim 0.1$.

The masses of the active neutrinos are produced through the usual type-I seesaw mechanism. Notice that, in this kind of formulations, the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix can always be achieved by a proper structure of the M_R matrix [57].

After the breaking of the flavor symmetry, the flavon field can be parametrized as

$$\phi(x) = \frac{1}{\sqrt{2}}(v_\phi + s(x))e^{ia(x)/v_\phi}, \quad (10)$$

with $s(x)$ a *CP*-even scalar and $a(x)$ the ALP. If all the interactions respect the $\mathrm{U}(1)_\phi \times Z_2$ symmetry, $a(x)$ is the massless Nambu-Goldstone boson emerging after the spontaneous breaking. In our model, however, a contribution to the mass of the ALP comes from the explicit soft breaking of the symmetry in Eq. (3),

$$V \supset V_{\text{soft}} = m^2 H_1 \cdot H_2 + \text{H.c.}, \quad (11)$$

With the help of the $\mathrm{U}(1)_\phi$ symmetry, the ALP field can be removed from all the terms in the Lagrangian through a rotation $e^{iQ_f a/v_\phi}$, except for the kinetic terms and this soft-breaking term. Thus, after EWSB, the soft-breaking interaction becomes

$$V_{\text{soft}} = m^2 v_{H_1} v_{H_2} \cos(2a(x)/v_\phi). \quad (12)$$

Expanding the cosine, we obtain naturally,

$$m_a^2 \simeq 4m^2 \frac{v_{H_1} v_{H_2}}{f_a^2} \sim \mathcal{O}(1) \text{ keV}^2, \quad (13)$$

where $f_a \sim \mathcal{O}(v_\phi) \sim 10^9$ GeV, $m \sim \mathcal{O}(\text{EW})$, $v_{H_1} \sim v_{\text{EW}}$ and, as stated before, $v_{H_2} \sim \epsilon^2 v_{\text{EW}}$. Alternatively, one may assume a hidden strong sector coupled to the ALP that generates its mass. In the following, we treat the ALP mass as a free parameter and, as preferred by the Xenon1T data, it should be around the keV scale.

The interaction between the pseudo-Nambu-Goldstone boson and the charged leptons, in the mass basis, is

$$-\mathcal{L}_{ae} = i \frac{\partial_\mu a}{2f_a} \bar{e}_i \gamma^\mu (V_{ij}^e + \gamma^5 A_{ij}^e) e_j. \quad (14)$$

The axial and vector couplings in Eq. (14) are defined as[1]

$$V_{ij}^e = \frac{1}{2}(U_R^{e\dagger} x_R U_R^e + U_L^{e\dagger} x_L U_L^e), \quad (15)$$

$$A_{ij}^e = \frac{1}{2}(U_R^{e\dagger} x_R U_R^e - U_L^{e\dagger} x_L U_L^e), \quad (16)$$

with x_L and x_R the diagonal 3×3 matrices whose elements are the charged-lepton $\mathrm{U}(1)_\phi$ charges and U_L^e, U_R^e the unitary transformations that diagonalize the mass matrices.[2] In general, Eqs. (15) and (16) induce FV effects which are subject to constraints from different experiments, as it is discussed in Sec. III.

B. Model II: general mixing

In model II, we generalize the previous structure to allow for arbitrary leptonic mixings. To do this, we consider the $\mathrm{U}(1)_\phi$ global symmetry as only part of a larger flavor symmetry, $\mathcal{F}$, that will determine the Yukawa structure with the observed hierarchy among generations in the lepton sector. In this way, the $\mathrm{U}(1)_\phi$ symmetry remains flavor dependent, but masses and mixings are not fixed by the $\mathrm{U}(1)_\phi$ charges.

As an example, we use the same $\mathrm{U}(1)_\phi \times Z_2$ charges as in model I although now we can take $v_\phi/\Lambda \simeq \mathcal{O}(1)$. The scalar potential and Yukawa terms remain as in Eqs. (3) and (4) but, in this case, we highlight that the coefficients c_{ij}^e and c_{ij}^ν are NOT forced to be $\mathcal{O}(1)$.

Adjusting them, different mixing patterns can be obtained. In particular, we are interested in the case of large PMNS-like mixing for charged leptons. As a typical benchmark model, we assume that the breaking of the symmetry $\mathcal{F}$ produces charged-lepton Yukawa couplings with PMNS-like mixing in the left- and right-handed sector. The couplings with the ALP are determined by Eqs. (15) and (16), but now

$$V_{ij}^e, A_{ij}^e = \frac{1}{2} U_{\text{PMNS}}^{e\dagger}(x_R \pm x_L) U_{\text{PMNS}}^e. \quad (17)$$

Then, for example, we can deduce the size of the axial 12-coupling to be as large as $A_{12}^e \simeq 0.56$. Note that $V_{ij}^e = 0$ by construction since $x_L = -x_R$.

III. CONSTRAINTS FROM LFV AND ASTROPHYSICS

Non-universal charges of the charged leptons under the $\mathrm{U}(1)_\phi$ global symmetry, together with nontrivial rotations to the mass basis, imply FV interactions between the ALP and these fermions. The absence of the anomalous coupling between the ALP and photons at tree level makes the search for ALPs by charged-lepton flavor-violating processes specially relevant. Detailed discussions about the phenomenology of *flavorful* ALPs can be found in [58,59].

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[1]For $i = j$, we can always redefine the fields to have $V_{ii}^e = 0$ [58].

[2]In our convention: $U_L^{e\dagger} M_e U_R^e = \text{Diag}(m_e, m_\mu, m_\tau)$.

TABLE I. Limits over the axion decay constant from lepton decays. `Belle-II` limits are derived from the simulated result at `Belle` [60] by rescaling the luminosity [58].

Lepton decay	BR limit	Experiment
	Present best limits	
BR($\mu \to ea$)	$<2.6 \times 10^{-6}$	Jodidio *et al.* [61]
BR($\mu \to ea$)	$<2.1 \times 10^{-5}$	`TWIST` [62]
BR($\mu \to ea\gamma$)	$<1.1 \times 10^{-9}$	Crystal Box [63]
BR($\tau \to ea$)	$<2.7 \times 10^{-3}$	`ARGUS` [64]
BR($\tau \to \mu a$)	$<4.5 \times 10^{-3}$	`ARGUS` [64]
	Projections of running experiments	
BR($\mu \to ea$)	$<1.3 \times 10^{-7}$	MEGII-fwd [58,65]
BR($\tau \to ea$)	$<8.4 \times 10^{-6}$	`Belle-II`
BR($\tau \to \mu a$)	$<1.6 \times 10^{-5}$	`Belle-II`
	Projections of planned experiments	
BR($\mu \to ea$)	$<7.3 \times 10^{-8}$	`Mu3e` [66]

Table I collects the experimental present limits and projected sensitivities for the search of ALPs through the detection of the FV process $\ell_i \to \ell_j a$. For an ALP mass around the keV, the branching ratio for the FV transition $\ell_i \to \ell_j a$ is given by

$$\text{BR}(\ell_i \to \ell_j a) = \frac{m_{\ell_i}^3}{16\pi\Gamma(\ell_j)} \frac{|C_{ij}^e|^2}{4f_a^2} \left(1 - \frac{m_a^2}{\ell_i^2}\right)^2, \quad (18)$$

with $|C_{ij}^e|^2 = |V_{ij}^e|^2 + |A_{ij}^e|^2$. For a given model, where the interaction between the ALP and charged leptons is fixed, the bounds in Table I on the $\ell_i \to \ell_j a$ transitions can be translated into bounds on f_a. Although all of them have been inspected, the strongest limits come from $\mu \to ea$.

Regarding astrophysics bounds, interesting limits can be derived from stellar evolution. In particular, the cooling of white dwarfs [67] and red giants [68,69] impose strong constraints over the ALP interactions to matter and radiation. For massless ALP, the limits at 95% C.L. are

$$f_a \gtrsim 2.3 \times 10^9 |C_{11}^e| \text{ GeV}, \quad (19)$$

$$f_a \gtrsim 1.2 \times 10^9 |C_{11}^e| \text{ GeV}. \quad (20)$$

For ALP masses above 1 keV, the cooling rate is Boltzmann suppressed and the limits above should be rescaled by the factor $\sqrt{\xi(m_a, T)/\xi(0, T)}$, where [69]

$$\xi(m_a, T) = \frac{1}{2\pi^2} \int_{m_a}^{\infty} \frac{E^2 \sqrt{E^2 - m_a^2}}{e^{E/T} - 1}. \quad (21)$$

Finally, we comment on an additional bound from the big bang nucleosynthesis (BBN). As discussed in [70], by constraining the effective number of relativistic neutrino species during BBN, non-negligible limits can be set over

the lepton-ALP coupling that compete with those from stellar evolution. When $m_a \sim [20, 1000]$ keV, the BBN bound is indeed the strongest one. In our scenario, however, with a few keV ALP mass, stellar evolution remains the most important phenomena.

IV. RESULTS

In [2], the authors conclude that an ALP satisfying

$$A_{11}^e \simeq 10^{-13} \frac{f_a}{m_e}, \quad \text{for } m_a \in [2, 3] \text{ keV}, \quad (22)$$

can reproduce the Xenon1T signal together with some reported anomalies in stellar cooling [67,69,71]. Such possibility can be realized in the context of anomaly-free DM ALPs, provided that they constitute only around the 10% of the total DM abundance ($\Omega_a/\Omega_{\text{DM}} \sim 0.1$).

Regarding the production of this anomaly-free ALP in the early Universe, two main mechanisms can be discussed: thermal and nonthermal production. The former is due to scatterings between leptons and the Higgs particles in the thermal plasma, and would require a reheating temperature $\sim\mathcal{O}(10^2 \text{ GeV})$, together with $m_{H_2} \sim \mathcal{O}(10^2 \text{ GeV})$. Such a relative low reheating temperature would not allow for thermal leptogesis, but other baryon generation mechanisms as Affleck-Dine [72] or resonant leptogenesis [73] would still be possible in a supersymmetric scenario. However, thermally produced warm DM is subject to important constraints from structure formation that might be in tension with the ALP masses considered here. In particular, the Lyman-α limit establishes that $m_a \geq 15$ keV, if $\Omega_a = \Omega_{\text{DM}}$ [74,75]. Although this bound is expected to be relaxed for our ALPs with $\Omega_a/\Omega_{\text{DM}} \simeq 0.1$ [76], it is not straightforward to rescale this constraint with the DM fraction and we leave the detailed computation for a future work.

Nonthermal production, on the other hand, can safely evade these constraints. We can use the so-called misalignment mechanism, where the ALP starts oscillating when the Hubble parameter is comparable to its mass [2,71] or we can also consider the ALP is produced by the decay of the inflaton. In both cases, a low reheating temperature is needed to avoid the thermal production of the ALP.

The discussion in [2,71] is restricted to astrophysical and cosmological constraints while flavor observables are not discussed. Here, we aim to highlight the role of flavor observables to (dis)prove this kind of models.

Figures 1 and 2 show the Xenon1T favored prediction for f_a, based on the result in Eq. (22) (black diamond). Similarly, current and expected sensitivity from Jodidio *et al.* [61] (green continuous line) and `MU3E` [66] (red dashed line) in dedicated searches for $\mu \to e\gamma$ are displayed as a function of m_a. We also show the projection of the proposal by Calibbi *et al.* [58], MEGII-fwd (yellow dashed line), for `MEGII` [65] to improve the detection of the
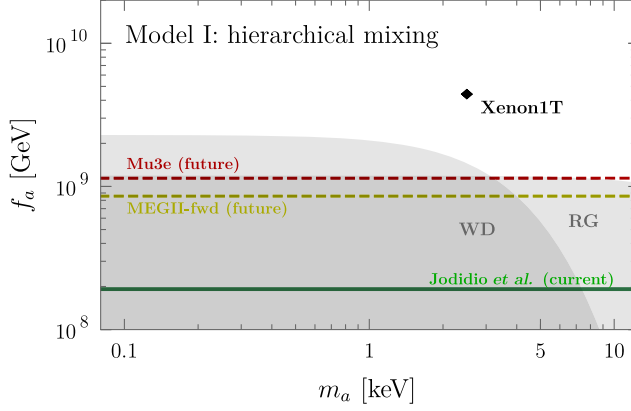
FIG. 1. Results for model I with hierarchical Yukawa couplings generated *à la* Froggatt-Nielsen.

process of interest, $\mu \to ea$. Finally, limits due to white dwarfs and red giants (gray shaded regions) also impose relevant bounds on our models [58].

From Fig. 1, we notice that testing model I (small mixing) with LFV observables remains quite challenging, even for future sensitivities. On the other hand, scenarios with larger mixing effects in the charged-lepton sector provide better prospects. For model II, in Fig. 2, we observe that while current limits are not sufficient to constrain the model, more stringent bounds coming from MU3E or the implementation of MEGII-fwd are enough to probe this formulation. One may then conclude that LFV can clearly complement astrophysics searches and, in some cases, go beyond them. Flavored ALP models provide a rich phenomenology to be investigated with present and future data.

Regarding the Higgs sector of our models, we briefly comment on its phenomenology here. A type-X 2HDM has been considered in our realizations, where H_1 couples to quarks and H_2 to leptons. In the absence of CP violation at tree level, the Higgs spectrum consist of two scalars, one pseudoscalar and two charged states. Our formulations are defined in the large $\tan\beta = v_{H_1}/v_{H_2}$ regime, where
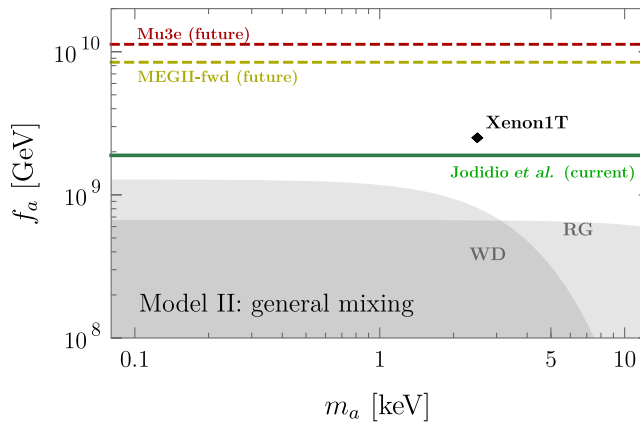

FIG. 2. Results for model II with general Yukawa matrices and mixing.

the pseudoscalar and the charged Higgses are essentially given by the leptophilic Higgs. We also assume the decoupling limit, where $\sin(\beta - \alpha) \simeq 1$ and the lightest scalar behaves as the SM Higgs and the observed mass can be reproduced with a mild tuning of the parameters in the Higgs potential.

The main constraints to this scenario are those on the mass of the charged Higgses and come from LEP, which establishes the model independent limit $M_{H^\pm} > 79.3$ GeV [77], the tau decay $\tau \to \mu \bar{\nu}\nu$ [78], the charged Higgs decay $H^\pm \to \tau^\pm \nu$ and the top decay $t \to H^+ b$, if $M_{H^+} \lesssim m_t - m_b$ [79]. The latter two bounds apply for $\tan\beta \lesssim 17$, so they are safely evaded in our case [80]. Similarly, our leptophilic Higgs states escape the strong constraints over the charged Higgs masses due to $B \to X_s \gamma$ [81] and to $B \to \tau\nu$ [82], common in type-II 2HDM [55]. The tau decay, in contrast, imposes the following limit for the considered $\tan\beta$ regime: $M_{H^\pm} \gtrsim 200$ GeV [55].

Regarding neutral states, we have a mild bound over the pseudoscalar mass, $M_A > 10$ GeV, from $B_s \to \mu\mu$ [83]. LHC searches of heavy neutral Higgses in the dilepton [84] and ditau [85,86] final state do not introduce relevant constraints in our model, mainly due to their limited H and A production in this environment with large $\tan\beta$ and small $\cos(\beta - \alpha)$, where both the gluon and weak-boson fusion production are strongly suppressed [55,87]. Therefore, present bounds on extra Higgses are easily satisfied in our model taking a scalar mass in the potential $m \gtrsim 200$ GeV, of the order of the electroweak scale.

V. CONCLUSIONS

In this paper, we considered the LFV effects from a keV scale flavor-dependent ALP which is motivated by recent Xenon1T results. We find that, for a general mixing in the lepton sector, the leptonic flavor changing experiments could confirm or exclude the possibility of explaining the Xenon1T result by an ALP, while being consistent with all phenomenological and astrophysical constraints. On the other hand, if the leptonic mixing originating from the Froggatt-Nielsen symmetry are small, CKM-like, the measurement of their LFV effects would constitute a challenge for future experiments.

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