

Axion-mediated forces and CP violation in left-right models

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We compute the CP-violating (CPV) scalar axion coupling to nucleons in the framework of baryon chiral perturbation theory and we apply the results to a class of left-right symmetric models. The correlated constraints with other CPV observables show that the predicted axion nucleon coupling is within the reach of axion-mediated force experiments as ARIADNE for M_{W_R} up to 1000 TeV.

Introduction. The axion experimental program has received an impressive boost in the last decade. Novel detection strategies, bridging distant areas of physics, promise to open for exploration the parameter space of the QCD axion in the not-so-far future, possibly addressing the issue of strong CP violation in the Standard Model (SM) via the Peccei-Quinn (PQ) mechanism [1–4] and the Dark Matter (DM) puzzle [5–7] (for updated reviews, see [8–10]). Standard axion searches often rely on highly model-dependent axion production mechanisms, as in the case of relic axions (haloscopes) or to a less extent solar axions (helioscopes); while traditional optical setups in which the axion is produced in the lab are still far from probing the standard QCD axion. A different experimental approach, as old as the axion itself [3], consists in searching for axion-mediated macroscopic forces [11]. Given the typical axion Compton wavelength $\lambda_a \sim 2$ cm ($10 \mu\text{eV}/m_a$), the presence of an even tiny *scalar* axion coupling to matter may coherently enhance the force between macroscopic bodies. The sensitivity of those experiments crucially depends on the (pseudo)scalar nature of the axion field, a matter of ultraviolet (UV) physics.

Within QCD the Vafa-Witten theorem [12] ensures that the axion vacuum expectation value (VEV) relaxes on the $\bar{\theta}_{\text{eff}} \equiv \langle a \rangle / f_a + \bar{\theta} = 0$ minimum, where $\bar{\theta}$ denotes the QCD topological term. However, extra CP violation in the UV invalidate the hypotheses of this theorem, and in general one expects a minimum with $\bar{\theta}_{\text{eff}} \neq 0$. While the CKM phase in the SM yields $\bar{\theta}_{\text{eff}} \simeq 10^{-18}$ [13], too tiny to be experimentally accessible, CPV phases from new physics can saturate the neutron Electric Dipole Moment (nEDM) bound $|\bar{\theta}_{\text{eff}}| \lesssim 10^{-10}$.

Another remarkable consequences of a non-zero $\bar{\theta}_{\text{eff}}$ is the generation of CPV scalar axion couplings to nucleons, \bar{g}_{aN} , which is probed in axion-mediated force experiments. In particular, given the nEDM bound on $\bar{\theta}_{\text{eff}}$ the scalar-pseudoscalar combination (also known as monopole-dipole interaction) offers the best chance for detecting the QCD axion. Additionally, the presence of a spin-dependent interaction allows to use Nuclear Magnetic Resonance

(NMR) to enhance the signal. This is the strategy pursued by ARIADNE [14, 15] which aims at probing the monopole-dipole force detected via a sample of nucleon spins. A similar approach is pursued by QUAX- $g_p g_s$ [16, 17], which employs instead electron spins. ARIADNE will probe $|\bar{\theta}_{\text{eff}}| \lesssim 10^{-10}$ for axion masses $1 \lesssim m_a / \mu\text{eV} \lesssim 10^4$, a range that is highly motivated by DM.

In this Letter, we provide a coherent framework for computing the CPV scalar axion coupling to nucleons in terms of new sources of CP violation beyond the SM. This is done in the framework of the baryon chiral Lagrangian so that, compared to previous works [11, 18–20], we consistently take into account the contributions of meson tadpoles (previously considered also in [19] using current algebra techniques), as well as isospin-breaking effects (previously considered in [20], although without meson tadpoles).

We then exemplify the results for a minimal PQ-extension of the Left-Right (LR) model with \mathcal{P} -parity, an extremely predictive scenario for CP violation. We follow closely the approach of Ref. [21] where a detailed study of the kaon CPV observables ε , ε' and the nEDM (d_n) in minimal LR scenarios was presented. It was shown there that the embedding of a PQ symmetry relaxes the lower bound on the LR scale just at the upper reach of the LHC. We show in this work that the search for the scalar axion coupling to nucleons provides correlated and complementary constraints, with a sensitivity to the LR scale stronger than other CPV observables. Remarkably, for a non-decoupled LR-scale we obtain a lower-bound on the \bar{g}_{aN} coupling, thus setting a target for present axion-mediated force experiments.

Axion couplings to matter. Including both CP-conserving and CPV couplings, the axion effective Lagrangian with matter fields ($f = p, n, e$) reads

$$\mathcal{L}_{af} = C_{af} \frac{\partial_\mu a}{2f_a} \bar{f} \gamma^\mu \gamma_5 f - \bar{g}_{af} a \bar{f} f, \quad (1)$$

where the first term can be rewritten in terms of a pseudoscalar density as $-g_{af} a \bar{f} i \gamma_5 f$, with $g_{af} = C_{af} m_f / f_a$. For protons and neutrons the adimen-

sional axion coupling coefficients are [22]

$$C_{ap} = -0.47(3) + 0.88(3)c_u - 0.39(2)c_d - K_a \quad (2)$$

$$C_{an} = -0.02(3) + 0.88(3)c_d - 0.39(2)c_u - K_a, \quad (3)$$

where $K_a = 0.038(5)c_s + 0.012(5)c_c + 0.009(2)c_b + 0.0035(4)c_t$, and where the (model-dependent) axion couplings to quarks c_q are defined via the Lagrangian term $c_q \frac{\partial_\mu a}{2f_a} \bar{q} \gamma^\mu \gamma_5 q$. The axion mass and decay constant are related by $m_a = 5.691(51) (10^{12} \text{ GeV}/f_a) \mu\text{eV}$ [23, 24].

The origin of the CPV scalar couplings to nucleons \bar{g}_{aN} ($N = p, n$) can be traced back to sources of either PQ or CP violation. These generically lead to a remnant $\bar{\theta}_{\text{eff}} \neq 0$ which induces CPV couplings. In the isospin symmetric limit one has [11]

$$\bar{g}_{aN} = \frac{\bar{\theta}_{\text{eff}}}{f_a} \frac{m_u m_d}{m_u + m_d} \langle N | (\bar{u}u + \bar{d}d) | N \rangle. \quad (4)$$

A shortcoming of Eq. (4) is that CPV physics can induce not only $\bar{\theta}_{\text{eff}}$, but also tadpoles for the π^0 , η_0 , η_8 meson fields, which yield extra contributions to \bar{g}_{aN} , as to other CPV observables such as d_n . A derivation of $g_{an,p}$ that takes all these effects into account at once can be obtained in the context of the baryon chiral Lagrangian with axion field. Following the approach and notation of [21] we find

$$\begin{aligned} \bar{g}_{an,p} \simeq & \frac{4B_0 m_u m_d}{f_a(m_u + m_d)} \left[\pm (b_D + b_F) \frac{\langle \pi^0 \rangle}{F_\pi} \right. \\ & + \frac{b_D - 3b_F}{\sqrt{3}} \frac{\langle \eta_8 \rangle}{F_\pi} - \sqrt{\frac{2}{3}} (3b_0 + 2b_D) \frac{\langle \eta_0 \rangle}{F_\pi} \\ & \left. - \left(b_0 + (b_D + b_F) \frac{m_{u,d}}{m_d + m_u} \right) \bar{\theta}_{\text{eff}} \right], \end{aligned} \quad (5)$$

where for clarity we neglected $m_{u,d}/m_s$ terms. Here $B_0 = m_\pi^2/(m_d + m_u)$ while the hadronic Lagrangian parameters $b_{D,F}$ are determined from the baryon octet mass splittings, $b_D \simeq 0.07 \text{ GeV}^{-1}$, $b_F \simeq -0.21 \text{ GeV}^{-1}$ at LO [25]. The value of b_0 is determined from the pion-nucleon sigma-term as $b_0 \simeq -\sigma_{\pi N}/4m_\pi^2$, and given the recent lattice results of Ref. [26] one obtains $b_0 \simeq -0.76 \pm 0.04 \text{ GeV}^{-1}$ at 90% C.L.

Eq. (5) represents the general result, including isospin-breaking effects, where $\bar{\theta}_{\text{eff}}$ and the meson VEVs are to be computed from a given source of CPV. We stress the importance of deriving them in the same formalism, given that exact cancelations can happen, as demonstrated in [21] for the case of d_n . This expression generically breaks the simple proportionality between \bar{g}_{aN} and d_n , as we shall see explicitly in the context of the PQ-extended LR model.

PQ-extended Left Right models. We begin by considering the minimal LR symmetric model [27–30], where the gauge group $\text{SU}(3)_C \times \text{SU}(2)_L \times \text{SU}(2)_R \times \text{U}(1)_{B-L}$ is spontaneously broken by a

scalar triplet VEV $\langle \Delta_R^0 \rangle = v_R$ and eventually by the VEVs of a bidoublet field $\langle \Phi \rangle = \text{diag}\{v_1, e^{i\alpha} v_2\}$, where $v^2 = v_1^2 + v_2^2 \ll v_R^2$ sets the electroweak scale and $\tan\beta \equiv t_\beta = v_2/v_1$. The single phase α is the source of the new CP violation. An important phenomenological parameter is the mixing between left and right gauge bosons, $\zeta \simeq -e^{i\alpha} \sin 2\beta M_{W_L}^2/M_{W_R}^2$. From the direct experimental limits on the LR scale one has $|\zeta| < 4 \times 10^{-4}$.

Born in order to feature the spontaneous origin of the SM parity breaking, the model is supplemented with a discrete LR exchange parity \mathcal{P} , assumed exact at high scale and broken spontaneously by v_R . \mathcal{P} exchanges the fermion representations $Q_L \leftrightarrow Q_R$, as well as the bidoublet $\Phi \leftrightarrow \Phi^\dagger$. As a result, the Yukawa Lagrangian

$$\mathcal{L}_Y = \bar{Q}_L (Y\Phi + \tilde{Y}\tilde{\Phi})Q_R + \text{h.c.}, \quad (6)$$

requires hermitean Y, \tilde{Y} . The diagonalization of the quark mass matrices gives rise to a new CKM matrix V_R in the W_R charged currents. Only for nonzero α the masses are non-hermitean and V_R departs from the standard V_L . A remarkable analytical form for V_R is given perturbatively in the small parameter $y = |s_\alpha t_{2\beta}| \lesssim 2m_b/m_t \simeq 0.05$ [31, 32]. While the left and right mixing angles can be considered equal for our purposes, the new CP external phases in V_R are relevant. For later convenience we denote them as θ_q , with $V_R = \text{diag}\{e^{i\theta_u}, e^{i\theta_c}, e^{i\theta_t}\} V_L \text{diag}\{e^{i\theta_d}, e^{i\theta_s}, e^{i\theta_b}\}$. Using the analytical form, one finds that all θ_q are small deviations of $O(y)$ around 0 or π , corresponding to 32 physically different sign combinations of the quark mass eigenvalues [21, 32]. For details on the relevant features of the minimal LR model we refer to [21, 33] and references therein.

There are two qualitatively different ways of implementing a $\text{U}(1)_{\text{PQ}}$ symmetry in LR models, following either the KSVZ [34, 35] or the DFSZ [36, 37] variant. The former is the most straightforward, as the field content of the minimal LR symmetric model remains uncharged under $\text{U}(1)_{\text{PQ}}$, and the pseudoscalar axion couplings to nucleons are given by Eqs. (2)–(3) with $c_q = 0$.

On the other hand, the construction of a LR-DFSZ model, with SM quarks carrying PQ charges, turns out to be less trivial [38]. This is due to a two-fold reason: *i*) the identification of the physical axion field requires the PQ symmetry to be orthogonal to the broken LR generators¹ and *ii*) chiral PQ charges $\mathcal{X}_{Q_L} \neq \mathcal{X}_{Q_R}$ forbid one of the Yukawa terms in Eq. (6), implying unphysical mass matrices. Hence, either the LR field content must be extended (e.g. with a second bidoublet Φ' [41]) or

¹ Existing works on LR-DFSZ models (see e.g. [39–42]) do not carry out the identification of the low-energy axion field, a necessary ingredient in order to compute its couplings to matter. This issue was recently addressed in the context of $\text{SO}(10) \times \text{U}(1)_{\text{PQ}}$ [43].

effective operators must be invoked in the Yukawa sector [42]. Finally, a complex singlet \mathcal{S} to decouple the PQ scale from v_R and v is needed. A consistent LR-DFSZ model can be constructed [38] and the axion couplings to quarks and charged-leptons are found to be

$$c_{u,c,t} = \frac{1}{3} \sin^2 \beta, \quad c_{d,s,b} = c_{e,\mu,\tau} = \frac{1}{3} \cos^2 \beta. \quad (7)$$

Here, $t_\beta^2 = (v_2^2 + v_2'^2)/(v_1^2 + v_1'^2)$ for renormalizable Yukawa terms involving two bidoublets, while in the case of one bidoublet (and effective Yukawas) the standard definition of t_β applies.

While the minimal LR model with \mathcal{P} is a predictive theory even in the strong CP sector [44, 45], the axion hypothesis can relax predictivity in the fermion as well as in the strong CP sector, if other fields as a second bidoublet are introduced. We stick below to the LR-KSVZ or the LR-DFSZ case with a single bidoublet and a nonrenormalizable Yukawa term. The axion washes out $\bar{\theta}$ (and renormalizations [44, 46]), and observables such as e.g. d_n and $\bar{g}_{an,p}$, are tightly predicted.

Important remarks then apply. Quark masses set as usual a perturbativity limit on t_β , mainly due to m_t/m_b : one finds $t_\beta \lesssim 0.5$ [47] or $\gtrsim 2$. The two ranges are equivalent in the minimal model (swapping Y and \bar{Y}) but they become physically different when the PQ symmetry acts on Φ . Within this perturbative domain the pseudoscalar axion coupling to nucleons Eqs. (2)–(3) can never be zero, the LR-DFSZ setup differing from the SM-DFSZ one where C_{an} vanishes for $t_\beta \simeq 1.35$, within its perturbative range [9]. Moreover, in LR-DFSZ the axion coupling to electrons in Eq. (7) can vanish for $\beta \rightarrow \pi/2$, with non-trivial implications for astrophysical limits.

CPV scalar axion couplings. After integrating out the LR interactions, a complete set of effective quark operators that violate CP is generated. In the baryon chiral Lagrangian they induce pion tadpoles and lead to CPV pion-nucleon couplings. Among all operators, $\mathcal{O}_1^{ud} = (\bar{u}u) (\bar{d}i\gamma_5 d)$ [48, 49] generates the leading contribution to d_n . We show here that it also generates the dominant contribution to $\bar{g}_{ap,n}$.

For the complete short-distance Lagrangian and the extended chiral Lagrangian in the presence of the effective operators we refer to [21]. Their Wilson coefficients, renormalized at the hadronic scale (1 GeV), are matched with their chiral representation. It is here sufficient to recall that the short-distance coefficients $C_i^{qq'}$ depend on the relevant CKM entries, which include the additional CP phases of V_R , and on the LR gauge mixing ζ .

In the spontaneously broken PQ phase, by an appropriate axion-dependent chiral rotation of the u , d , s quarks the topological term plus the axion field can be properly absorbed into the LR extended meson Lagrangian [21, 49]. One then computes the shift of the chiral vacuum for the meson and axion fields, due to the presence of the CPV operators.

When only the leading $\mathcal{C}_1^{ud} - \mathcal{C}_1^{du} \equiv \mathcal{C}_1^{[u,d]}$ term is considered, the induced VEVs are readily obtained [21, 48, 49],

$$\begin{aligned} \frac{\langle \pi^0 \rangle}{F_\pi} &\simeq \frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[u,d]} \frac{c_3}{B_0 F_\pi^2} \frac{m_u + m_d + 4m_s}{m_u m_d + m_d m_s + m_s m_u} \\ \frac{\langle \eta_8 \rangle}{F_\pi} &\simeq \frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[u,d]} \frac{\sqrt{3}c_3}{B_0 F_\pi^2} \frac{m_d - m_u}{m_u m_d + m_d m_s + m_s m_u} \\ \bar{\theta}_{\text{eff}} &\simeq \frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[u,d]} \frac{2c_3}{B_0 F_\pi^2} \frac{m_d - m_u}{m_u m_d}, \end{aligned} \quad (8)$$

with $\langle \eta_0 \rangle = 0$. The axion no longer cancels the original $\bar{\theta}$ term, leaving a calculable $\bar{\theta}_{\text{eff}}$. As expected, the pion VEV is isospin odd ($u \leftrightarrow d$), while the other VEVs are even. The short distance coefficients $\mathcal{C}_1^{qq'}$ are listed in [21], and the unknown low-energy constant c_3 is estimated in the large N limit as $c_3 \sim F_\pi^4 B_0^2/4$. Analogously, for the doubly Cabibbo suppressed \mathcal{O}_1^{us} operator we find

$$\begin{aligned} \frac{\langle \pi^0 \rangle}{F_\pi} &\simeq \frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[u,s]} \frac{c_3}{B_0 F_\pi^2} \frac{2m_d + 2m_s - m_u}{m_u m_d + m_d m_s + m_s m_u} \\ \frac{\langle \eta_8 \rangle}{F_\pi} &\simeq \frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[u,s]} \frac{\sqrt{3}c_3}{B_0 F_\pi^2} \frac{2m_d + m_u}{m_u m_d + m_d m_s + m_s m_u} \\ \bar{\theta}_{\text{eff}} &\simeq \frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[u,s]} \frac{2c_3}{B_0 F_\pi^2} \frac{m_s - m_u}{m_u m_s}. \end{aligned} \quad (9)$$

One notices in both Eqs. (8)–(9) the m_s/m_d enhancement of $\langle \pi^0 \rangle$ over the other meson VEV.

It was observed in [21] that computing via Eq. (8) the CPV couplings, $\bar{g}_{n,p\pi}$ and $\bar{g}_{n\Sigma^- K^+}$, the former vanishes identically. On the other hand, when \mathcal{O}_1^{us} is considered, $\bar{g}_{n\Sigma^- K^+}$ cancels in turn. In either case the meson VEVs cancel exactly the contribution of $\bar{\theta}_{\text{eff}}$. We double checked this result using the basis of Ref. [25], which makes such a cancellation transparent. This feature suppresses the pion loop contribution of \mathcal{O}_1^{ud} to the nEDM and makes the contribution of \mathcal{O}_1^{us} numerically relevant, albeit still subleading. The destructive interference among the two operators weakens the strength of the LR contribution to the nEDM.

Such a cancellation is not present for the CPV axion-nucleon couplings $\bar{g}_{an,p}$, obtained via Eq. (5) using (8)–(9). The result, dominated by \mathcal{O}_1^{ud} , is

$$\begin{aligned} \bar{g}_{an,p} &\simeq -\frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[u,d]} \frac{8c_3}{F_\pi^2 f_a (m_d + m_u)} \\ &\times \left[b_0 (m_d - m_u) \mp (b_D + b_F) m_{u,d} \right]. \end{aligned} \quad (10)$$

A few comments on Eqs. (5) and (10) are in order. The chiral approach detailed in Ref. [21] allows us to consistently derive and account for the meson and axion tadpole contributions, thus properly addressing interference and comparison among the various contributions. It further includes LO isospin-breaking effects that enter through the pion VEV (via the $b_{D,F}$ couplings) and from the $\bar{\theta}_{\text{eff}}$ term (due to the quark mass dependence of the chiral rotation

that absorbed the topological term). Within the range of values of the hadronic parameters here considered it leads to a \bar{g}_{ap} coupling about 60% larger than \bar{g}_{an} . Finally, the results in Eqs. (5)–(10) are general enough to apply to any PQ completion of the effective LR scenario, since the model-dependent derivative axion couplings do not enter.

Axion-mediated forces. Monopole-dipole forces turn out to be the best combination to test the QCD axion. In fact, monopole-monopole interactions are doubly suppressed in $\bar{\theta}_{\text{eff}}$, while dipole-dipole forces have large backgrounds from ordinary magnetic forces. State of the art limits on monopole-dipole forces can be found in Ref. [50]. The resulting lower bounds are at most at the level of $f_a \gtrsim \sqrt{\bar{\theta}_{\text{eff}}} 10^{13}$ GeV. Better constraints are actually obtained by combining limits on monopole-monopole interactions with astrophysical limits of pseudoscalar couplings [51].

A new detection concept, by the ARIADNE collaboration [14, 15], plans to use NMR techniques to probe the axion field sourced by unpolarized Tungsten ^{184}W and detected by laser-polarized ^3He . In its current version, the experiment is sensitive to $\bar{g}_{a^{184}\text{W}} g_{a^3\text{He}}$. The CPV coupling axion coupling to Tungsten is approximated by $\bar{g}_{a^{184}\text{W}} \simeq 74(\bar{g}_{ap} + \bar{g}_{ae}) + 110\bar{g}_{an}$ [10], where for the QCD axion $\bar{g}_{ae} = 0$. It is convenient to define an average coupling to nucleons (weighting isospin breaking) as

$$\bar{g}_{aN} \equiv \frac{74\bar{g}_{ap} + 110\bar{g}_{an}}{184}. \quad (11)$$

The CP-conserving term, $g_{a^3\text{He}} = g_{an}$, is only sensitive to neutrons because protons and electrons are paired in the detection sample. Thanks to NMR, ARIADNE can improve the sensitivity of previous searches and astrophysical limits by up to two orders of magnitude in $(\bar{g}_{aN}g_{an})^{1/2}$ (for $m_a \in [1, 10^4] \mu\text{eV}$ depending on the spin relaxation time), before passing to a scaled-up version with a larger ^3He cell reaching liquid density.

CPV probes of LR scale. To analyze the predicted $(\bar{g}_{aN}g_{an})^{1/2}$ as a function of M_{WR} , we study the four CPV observables (ε , ε' , d_n , \bar{g}_{aN}), while marginalizing on $\tan\beta$, the CP phase α , and the sign combinations. As in Refs. [21, 52], we introduce a parameter h_i for each observable, normalizing the LR contributions to the experimental central value (ε , ε') or upper bound (d_n). For the latter we take the updated 90% C.L. result $d_n < 1.8 \times 10^{-26}$ ecm [53]. The LR contributions to the indirect CPV parameter ε in kaon mixing was thoroughly analyzed in [52] to which we refer for details. As for the direct CPV parameter ε' , the latest lattice result [54] for the $K \rightarrow \pi\pi$ matrix element of the leading QCD penguin operator supports the early chiral quark model prediction [55, 56], confirmed by the resummation of the pion rescattering [57], as well as more recent chiral Lagrangian

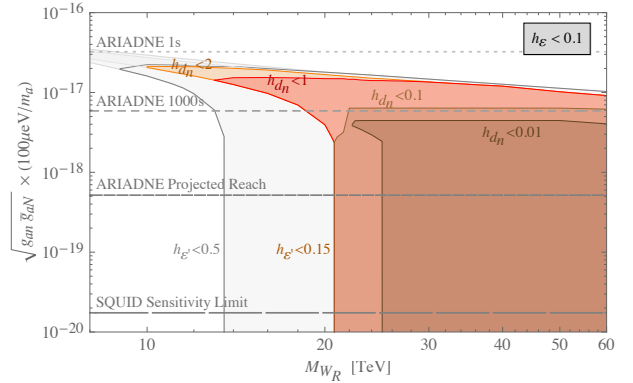


FIG. 1. Regions in the LR-DFSZ model of the CPV axion nucleon coupling probed by ARIADNE.

reassessments [58, 59], including a detailed analysis of isospin breaking. All of the above point to a SM prediction in the ballpark of the experimental value, albeit with a large error [60]. We consider below two benchmark cases: 50% and 15% of ε' accounted for by LR physics [61, 62]. The model-dependent pseudoscalar coupling g_{an} in the monopole-dipole interaction is taken for the case of the LR-DFSZ setup. Similar results are obtained for LR-KSVZ, for which however g_{an} is compatible with zero, Eq. (3).

The average nucleon coupling in Eq. (11) is computed using Eq. (10). With the updated d_n bound and including the strange quark contributions, we obtain

$$\begin{aligned} \bar{g}_{aN} &= \frac{|\zeta|}{10^{-5}} \left[6.4 \sin \alpha_{ud} + 0.7 \sin \alpha_{us} \right] \frac{m_a}{100 \mu\text{eV}} 10^{-12} \\ h_{d_n} &= \frac{|\zeta|}{10^{-5}} \left[7.1 \sin \alpha_{ud} - 3.4 \sin \alpha_{us} \right] \\ h_{\varepsilon'} &= \frac{|\zeta|}{10^{-5}} \left[9.2 \sin \alpha_{ud} + 9.2 \sin \alpha_{us} \right], \end{aligned} \quad (12)$$

where $\alpha_{qq'} = \alpha - \theta_q - \theta_{q'}$. We recall that all phases θ_q depend on a single parameter. Also, $\alpha_{ud} \simeq \alpha_{us}$ modulo π for $M_{WR} \lesssim 30$ TeV from the h_ε constraint [52]. There is clearly a tight correlation between the above observables.

The allowed regions of $(\bar{g}_{aN}g_{an})^{1/2}$ as a function of M_{WR} are shown in Fig. 1, together with the reach of three different phases of ARIADNE (1s, 1000s, projected) [14, 15] and the SQUID sensitivity limit. We scale the coupling combination by $f_a \sim 1/m_a$, making the prediction independent from it. With this normalization the experiment sensitivities vary mildly with m_a . In the plot we show their best reach, attained for $m_a \sim 10^{2 \pm 3} \mu\text{eV}$. Present limits from astrophysics [51] and monopole-dipole experiments [50] lie above the plot and are hence ineffective to probe the LR scale.

The predicted regions depend on the constraints on h_ε , $h_{\varepsilon'}$ and h_{d_n} . In the colored area the LR contribution to ε' is allowed up to 15%, while in light gray one it is extended to 50%, given the present theoretical uncertainties. In both cases, a lower bound

on \bar{g}_{aN} arises, for $M_{W_R} \lesssim 20$ or 13 TeV respectively. For $m_a \approx 100 \mu\text{eV}$ a positive detection from ARIADNE below 2×10^{-18} would falsify the LR-DFSZ scenario below these scales. Either a rejection of the LR-DFSZ model or a sharp upper bound on M_{W_R} arises for a measurement above 10^{-17} . Given the probed combination, the upper boundary of the shaded region decreases as $1/M_{W_R}$ so that within the ARIADNE sensitivity the model provides possible signals up to $M_{W_R} \sim 1000$ TeV. Standard flavour observables, which decouple as $1/M_{W_R}^2$, have a more limited reach.

The effect of the present and future constraints on d_n are shown with increasingly darker shadings, from a most conservative $h_{d_n} < 2$ (accounting for hadronic uncertainties), to a most stringent future bound of $h_{d_n} < 0.01$. The bounds on d_n limit from above the predicted axion-mediated force. For instance $h_{d_n} < 0.1$ implies a prediction at the level of the ARIADNE 1000s sensitivity.

To conclude, we showed that axion-mediated forces provide a powerful probe of the CPV structure and scale of minimal LR-PQ scenarios. It is amusing that the first hints of high-energy parity restoration may possibly be revealed in a condensed matter lab.

ACKNOWLEDGMENTS

The work of LDL is supported by the Marie Skłodowska-Curie Individual Fellowship grant AXIONRUSH (GA 840791).

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