Axion-Mediated Forces, *CP* Violation, and Left-Right Interactions

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We compute the *CP*-violating scalar axion coupling to nucleons in the framework of baryon chiral perturbation theory and we apply the results to the case of left-right symmetry. The correlated constraints with other *CP*-violating observables show that the predicted axion nucleon coupling is within the reach of present axion-mediated force experiments for M_{W_R} up to 1000 TeV.

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Introduction.—The axion experimental program has received an impressive boost in the past 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 Refs. [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 $\lambda_a \sim 2\ \text{cm}\,(10\ \mu\text{eV}/m_a)$, an even tiny *scalar* axion coupling to matter may coherently enhance the force between macroscopic bodies. The sensitivity of these 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 Cabibbo-Kobayashi-Maskawa (CKM) phase in the SM yields $\bar{\theta}_{\text{eff}} \simeq 10^{-18}$ [13], too tiny to be experimentally accessible, *CP*-violating (CPV) phases from new physics can saturate the neutron electric dipole moment (nEDM) bound $|\bar{\theta}_{\text{eff}}| \lesssim 10^{-10}$.

Another remarkable consequence of a nonzero $\bar{\theta}_{\text{eff}}$ 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 us to use nuclear magnetic resonance (NMR) to enhance the signal. This is the strategy pursued by the ARIADNE experiment [14,15], which aims at probing the monopole-dipole force via a sample of nucleon spins. A similar approach is pursued by QUAX-$g_p g_s$ [16,17], using instead electron spins. ARIADNE will probe $|\bar{\theta}_{\text{eff}}|$ below 10^{-10} for axion masses $1 \lesssim m_a/\mu\text{eV} \lesssim 10^4$, a range 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 that allows us to compute all contributions of meson tadpoles and $\bar{\theta}_{\text{eff}}$ at once, as well as isospin-breaking effects. In comparison to previous works [11,18–20], the contributions of the pion tadpole induced by the QCD dipole operator was estimated in Ref. [18] by naive dimensional analysis and in Ref. [19] using current algebra techniques, while isospin breaking was considered in Ref. [20] for $\bar{\theta}_{\text{eff}}$ without meson tadpoles. Our result is general and can be systematically applied to any bosonic representation of *P*- and *CP*-violating effective operators induced in extensions of the SM.

We detail our approach in the case of effective operators from right-handed (RH) currents, and then apply the results in the minimal left-right symmetric model (LRSM) endowed with a PQ symmetry and $\mathcal{P}$ parity as LR symmetry. This is an extremely predictive and motivated case for neutrino masses and additional CP violation, with an active collider physics program [21]. We build on the approach detailed in Ref. [22], which presented a study of the kaon CPV observables ε, ε' and the nEDM (d_n) in minimal LR scenarios. It was found there that the embedding of a PQ symmetry relaxes the lower bound on the LR scale just at the upper reach of the LHC. In this work we show that the present 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 nondecoupled LR scale we obtain a lower bound on the $\bar{g}_{aN}$ coupling, thus setting a target for present axion-mediated force experiments.

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

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

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 adimensional axion coupling coefficients are [23]

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

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

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 $c_q(\partial_\mu a/2f_a)\bar{q}\gamma^\mu\gamma_5 q$. The axion mass and decay constant are related by $m_a = 5.691(51)(10^{12}\,\mathrm{GeV}/f_a)\,\mu\mathrm{eV}$ [24,25].

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 $\bar{\theta}_{\rm eff} \neq 0$ which induces CPV couplings. One finds for the isospin singlet component of the matrix element [11]

$$\bar{g}_{aN} = \frac{\bar{\theta}_{\rm eff}}{f_a}\frac{m_u m_d}{m_u + m_d}\frac{\langle N|\bar{u}u + \bar{d}d|N\rangle}{2}, \tag{4}$$

where we included a $1/2$ factor missed in Ref. [11]. A shortcoming of Eq. (4) is that CPV physics can induce not only $\bar{\theta}_{\rm eff}$, but also shifts the chiral vacuum, inducing tadpoles for the π^0, η_0, η_8 meson fields. These in turn yield extra contributions to $\bar{g}_{aN}$, as to other CPV observables such as d_n. A derivation of $g_{an,p}$ taking all these effects consistently into account is here obtained in the context of the baryon chiral Lagrangian with axion field, as described below. We find

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

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, $b_D \simeq 0.07\ \mathrm{GeV}^{-1}$, $b_F \simeq -0.21\ \mathrm{GeV}^{-1}$ at the leading order (LO) [26]. The value of b_0 is determined from the pion-nucleon σ-term as $b_0 \simeq -\sigma_{\pi N}/4m_\pi^2$. From the precise determination in Refs. [27,28], one obtains $b_0 \simeq -0.76 \pm 0.04\ \mathrm{GeV}^{-1}$ at 90% C.L. Given $\sigma_{\pi N} \equiv \langle N|\bar{u}u + \bar{d}d|N\rangle(m_u + m_d)/2$, the isospin symmetric $b_0\bar{\theta}_{\rm eff}$ term reproduces exactly Eq. (4).

Equation (5) represents our general result, including isospin-breaking effects, where $\bar{\theta}_{\rm eff}$ and the meson VEVs are meant to be computed from a given source of CPV. In general $\bar{g}_{aN}$ and d_n are not proportional, as it would follow from Eq. (4). Exact cancellations among the VEVs can happen for d_n [22,29].

Axion coupling and RH currents.—As a paradigmatic application, we explicitly compute the above CPV axion-matter coupling in the case of RH currents, which arise in a wide class of models beyond the SM. Heavy RH currents lead generally to four quark operators that violate P and CP as $\mathcal{O}_1^{qq'} = (\bar{q}q)(\bar{q}'i\gamma_5 q')$, $q = u,\, d,\, s$ [22,29–32]. Such operators induce meson tadpoles and allow for a nonvanishing correlator with the topological $G\tilde{G}$ term, thus shifting both chiral and axion vacua [19]. At the leading order in momentum expansion the operators $\mathcal{O}_1^{qq'}$ are represented in the low-energy meson Lagrangian by combinations of $[U^\dagger]_{qq}[U]_{q'q'}$ terms, where the usual 3×3 matrix U represents nonlinearly the meson nonet under $U(3)_L \times U(3)_R$ rotations. By a proper $U(3)_A$ field rotation, the axion field is also included in the meson and baryon chiral Lagrangians. Complete notation and details are found in Appendix D of Ref. [22]. Rotating away the axion and meson tadpoles, the new CPV axion-nucleon scalar couplings of Eq. (5) are induced from the baryon Lagrangian.

In LR effective setups the operator $\mathcal{O}_1^{ud}$ generates typically the leading contribution to d_n. We show in this work that it also generates the dominant contribution to $\bar{g}_{ap,n}$. We denote its low scale Wilson coefficient as C_1^{ud}, and similarly for other flavors. When O_1^{ud} is considered, we find [22,30,32]

$$\frac{\langle\pi^0\rangle}{F_\pi} \simeq \frac{G_F}{\sqrt{2}}C_1^{[ud]}\frac{c_3}{B_0F_\pi^2}\frac{m_u + m_d + 4m_s}{m_u m_d + m_d m_s + m_s m_u},$$
$$\frac{\langle\eta_8\rangle}{F_\pi} \simeq \frac{G_F}{\sqrt{2}}C_1^{[ud]}\frac{\sqrt{3}c_3}{B_0F_\pi^2}\frac{m_d - m_u}{m_u m_d + m_d m_s + m_s m_u},$$
$$\bar{\theta}_{\rm eff} \simeq \frac{G_F}{\sqrt{2}}C_1^{[ud]}\frac{2c_3}{B_0F_\pi^2}\frac{m_d - m_u}{m_u m_d}, \tag{6}$$

where $\mathcal{C}_1^{[ud]} \equiv \mathcal{C}_1^{ud} - \mathcal{C}_1^{du}$ and $\langle \eta_0 \rangle = 0$. The axion VEV 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 low-energy constant c_3 is estimated in the large N limit as $c_3 \sim F_\pi^4 B_0^2/4$. Another estimate, based on SU(3) chiral symmetry, is given in Ref. [29]. Analogously, for $\mathcal{O}_1^{us}$ we find

$$\frac{\langle \pi^0 \rangle}{F_\pi} \simeq \frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[us]} \frac{c_3}{B_0 F_\pi^2} \frac{2m_d + 2m_s - m_u}{m_u m_d + m_d m_s + m_s m_u},$$
$$\frac{\langle \eta_8 \rangle}{F_\pi} \simeq \frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[us]} \frac{\sqrt{3} c_3}{B_0 F_\pi^2} \frac{2m_d + m_u}{m_u m_d + m_d m_s + m_s m_u},$$
$$\bar{\theta}_{\text{eff}} \simeq \frac{G_F}{\sqrt{2}} \mathcal{C}_1^{[us]} \frac{2c_3}{B_0 F_\pi^2} \frac{m_s - m_u}{m_u m_s}. \quad (7)$$

One notices in both Eqs. (6) and (7) the m_s/m_d enhancement of $\langle \pi^0 \rangle$ over the other meson VEV.

As observed in Refs. [22,29], the CPV coupling $\bar{g}_{np\pi}$ computed using the VEVs (6) 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 against $\bar{\theta}_{\text{eff}}$, a result which is made transparent in the basis of Ref. [26].

Such a cancellation is not present for the CPV axion-nucleon couplings $\bar{g}_{an,p}$, obtained via Eq. (5) using Eqs. (6) and (7), so that the typically unsuppressed $\mathcal{O}_1^{ud}$ operator dominates. In the large m_s limit the complete result can be written as

$$\bar{g}_{an,p} \simeq -\frac{G_F}{\sqrt{2}} \frac{8 c_3 b_0}{F_\pi^2 f_a (m_d + m_u)} \times \begin{cases} m_d(\mathcal{C}_1^{[ud]} + \mathcal{C}_1^{[us]}) - m_u \mathcal{C}_1^{[ud]} b \\ m_d(\mathcal{C}_1^{[ud]} + \mathcal{C}_1^{[us]}) b - m_u \mathcal{C}_1^{[ud]}, \end{cases} \quad (8)$$

where $b = (b_0 + b_D + b_F)/b_0 \simeq 1.2$. A few comments on Eqs. (5) and (8) are in order. The chiral approach 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. Within the range of 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)–(8) are general enough to apply to any axion model with effective RH currents, since the model-dependent derivative axion couplings do not enter the scalar coupling.

Experimental probes for $\bar{g}_{an,p}$.—At present, the best sensitivity on the QCD axion exploiting axion-mediated forces is obtained by combining limits on monopole-monopole interactions with astrophysical limits of pseudo-scalar couplings [33]. On the other hand, monopole-dipole forces will become the best constraining combination in laboratory experiments. 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. [34]: the resulting lower bounds are at most at the level of $f_a \gtrsim \sqrt{\bar{\theta}_{\text{eff}}}\, 10^{13}$ GeV.

A new detection concept by Arvanitaki and Geraci [14], exploited by the ARIADNE Collaboration [15], plans to use NMR techniques to probe the axion field sourced by unpolarized tungsten ^{184}W and detected by laser-polarized ^{3}He. 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 $\bar{g}_{ae} = 0$ at tree level. It is convenient to define an average coupling to nucleons (weighting isospin breaking) as

$$\bar{g}_{aN} \equiv \frac{74\bar{g}_{ap} + 110\bar{g}_{an}}{184}. \quad (9)$$

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 2 orders of magnitude in $(\bar{g}_{aN} g_{an})^{1/2}$ (for $m_a \in [1, 10^4]$ μeV depending on the spin relaxation time), before passing to a scaled-up version with a larger ^{3}He cell reaching liquid density.

To provide an example of the testing power of these future experiments, as a definite model of RH currents we consider the paradigmatic case of the LR symmetric model, with a PQ symmetry.

Application to left-right models.—In the minimal LRSM [35–39], the gauge group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ 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, $\zeta \simeq -e^{i\alpha} \sin 2\beta M_{W_L}^2/M_{W_R}^2$, bound to $|\zeta| < 4 \times 10^{-4}$ from direct search limits on W_R.

Born in order to feature the spontaneous origin of the SM parity breaking, the model is endowed with the discrete parity $\mathcal{P}$, assumed exact at high scale and broken spontaneously by v_R. $\mathcal{P}$ exchanges the gauge groups, the fermion representations $Q_L \leftrightarrow Q_R$, and conjugates the bidoublet $\Phi \leftrightarrow \Phi^\dagger$. As a result, the Yukawa Lagrangian $\mathcal{L}_Y = \bar{Q}_L(Y\Phi + \tilde{Y}\tilde{\Phi})Q_R + \text{H.c.}$ requires Hermitian Y, $\tilde{Y}$. The diagonalization of quark masses gives rise to a new CKM matrix V_R in the W_R charged currents. Only for nonzero α the masses are non-Hermitian and V_R departs from the standard V_L. An analytical form for V_R is found perturbatively in the small parameter $y = |s_\alpha t_{2\beta}| \lesssim 2m_b/m_t \simeq 0.05$ [40,41]. While the left and right mixing angles can be considered equal for our aims, V_R has new external *CP* phases. For later convenience we denote them as θ_q, with

$V_R = \mathrm{diag}\{e^{i\theta_u}, e^{i\theta_c}, e^{i\theta_t}\} V_L \mathrm{diag}\{e^{i\theta_d}, e^{i\theta_s}, e^{i\theta_b}\}$. 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 [22,41]. For details on the relevant features of the minimal LR model, we refer to Refs. [21,22] and references therein.

There are two qualitatively different ways of implementing a $U(1)_{PQ}$ symmetry in LR models, following either the Kim-Shifman-Vainshtein-Zakharov (KSVZ) [42,43] or the Dine-Fischler-Srednicki-Zhitnitsky (DFSZ) [44,45] variant. In the former, the field content of the minimal LRSM remains uncharged under $U(1)_{PQ}$, and the pseudoscalar axion couplings to nucleons are given by Eqs. (2) and (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. This is due mainly to the fact that chiral PQ charges $\mathcal{X}_{Q_L} \neq \mathcal{X}_{Q_R}$ forbid one of the Yukawa terms in $\mathcal{L}_Y$, implying unphysical mass matrices. Hence, either the LR field content must be extended [46,47] (e.g., with a second bidoublet) or effective operators must be invoked in the Yukawa sector [48,49]. Finally, a complex singlet $\mathcal{S}$ to decouple the PQ scale from v_R and v is needed. A complete ultraviolet LR DFSZ model description is not needed here [50]; it is enough to report the axion couplings to quarks and charged-leptons:

$$c_{u,c,t} = \frac{1}{3}\sin^2\beta, \qquad c_{d,s,b} = c_{e,\mu,\tau} = \frac{1}{3}\cos^2\beta. \tag{10}$$

While the minimal LR model with $\mathcal{P}$ is a predictive theory even in the strong CP sector [51,52], 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. Below we stick 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 [51,53]), and observables such as, e.g., d_n and $\bar{g}_{an,p}$, are tightly predicted.

With this choice, quark masses set as usual a perturbativity limit on t_β, mainly due to m_t/m_b: one finds $t_\beta \lesssim 0.5$ [54] or $\gtrsim 2$. The two ranges are equivalent in the minimal model (swapping Y and $\tilde{Y}$), but they become physically different when the PQ symmetry acts on Φ. Within this perturbative domain the pseudoscalar axion coupling to nucleons Eqs. (2) and (3) can never vanish.

Axion and CPV probes of LR scale.—The RH currents in the LRSM induce the axion couplings described above. For details on the LRSM short-distance and the extended chiral Lagrangian, we refer to Ref. [22]. We just recall that the short-distance coefficients $\mathcal{C}_i^{qq'}$ depend on the relevant CKM entries, carrying the additional CP phases of V_R, and on the LR gauge mixing ζ. The $\mathcal{C}_i^{qq'}$ are renormalized at the 1 GeV hadronic scale and matched with the chiral low-energy constants.

To analyze the predicted $(g_{an}\bar{g}_{aN})^{1/2}$ as a function of M_{W_R}, we study together the four CPV observables $(\varepsilon, \varepsilon', d_n, \bar{g}_{aN})$, while marginalizing on $\tan\beta$, the CP phase α, and the 32 signs. As in Refs. [22,55], 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 $d_n < 1.8 \times 10^{-26}\ e\,\mathrm{cm}$ [56]. The LR contributions to the indirect CPV parameter ε in kaon mixing was thoroughly analyzed in Ref. [55], to which we refer the reader for details. For the direct CPV parameter ε' the latest lattice result [57] for the $K \to \pi\pi$ matrix element of the leading QCD penguin operator supports the early chiral quark model prediction [58,59], confirmed by the resummation of the pion rescattering [60], as well as more recent chiral Lagrangian reassessments [61,62], 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 [63]. We consider below two benchmark cases: 50% and 15% of ε' induced by LR physics [64,65].

The average CPV nucleon coupling in Eq. (9) is computed using Eq. (8). With the updated d_n bound and including the strange quark contributions, we obtain

$$\begin{aligned} \bar{g}_{aN} &= \frac{|\zeta|}{10^{-5}}[6.4\sin\alpha_{ud} + 0.7\sin\alpha_{us}]\frac{m_a}{100\ \mu\mathrm{eV}}10^{-12}, \\ h_{d_n} &= \frac{|\zeta|}{10^{-5}}[7.1\sin\alpha_{ud} - 3.4\sin\alpha_{us}], \\ h_{\varepsilon'} &= \frac{|\zeta|}{10^{-5}}[9.2\sin\alpha_{ud} + 9.2\sin\alpha_{us}], \end{aligned} \tag{11}$$

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 $M_{W_R} \lesssim 30$ TeV from the h_ε constraint [55], which plays an important role in enforcing a tight correlation between the above observables. The subleading role of the Cabibbo suppressed us Wilson coefficient in $\bar{g}_{aN}$ is clear, unlike the case of d_n where the leading ud contribution is canceled as mentioned above [22].

The model-dependent pseudoscalar coupling g_{an} in the monopole-dipole interaction is taken for the LR DFSZ case via Eq. (10). Similar results are obtained for LR KSVZ, for which, however, g_{an} is compatible with zero; see Eq. (3).

In Fig. 1 we show the allowed regions of $(g_{an}\bar{g}_{aN})^{1/2}$ as a function of M_{W_R}, together with the reach of three phases of ARIADNE (1 s, 1000 s, projected) [14,15] and the SQUID sensitivity limit. We scale the coupling combination by $f_a \propto 1/m_a$, making the prediction independent from it. With this normalization the experiment sensitivities vary mildly with m_a, and we show their best reach, attained for $m_a \sim 10^{2-3}\ \mu$eV. Present limits from astrophysics [33] and monopole-dipole experiments [34] 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 we relax it to 50%,

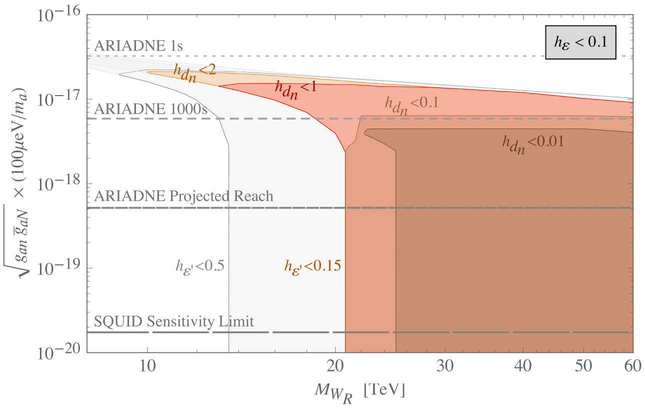

FIG. 1. Regions in the LR DFSZ model of the CPV axion nucleon coupling probed by ARIADNE.

given the present theoretical uncertainties. In either case, a lower bound on $\bar{g}_{aN}$ arises, for $M_{W_R} \lesssim 20$ or 13 TeV, respectively. The origin of this lower bound is traced to the fact that, in the LRSM with $\mathcal{P}$, for a few TeV M_{W_R} the CPV effects cannot be eliminated by taking $\alpha \to 0$: an exceedingly large contribution to h_ε would remain from the CKM phase in V_R; thus a destructive interference from additional *CP* phases is required [55]. Thus, for instance, a positive detection from ARIADNE below 2×10^{-18} with $m_a \approx 100\ \mu$eV would falsify such a TeV-scale LR DFSZ scenario. Instead, a measurement above 10^{-17} would result in a rejection of the LR DFSZ model or a sharp upper bound on M_{W_R}, at the reach of a future collider.

Given the square root in $(g_{an}\bar{g}_{aN})^{1/2}$, the probed observable depends mildly on the new physics scale. Indeed, the upper boundary of the shaded region decreases as $1/M_{W_R}$, and we find that within the ARIADNE sensitivity the model provides possible signals up to $M_{W_R} \sim 1000$ TeV. Standard flavor observables, decoupling 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 1000 s sensitivity.

To conclude, we provided a complete and consistent calculation of the CPV axion couplings to matter and applied it to the case RH currents, showing 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.

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