

Astrophobic Axions

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We propose a class of axion models with generation-dependent Peccei-Quinn charges for the known fermions that allow one to suppress the axion couplings to nucleons and electrons. Astrophysical limits are thus relaxed, allowing for axion masses up to $O(0.1)$ eV. The axion-photon coupling remains instead sizable, so that next-generation helioscopes will be able to probe this scenario. Astrophobia unavoidably implies flavor-violating axion couplings so that experimental limits on flavor-violating processes can provide complementary probes. The astrophobic axion can be a viable dark matter candidate in the heavy mass window and can also account for anomalous energy loss in stars.

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Introduction.—One of the main mysteries of the standard model (SM) is the absence of CP violation in strong interactions. The most elegant solution is provided by the Peccei-Quinn (PQ) mechanism [1,2], which predicts the axion as a low-energy remnant [3,4]. The axion is required to be extremely light and decoupled, and in a certain mass range it can provide a viable dark matter (DM) candidate. The Kim-Shifman-Vainshtein-Zakharov (KSVZ) [5,6] and Dine-Fischler-Srednicki-Zhitnitsky (DFSZ) [7,8] axion models are frequently used as benchmarks to assess experimental sensitivities and to derive astrophysical bounds. However, constraining axion properties solely on the basis of standard benchmarks can be too restrictive, and exploring alternative models whose properties can sizably deviate from those of the KSVZ and DFSZ models is highly desirable. While it is conceptually easy to build models with suppressed axion-electron couplings g_{ae} [5,6,9] or axion-photon couplings $g_{a\gamma}$ [10–12], it is generally believed that a robust prediction of all axion models is an unsuppressed axion-nucleon coupling g_{aN} . This is particularly important, because g_{aN} is responsible for the often-quoted bound on the axion mass $m_a \lesssim 20$ meV from the neutrino burst duration of SN1987A [13,14]. In this Letter, we argue that a strong suppression of g_{aN} is instead possible in a class of DFSZ-like models with generation-dependent PQ charges. Additional strong bounds

on m_a are obtained if g_{ae} is unsuppressed, since this can affect white-dwarf (WD) cooling rates and red giant (RG) evolution [14]. In our scenario, a suppression of g_{ae} can be also arranged. Thus, nucleophobia allows one to relax the SN bound, and electrophobia allows one to evade the WD and RG constraints, rendering viable $m_a \sim O(0.1)$ eV. We denote such an axion as *astrophobic*, although $g_{a\gamma}$ remains generically sizable and could still affect the evolution of horizontal branch (HB) stars. Astrophobic axions are interesting in many respects: (i) They render viable a parameter space region well beyond the standard DFSZ and KSVZ benchmarks yet still within the reach of the planned International Axion Observatory (IAXO) helioscope [15]. (ii) Nucleophobia necessarily implies flavor-violating (FV) axion couplings to the quarks so that complementary searches can be carried out in flavor experiments. (iii) Astrophobic axions can be nonstandard DM in the heavy mass window [16–18] and (iv) can account for various hints of anomalous star energy losses [19,20].

Axion coupling to nucleons.—Let us first recall why g_{aN} cannot be suppressed in KSVZ and DFSZ models. The relevant terms for this discussion are

$$\mathcal{L}_a \supset \frac{\alpha_s}{8\pi} \frac{a}{f_a} G_{\mu\nu}^a \tilde{G}^{a,\mu\nu} + \frac{\alpha}{8\pi N} \frac{a}{f_a} F_{\mu\nu} \tilde{F}^{\mu\nu} + \frac{\partial_\mu a}{2f_a} \sum_{Q=U,D} \left[\bar{Q}_L \frac{X_{Q_L}}{N} \gamma^\mu Q_L + \bar{Q}_R \frac{X_{Q_R}}{N} \gamma^\mu Q_R \right], \quad (1)$$

where $N(E)$ are the QCD (QED) anomaly coefficients, $f_a = v_a/(2N)$ with $v_a = \sqrt{2}\langle\phi\rangle$ the vacuum expectation value (VEV) of the PQ symmetry-breaking singlet field,

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$\tilde{G}^{a,\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}G_{\rho\sigma}^a$, $\tilde{F}^{\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma}F_{\rho\sigma}$, and $Q_{L,R} = U_{L,R}$, $D_{L,R}$ are vectors containing the left-handed (LH) and right-handed (RH) quarks of the three generations (capital letters denote matrix quantities; q, u, d are used otherwise). In the KSVZ model the PQ charge matrices $X_{Q_{L,R}}$ vanish, while in the DFSZ model they are nonzero but generation blind; hence, the current in Eq. (1) is not dependent on the quark basis. The axion-gluon term can be removed via a field-dependent chiral rotation of the first-generation quarks $q = u, d$: $q_{L,R} \rightarrow e^{\mp i(a/2f_a)f_q}q_{L,R}$ with $f_u + f_d = 1$. Defining $z = m_u/m_d$ and choosing $f_u = 1/(1+z) \simeq 2/3$ avoids axion-pion mixing. As a result of this rotation, the coefficient of the QED term gets shifted as $E/N \rightarrow E/N - f_\gamma(z)$ with $f_\gamma \simeq 1.92$, while the axion coupling to the first-generation quarks becomes

$$\mathcal{L}_{aq} = \frac{\partial_\mu a}{2f_a} \sum_{q=u,d} \left[\bar{q}\gamma^\mu\gamma_5 \left(\frac{X_{q_R} - X_{q_L}}{2N} - f_q \right) q \right]. \quad (2)$$

The vector couplings vanish because of the equation of motion. For the axial-vector couplings, it is common to denote $C_q^0 = (X_{q_R} - X_{q_L})/(2N)$. Matching Eq. (2) with the nonrelativistic axion-nucleon Lagrangian allows one to extract the axion-nucleon couplings [21], which are defined via $\partial_\mu a/(2f_a)C_N \bar{N}\gamma^\mu\gamma_5 N$ with $N = p, n$. We recast the results in terms of the two linear combinations

$$C_p + C_n = 0.50(5)(C_u^0 + C_d^0 - 1) - 2\delta_s, \quad (3)$$

$$C_p - C_n = 1.273(2) \left(C_u^0 - C_d^0 - \frac{1}{3} \right), \quad (4)$$

where the two numbers in parentheses correspond to $f_u + f_d = 1$ (exact) and $f_u - f_d \simeq 1/3$ (approximate), while δ_s is a correction appearing in the DFSZ model which is dominated by the s -quark contribution. In the models below, using the results from Ref. [21] and allowing for the largest possible values of $C_{s,c,b,t}^0$, we have $|\delta_s| \lesssim 0.04$. Equation (3) makes clear why it is difficult to implement axion-nucleon decoupling. For the KSVZ model, $C_u^0 = C_d^0 = 0$ and the model-independent contribution survives. For the DFSZ model, $C_u^0 + C_d^0 = N_l/N$, with N_l the contribution to the QCD anomaly of the light first-generation quarks. Hence, for generation-blind charges, $C_u^0 + C_d^0 = 1/3$ is an exact result.

The nucleophobic axion.—We take as the defining condition for the nucleophobic axion the (approximate) vanishing of Eqs. (3) and (4). Remarkably, since the axion-pion coupling is proportional to $C_p - C_n$ [22], nucleophobic axions are also pionphobic. We start by studying Eq. (3). Neglecting δ_s , $C_p + C_n = 0$ implies $C_u^0 + C_d^0 = N_l/N = 1$. This can be realized in two ways: Either (i) the contributions of the two heavier generations cancel each other ($N_2 = -N_3$ and $N_l = N_1$), or (ii) they vanish identically, in which case it

is convenient to assign $N_l = N_3$ and, hoping that no confusion will arise with generation ordering, require for the heavier generations $N_1 = N_2 = 0$. (This second case was also identified in Ref. [23].) Clearly, both cases require generation-dependent PQ charges. A generic matrix of charges for a LH or RH quark q can be written as $X_Q = X_q^0 I + X_q^8 \lambda_8 + X_q^3 \lambda_3$ with $I = \text{diag}(1, 1, 1)$ the identity in generation space, while $\lambda_8 = \text{diag}(1, 1, -2)$ and $\lambda_3 = \text{diag}(1, -1, 0)$ are proportional to the corresponding SU(3) matrices. Since we are mainly interested in a proof of existence for nucleophobic axions, we introduce some simplification: We assume just two Higgs doublets $H_{1,2}$ (with PQ charges $X_{1,2}$ and hypercharge $Y = -1/2$), and we consider only PQ charge assignments that do not forbid any of the SM Yukawa operators. Under these conditions, it can be shown that two generations must have the same charges [24], and we can then drop the SU(2)-breaking λ_3 term. The matrix $X_Q = X_q^0 I + X_q^8 \lambda_8$ then respects a SU(2) symmetry acting on the generation indices $\{1, 2\}$, and we henceforth refer to such a structure as $2 + 1$. To study which Yukawa structures can enforce the condition $N = N_l$, it is then sufficient to consider just one generation in 2 together with the generation in 1 carrying index $\{3\}$ and write

$$\begin{aligned} \bar{q}_2 u_2 H_1, & \quad \bar{q}_3 u_3 H_a, & \quad \bar{q}_2 u_3 H_b, & \quad \bar{q}_3 u_2 H_{1+a-b}, \\ \bar{q}_2 d_2 \tilde{H}_c, & \quad \bar{q}_3 d_3 \tilde{H}_d, & \quad \bar{q}_2 d_3 \tilde{H}_{d+a-b}, & \quad \bar{q}_3 d_2 \tilde{H}_{c-a+b}, \end{aligned} \quad (5)$$

where $\tilde{H} = i\sigma_2 H^*$. Assigning H_1 to the first term is without a loss of generality, while all the other Higgs indices must take values in $\{1, 2\}$. It is easy to verify that in each line the charges of the first three quark bilinears determine the fourth one, e.g., $X(\bar{q}_3 u_2) = X(\bar{q}_2 u_2) + X(\bar{q}_3 u_3) - X(\bar{q}_2 u_3)$, while the third term in the second line is obtained by equating $X_{q_3} - X_{q_2}$ as extracted from the second and third terms of both lines. It is now straightforward to classify all the possibilities that yield $N_l/N = 1$. Denoting the Higgs ordering in the two lines of Eq. (5) with their indices, e.g., $(H_1, H_2, H_1, H_2)_u \sim (1212)_u$, we have, respectively, for $(i_{1,2})N_1 = N_2 = -N_3$ and $(ii_{1,2})N_1 = N_2 = 0$

$$\begin{aligned} (i_1): & (1212)_u(2121)_d, & (i_2): & (1221)_u(2112)_d, \\ (ii_1): & (1111)_u(1221)_d, & (ii_2): & (1221)_u(1111)_d. \end{aligned} \quad (6)$$

It is easy to verify that in $(i_{1,2})$ $2N_l = 2N_2 = X_{u_{2R}} + X_{d_{2R}} - X_{u_{2L}} - X_{d_{2L}} = X_2 - X_1$ with $N_3 = -N_2$, in (ii_1) $2N_l = 2N_3 = X_2 - X_1$, and in (ii_2) $2N_l = 2N_3 = -X_2 + X_1$ with, in both of the last cases, $N_1 = N_2 = 0$. Let us now discuss the second condition $C_p - C_n \approx 0$. We denote by $\tan\beta = v_2/v_1$ the ratio of the $H_{1,2}$ VEVs, and we introduce the shorthand notation $s_\beta = \sin\beta$, $c_\beta = \cos\beta$. The ratio $X_1/X_2 = -\tan^2\beta$ is fixed by the requirement that the PQ Goldstone boson is orthogonal to the Goldstone boson eaten

up by the Z boson [8], and the charge normalization is given in terms of the light quark anomaly as $X_2 - X_1 = \pm 2N_l$. Here and below, the upper sign holds for (i_{1,2}) and (ii₁) and the lower sign for (ii₂). From Eq. (6), it follows that in all cases $C_u^0 - C_d^0 = -(1/2N)(X_1 + X_2) = \pm(s_\beta^2 - c_\beta^2)$. The second condition for nucleophobia $C_u^0 - C_d^0 = 1/3$ is then realized for $s_\beta^2 = 2/3$ in (i_{1,2}) and (ii₁) and for $s_\beta^2 = 1/3$ in (ii₂). We learn that, even under some restrictive assumptions, there are four different ways to enforce nucleophobia. More possibilities would become viable by allowing for PQ charges that forbid some Yukawa operator [24]. Note that $C_p + C_n \approx 0$ is enforced just by charge assignments, while $C_p - C_n \approx 0$ requires a specific choice $\tan\beta \approx 2^{\pm 1/2}$. For both these values, the top Yukawa coupling remains perturbative up to the Planck scale; however, we stress that they should be understood as relative to the physical VEVs rather than resulting from a tree-level scalar potential, since the large v_a would destabilize any lowest-order result for $v_{1,2}$. This is, of course, a naturalness issue common to all invisible axion models. Finally, to render the axion invisible, $H_{1,2}$ need to be coupled via a non-Hermitian operator to the scalar singlet ϕ with PQ charge X_ϕ . This ensures that the PQ symmetry gets spontaneously broken at the scale $v_a \gg v_{1,2}$, suppressing all axion couplings. There are two possibilities: $H_2^\dagger H_1 \phi$, in which case $|X_\phi| = 2N_l = 2N$, the axion field has the same periodicity as the θ term, and the number of domain walls (DWs) is $N_{\text{DW}} = 1$, or $H_2^\dagger H_1 \phi^2$, in which case $|X_\phi| = N_l = N$ and $N_{\text{DW}} = 2$. In contrast, in DFSZ models $|X_\phi| = 2N/3$, $(2N/6)$ yield $N_{\text{DW}} = 3$, (6) and a DW problem is always present.

Flavor-changing axion couplings.—Generation-dependent PQ charges imply FV axion couplings. Plugging $X_Q = X_{q_1}^0 I + X_{q_3}^8 \lambda_8$ into Eq. (1), it is readily seen that a misalignment between the Yukawa and the PQ charge matrix becomes physical. Since we are mostly interested in the light quark couplings, we single out X_{q_1} for case (i) and X_{q_3} for (ii):

$$X_Q = X_{q_1} I - 3X_{q_3}^8 \Lambda = X_{q_3} I + 3X_{q_1}^8 \Lambda', \quad (7)$$

with $3X_{q_1}^8 = X_{q_1} - X_{q_3}$, $\Lambda = \frac{1}{3}(I - \lambda_8) = \text{diag}(0, 0, 1)$, and $\Lambda' = \frac{1}{3}(2I + \lambda_8) = \text{diag}(1, 1, 0)$. In case (i), the matrices of couplings in the Yukawa basis read

$$C_Q^{0V} = -\frac{3}{2N} [X_{q_R}^8 W_{Q_R} + X_{q_L}^8 W_{Q_L}], \quad (8)$$

$$C_Q^0 + \Delta C_Q^0 = C_{q_1}^0 I - \frac{3}{2N} [X_{q_R}^8 W_{Q_R} - X_{q_L}^8 W_{Q_L}], \quad (9)$$

where for C_Q^{0V} the equations of motion imply vanishing diagonal entries but do not imply vanishing off-diagonal ones, $C_Q^0 = C_{q_1}^0 I$ with $C_{q_1}^0$ defined below Eq. (2), and,

denoting by V_Q the unitary rotations to the diagonal Yukawa basis, $W_Q = V_Q^\dagger \Lambda V_Q$. While in the models discussed here W_{Q_R} and W_{Q_L} are never simultaneously present, this is possible in more general cases [24]. It is now convenient to single out the diagonal (denoted by δ) and off-diagonal (denoted by ω) entries in $W_Q = \delta_Q + \omega_Q$:

$$\begin{aligned} (\delta_Q)_{ij} &= \delta_{q_i} \delta_{ij}, & \sum_i \delta_{q_i} &= 1, \\ (\omega_Q)_{ii} &= 0, & |(\omega_Q)_{ij}|^2 &= \delta_{q_i} \delta_{q_j}, \end{aligned} \quad (10)$$

where the condition on δ_q follows from $\text{Tr}(W_Q) = 1$, the one on ω_Q from the vanishing of the principal minors for the rank one matrix W_Q , and δ_{ij} in the first relation is the usual Kronecker symbol. In (ii), the couplings are given by Eqs. (8) and (9) by replacing $C_{q_1}^0 \rightarrow C_{q_3}^0$, $(-3) \rightarrow (+3)$, and $W_Q \rightarrow W'_Q = V_Q^\dagger \Lambda' V_Q$, while the two conditions read $\sum_i \delta'_{q_i} = 2$ and $|(\omega'_Q)_{ij}|^2 = (1 - \delta'_{q_i})(1 - \delta'_{q_j})$. Information on the LH matrices can be obtained from the Cabibbo-Kobayashi-Maskawa (CKM) matrix: $V_{U_L}^\dagger V_{D_L} = V_{\text{CKM}} \approx I$ implies $V_{U_L} \approx V_{D_L}$ and hence $W_{U_L} \approx W_{D_L}$. Therefore, to a good approximation we can define a single set of LH parameters $\delta_L = \delta_{u_L} \approx \delta_{d_L}$. In contrast, we have no information about the RH matrices. In general, $W_{U_R} \neq W_{D_R}$ so that δ_{u_R} , δ_{d_R} are two independent sets. Corrections to the diagonal axial couplings due to quark mixing are listed in Table I. Corrections to the second condition for nucleophobia can be always compensated by changing appropriately the value of $\tan\beta$ to maintain $C_p - C_n \approx 0$. However, this is not so for the first condition, for which large corrections would spoil $C_p + C_n \approx 0$. Actually, only for (ii₁) can a relatively small correction improve nucleophobia, and this is because in this case C_s^0 , which determines the sign of δ_s in Eq. (3), is negative ($C_s^0 = -s_\beta^2$), rendering possible a tuned cancellation $-0.50\delta'_{d_{3R}} + 2|\delta_s| \approx 0$. Thus, nucleophobia generically requires quark Yukawa and PQ charge matrices approximately aligned (for recent attempts to connect axion physics to flavor dynamics, see [25–27]).

Electrophobia.—Electrophobia can be implemented exactly (at the lowest loop order) or approximately (modulo

TABLE I. Contributions from the quarks to E/N and corrections to the nucleophobic axion couplings due to quark mixings. The (off-diagonal) vector couplings C_q^{0V} are equal in modulus to the axial-vector ones.

	E_Q/N	ΔC_u^0	ΔC_d^0	$ C_u^0 _{i \neq j}$	$ C_d^0 _{i \neq j}$
(i ₁)	$-4/3 + 6s_\beta^2$	$-\delta_{1L}$	$-\delta_{1L}$	ω_L	ω_L
(i ₂)	$-4/3 + 6s_\beta^2$	$-\delta_{u_{1R}}$	$-\delta_{d_{1R}}$	ω_{u_R}	ω_{d_R}
(ii ₁)	$2/3 + 6s_\beta^2$	0	$-\delta'_{d_{3R}}$	0	ω'_{d_R}
(ii ₂)	$8/3 - 6s_\beta^2$	$-\delta'_{u_{3R}}$	0	ω'_{u_R}	0

TABLE II. Contributions of the leptons and total values of E/N in four representative models, selected by the (arbitrary) choice that the electron couples to \tilde{H}_1 . The numerical values of C_e^0 are given in parentheses, and the corrections ΔC_e^0 can come from RH or LH mixings.

	E_L/N	E/N	C_e^0	ΔC_e^0
M1: (i) + (i _l)	$2 - 6s_\beta^2$	2/3	$-s_\beta^2(-2/3)$	$+\delta_{e_1}$
M2: (ii ₁) + (i _l)	$2 - 6s_\beta^2$	8/3	$-s_\beta^2(-2/3)$	$+\delta_{e_1}$
M3: (ii ₂) + (ii _l)	$-4 + 6s_\beta^2$	-4/3	$s_\beta^2(1/3)$	$-\delta'_{e_3}$
M4: (ii ₁) + (ii _l)	$4 - 6s_\beta^2$	14/3	$-s_\beta^2(-2/3)$	$+\delta'_{e_3}$

lepton mixing corrections) by introducing an additional Higgs doublet uncharged under the PQ symmetry and by coupling it, respectively, to all the leptons or just to the electron. However, electrophobia can also be implemented without enlarging the Higgs sector at the cost of tuning a cancellation between C_e^0 and a mixing correction. This requires large lepton mixings and one fine-tuning. Given that large mixings do characterize the lepton sector, at least the first requirement is not unnatural. It is easy to verify that in all the following cases a cancellation is possible: We can assign the electron (i_l) to the doublet in $2 + 1$ or (ii_l) to the singlet, and in both cases we can consider $(12\dots)_l$ or $(21\dots)_l$ structures and combine these possibilities with the four quark cases. Moreover, for $(abab)_l$ type of structures electrophobia is enforced by a cancellation from LH mixing, while for $(abba)_l$ from RH mixing. All in all, there are $2 \times 2 \times 4 \times 2 = 32$ physically different astrophobic models. However, as regards the axion-photon coupling, there are only four different values of E/N . They are listed in Table II for four representative models.

Phenomenology of the heavy axion window.—We denote as $g_{af} = C_f m_f / f_a$ the axion coupling to $f = p, n, e$, including corrections from mixing effects, and by $g_{a\gamma} = \alpha / (8\pi f_a) (E/N - f_\gamma)$ the coupling to photons. The most relevant astrophysical bounds are [13,14] as follows: (a) $|g_{a\gamma}| < 6.6 \times 10^{-11} \text{ GeV}^{-1}$ (95% C.L.) from the evolution of HB stars [28]; (b) $|g_{ae}| < 2.7 \times 10^{-13}$ ($< 4.3 \times 10^{-13}$) (95% C.L.) from the shape of the WD luminosity function [29] (from RG evolution [30]); (c) $g_{ap}^2 + g_{an}^2 < 3.6 \times 10^{-19}$ from the SN1987A neutrino burst duration [20]—large uncertainties in estimating SN axion emissivity [31,32] prevent assigning a reliable statistical significance to this limit; (d) structure-formation arguments also provide hot DM (HDM) limits on the axion mass: In benchmark models, $m_a \lesssim 0.8 \text{ eV}$ [13,33,34]. However, nucleophobic axions are also pionophobic, and the main thermalization process $\pi\pi \rightarrow \pi a$ is then suppressed, relaxing the HDM bound. This implies that large-volume surveys like EUCLID [35] cannot probe astrophobic axions.

The main results for astrophobic axions are summarized in Fig. 1 and compared to the KSVZ and DFSZ benchmarks.

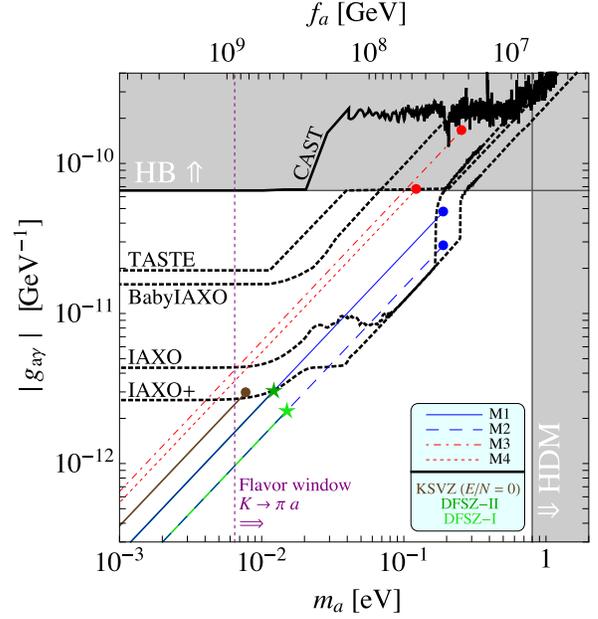


FIG. 1. Axion-photon coupling $|g_{a\gamma}|$ for the astrophobic models in Table II as a function of m_a . The DFSZ-I,II (respectively, with $E/N = 8/3, 2/3$) and KSVZ benchmarks are also shown for comparison.

The lines are broken at • marks, which indicate the upper bounds on m_a from SN1987A, and ★ marks, corresponding to the combined SN and WD constraints for DFSZ models. As anticipated, for the KSVZ and DFSZ models, axion masses above $m_a \sim 10^{-2} \text{ eV}$ are precluded by the SN and WD limits (dark brown bullet for the KSVZ model and green stars for the DFSZ model). For astrophobic axions, the SN and WD bounds get significantly relaxed [they cannot evaporate completely because of the contribution δ_s in Eq. (3) to g_{aN}]. We obtain $m_a < 0.20 \text{ eV}$ for M1 and M2 (blue bullets), $m_a < 0.25 \text{ eV}$ for M3, and $m_a < 0.12 \text{ eV}$ for M4 (red bullets).

Searches with helioscopes.—Helioscopes are sensitive to $g_{a\gamma}$, which is not particularly suppressed in astrophobic models. The solid black line in Fig. 1 shows the present limits from the CERN Axion Solar Telescope (CAST) [36], while the dotted black lines show the projected sensitivities of next-generation helioscopes. While the improvement in mass reach will be limited for the Troitsk Axion Solar Telescope Experiment (TASTE) [37] and BabyIAXO [38], we see that IAXO [15,39] and its upgrade IAXO+ [20] will be able to cover the whole interesting region up to $m_a \sim 0.2 \text{ eV}$.

Flavor violation.—The strongest limits on FV axion couplings come from $K^+ \rightarrow \pi^+ a$ [40]. Comparing the model prediction with the current limit [41] gives

$$\mathcal{B}_{K^+ \rightarrow \pi^+ a} \simeq 10^{-2} \left(\frac{m_a}{0.2 \text{ eV}} \right)^2 \omega_{12}^2 \lesssim 7.3 \times 10^{-11}, \quad (11)$$

where $\omega_{12}^2 = |\omega_{12}|^2 = \delta_{1L}\delta_{2L}, \delta_{d_{1R}}\delta_{d_{2R}}$ for $(i_{1,2})$ and $\omega_{12}^2 = |\omega'_{32}|^2$ for (ii_1) , while in (ii_2) the branching ratio vanishes

(see Table I). For models (i_{1,2}), CKM-like entries in $V_{d_{L,R}}$ would yield $\omega_{12}^2 \gtrsim 10^{-8}$ saturating the limit. This implies that NA62, which is expected to improve the limit by a factor of ~ 70 [42,43], can probe these models (the exploratory “flavor window” corresponds to the vertical magenta line in Fig. 1). For (ii₁), assuming CKM-like mixings the result would exceed the limit Eq. (11) by 6 orders of magnitude, so that this case is rather unrealistic. If we allow for only two Higgs doublets, electrophobic models necessarily have FV axion couplings to leptons. The strongest limits come from searches for $\mu \rightarrow e\gamma$ [44,45]. They yield $f_a/\omega_{e\mu} \gtrsim 2 \times 10^9$ GeV, which implies $m_a \lesssim 2.7 \times 10^{-3}/\omega_{e\mu}$ eV. Recalling that $\omega_{e\mu} = \sqrt{\delta_e \delta_\mu}$ and that $\delta_e \sim O(1)$ is needed to cancel C_e^0 , we need to impose $\delta_\mu \lesssim 10^{-4}$ to avoid overconstraining the large mass window. This implies $\delta_\tau \sim O(1)$, from which we can predict $\mathcal{B}_{\tau \rightarrow ea} \simeq 7 \times 10^{-6} (m_a/0.2 \text{ eV})^2$, about 3 orders of magnitude below the present bound [46].

Axion DM in the heavy mass window.—For $m_a \sim 0.2$ eV, the misalignment mechanism cannot fulfill $\Omega_a \simeq \Omega_{\text{DM}}$. In postinflationary scenarios, if $N_{\text{DW}} > 1$ [16,17], an additional contribution from topological defects can concur to saturate Ω_{DM} . This requires an explicit PQ breaking to trigger DW decays and a fine-tuning not to spoil the solution to the strong CP problem. For $N_{\text{DW}} = 1$, a contribution to Ω_a can come from axion production via parametric resonance in the oscillations of the axion field radial mode [18], in which case $f_a \lesssim 10^{18}$ GeV is needed. In both cases, the lower values of f_a allowed by the astrophobic models can help to match the required conditions.

Stellar cooling anomalies.—Estimates for anomalous star energy losses [19,20] can be more easily accommodated in astrophobic axion models. Reference [20] finds the best-fit point for extra axion cooling $g_{a\gamma} \sim 0.14 \times 10^{-10}$ GeV⁻¹ and $g_{ae} \sim 1.5 \times 10^{-13}$, which is disfavored in the DFSZ model but is comfortably within the allowed parameter space of the astrophobic axion.

Conclusions.—We have discussed a class of DFSZ-like axion models with generation-dependent PQ charges that allows relaxation of the SN1987A bound on g_{aN} and the WD and RG limit on g_{ae} and to extend the viable axion mass window up to $m_a \sim 0.2$ eV. This scenario is characterized by compelling connections with flavor physics. Complementary information for direct axion searches can be provided by experimental searches for FV meson and lepton decays, and, conversely, the discovery of this type of astrophobic axion would provide evidence that the quark Yukawa matrices are approximately diagonal in the interaction basis, conveying valuable information on the SM flavor structure. While we have restricted our analysis to PQ charge assignments which do not forbid any of the SM Yukawa operators, it would be interesting to relax this condition and explore the extent to which the PQ symmetry

could play a role as a flavor symmetry in determining specific textures for the SM Yukawa matrices.

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