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 $g_{a\gamma}$ is unsuppressed, since this can affect white-dwarf (WD) cooling rates and red giant (RG) evolution [14]. In our scenario, a suppression of $g_{a\gamma}$ 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

$$\mathcal{L}_a \supset \frac{\alpha_e}{8\pi} a F_{\mu\nu} \tilde{G}^\mu_{\nu a} + \frac{\alpha}{8\pi} \frac{E}{N} F_{\mu\nu} \tilde{F}^\mu_{\nu a} + \frac{\beta}{2f_a} \sum_{Q=U,D} \left[ \tilde{Q}_L X_{Qa} \gamma^\mu Q_L + \tilde{Q}_R X_{Qa} \gamma^\mu Q_R \right],$$

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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\[ \tilde{C}^{a,\mu \nu} = \frac{1}{2} \epsilon^{\mu \nu \rho \sigma} G^{a}_{\rho \sigma}, \quad \tilde{F}^{\mu \nu} = \frac{1}{2} \epsilon^{\mu \nu \rho \sigma} F_{\rho \sigma}, \quad Q^{a}_{L,R} = U^{a}_{L,R}. \]

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-glue term can be removed via a field-dependent chiral rotation of the first-generation quarks \( q = u, \ d: \ q_{L,R} \to e^{\mp i(q/2a)} + q_{L,R} \) with \( f_u + f_d = 1 \). Defining \( z = m_u/m_d \) and choosing \( f_u = 1/(1 + z) \approx 2/3 \) avoids axion-pion mixing. As a result of this rotation, the coefficient of the QED term gets shifted as \( E/N \to E/N - f_{\gamma}(z) \) with \( f_{\gamma} \approx 1.92 \), while the axion coupling to the first-generation quarks becomes

\[ L_{\text{aq}} = \frac{\partial_{\mu} a_{\mu}}{2 f_{\gamma} a} \sum_{q \neq u,d} \left[ \frac{\bar{q}^\nu \gamma_5 (X_{q_\mu} - X_{q_L} - f_{\gamma} q)}{2N} \right]. \]

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

\[ C_p + C_n = 0.50(5) (C^0_q + C^0_d - 1) - 2 \delta_q, \]

\[ C_p - C_n = 1.273(2) \left( C^0_q - C^0_d - \frac{1}{3} \right), \]

where the two numbers in parentheses correspond to \( f_u + f_d = 1 \) (exact) and \( f_u - f_d \approx 1/3 \) (approximate), while \( \delta_q \) 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 \( C^0_{s,c,b,t} \), we have \( |\delta_q| \lesssim 0.04 \). Equation (3) makes clear why it is difficult to implement axion-nucleon decoupling. For the KSVZ model, \( C^0_q = 0 = 0 \) and the model-independent contribution survives. For the DFSZ model, \( C^0_q + C^0_d = 1/3 \) is an exact result.

The nucleophobic axion.—We take as the defining condition for the nucleophobic axion the (approximately) vanishing of Eqs. (3) and (4). Remarkably, since the axion-pion coupling is proportional to \( C_p - C_n \) [22], nucleophobic axions are also pionophobic. We start by studying Eq. (3). Neglecting \( \delta_q \), \( C_p + C_n = 0 \) implies \( C^0_q + C^0_d = N_1/N = 1 \). This can be realized in two ways: Either (i) the contributions of the two heavier generations cancel each other (\( N_2 = -N_3 \) and \( N_1 = N_3 \)), or (ii) they vanish identically, in which case it is convenient to assign \( N_1 = N_3 \) 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 \( X_Q = X^0_Q I + X^8_Q \delta_8 + X^3_Q \lambda_3 \) \( (I = \text{diag}(1,1,1)) \) the identity in generation space, while \( \lambda_8 = \text{diag}(1,1,-2) \) and \( \lambda_3 = \text{diag}(-1,0) \) 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 \( \lambda_3 \) term. The matrix \( X_Q = X^0_Q I + X^8_Q \delta_8 \) then respects a SU(2) symmetry acting on the generation indices \( \{1,2\} \), and we henceforth refer to such a structure as \( 2 + I \). To study which Yukawa structures can enforce the condition \( N = N_1 \), it is then sufficient to consider just one generation in 2 together with the generation in \( I \) carrying index \( \{1\} \) and write

\[ \begin{align*}
q_1^u H_1, & \quad q_3^u H_a, & \quad q_3^d H_b, & \quad q_3^d H_{1+a-b}, \\
q_2^d H_c, & \quad q_3^d H_d, & \quad q_2^d H_{d+a-b}, & \quad q_3^d H_{c+a-b},
\end{align*} \]

where \( \tilde{H} = i e^2 H^* \). 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., \( X(q_3 u_2) = X(q_2 u_2) + X(q_3 u_3) - X(q_2 u_3) \), 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 \( N_1/N = 1 \). Denoting the Higgs ordering in the two lines of Eq. (5) with their indices, e.g., \((H_1, H_2, H_{1,2}, H_{a,b,c}) \) \( \sim (1212)_u \), we have, respectively, for \((i_{1,2})N_1 = N_2 = -N_3 \) and \((i_{1,2})N_1 = N_2 = 0 \)

\[ \begin{align*}
(i_{1,2})_1 & \sim (1212)_u (2112)_d, \\
(i_{1,2})_2 & \sim (2212)_u (2112)_d.
\end{align*} \]

It is easy to verify that in \((i_{1,2})_2 2N_1 = 2N_2 = X_{u_2} + X_{d_2} - X_{u_2} - X_{d_2} = X_2 - X_1 \) with \( N_1 = -N_2 \), in \((i_{1,2})_2 2N_1 = 2N_2 = X_2 - X_1 \), and in \((i_{1,2})_2 2N_1 = 2N_2 = X_2 - X_1 \), with 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 Higgs 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.

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up by the Z boson [8], and the charge normalization is given in terms of the light quark anomaly as \( X_2 - X_1 = \pm 2N_f \). Here and below, the upper sign holds for (i) and (ii) and the lower sign for (ii). From Eq. (6), it follows that in all cases \( C^0_u - C^0_d = -(1/2N)(X_1 + X_2) = \pm (s^2_\beta - c^2_\beta) \). The second condition for nucleophobia \( C^0_u - C^0_d = 1/3 \) is then realized for \( s^2_\beta = 2/3 \) in (i) and (ii) and for \( s^2_\beta = 1/3 \) in (ii). 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 \( v_d \) 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, \( \omega_q \) should be understood as relative to the physical VEVs rather than resulting from a tree-level scalar potential, since the \( \delta \) would destabilize any lowest-order result for \( \omega_q \), the axion field has big Eq. (8) and (9) by replacing \( C^i_q \to C^i_q \), \( (3) \to (4,3) \), and \( W_Q \to W_Q' = V_Q^L \Lambda Q = V_Q^L \Lambda Q = V_Q^L \Lambda Q = V_Q^L \Lambda Q \), where the two conditions read \( \sum_i \delta_{ij} = 2 \) and \( |(\omega_q^i)^2| = (1 - \delta_{ij})(1 - \delta_{ij}) \). Information on the LH matrices can be obtained from the Cabibbo-Kobayashi-Maskawa (CKM) matrix: \( V_{U_L}^d = V_{CKM} \approx I \) implies \( V_{U_L} \approx V_{L} \) and hence \( W_{U_L} \approx W_{D_L} \). Therefore, to a good approximation we can define a single set of LH parameters \( \delta_L = \delta_{L_L} = \delta_{L_R} \). In contrast, we have no information about the RH matrices. In general, \( W_{U_R} \neq W_{D_R} \), so that \( \delta_{R_L} \neq \delta_{R_R} \) 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 (ii) (ii) can a relatively small correction improve nucleophobia, and this is because in this case \( C^i_q \), which determines the sign of \( \delta_q \) in Eq. (3), is negative \( C^i_q = -s^2_\beta \), rendering possible a tuned cancellation \(-0.50 \delta_{ij}^2 + 2|\delta_q| \approx 0 \). 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]).

**Flavor-changing axion couplings.**—Generation-dependent PQ charges imply FV axion couplings. Plugging \( X_Q = X_Q^0 + X_Q^\lambda \delta_q \) 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 \( X_{q_i} \) for (i) and \( X_{q_i} \) for (ii):

\[
X_Q = X_{q_i} I - 3X^8_q \Lambda = X_{q_i} I + 3X^8_q \Lambda',
\]

(7)

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

\[
C^W_Q = -\frac{3}{2N}[X^8_q W_{q_u} + X^8_q W_{q_L}],
\]

(8)

\[
C^W_Q + \Delta C^W_Q = C^W_Q \pm \frac{3}{2N}[X^8_q W_{q_u} - X^8_q W_{q_u}],
\]

(9)

where for \( C^W_Q \) the equations of motion imply vanishing diagonal entries but do not imply vanishing off-diagonal ones, \( C^W_Q \equiv C^W_{q_1} \) with \( C^W_{q_1} \) defined below Eq. (2), and, denoting by \( V_Q \) the unitary rotations to the diagonal Yukawa basis, \( W_Q = V_Q^L V_Q^R \). While in the models discussed here \( W_{q_u} \) and \( W_{q_d} \) are never simultaneously present, this is possible in more general cases [24]. It is now convenient to single out the diagonal (denoted by \( \delta \)) and off-diagonal (denoted by \( \omega \)) entries in \( W_Q = \delta_Q + \omega_Q \):

\[
(\delta_Q)_{ij} = \delta_{ij}, \quad \sum_i \delta_{ij} = 1,
\]

\[
(\omega_Q)_{ij} = 0, \quad |(\omega_Q)_{ij}|^2 = \delta_{ij},
\]

(10)

where the condition on \( \delta_q \) follows from \( \text{Tr}(W_Q) = 1 \), the one on \( \omega_q \) from the vanishing of the principal minors for the rank one matrix \( W_Q \), and \( \delta_{ij} \) in the first relation is the usual Kronecker symbol. In (ii), the couplings are given by Eqs. (8) and (9) by replacing \( C^i_q \to C^i_q \), \( (3) \to (4,3) \), and \( W_Q \to W_Q' = V_Q^L \Lambda Q \), while the two conditions read \( \sum_i \delta_{ij} = 2 \) and \( |(\omega_q^i)^2| = (1 - \delta_{ij})(1 - \delta_{ij}). \) Information on the LH matrices can be obtained from the Cabibbo-Kobayashi-Maskawa (CKM) matrix: \( V_{U_L}^d = V_{CKM} \approx I \) implies \( V_{U_L} \approx V_{L} \) and hence \( W_{U_L} \approx W_{D_L} \).

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

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**Table I.** Contributions from the quarks to \( E/N \) and corrections to the nucleophobichaxion couplings due to quark mixings. The (off-diagonal) vector couplings \( C^0 \) are equal in modulus to the axial-vector ones.

| \((i_i)_l\) | \(+4/3 + 6x^2_q\) | \(-\delta_{ij}L\) | \(-\delta_{ij}L\) | \(\omega_{ij}L\) | \(\omega_{ij}L\) |
|-----------|-----------------|-----------------|-----------------|-----------------|-----------------|
| 0         | \(-2/3 + 6x^2_q\) | \(-\delta_{ij}L\) | \(-\delta_{ij}L\) | \(\omega_{ij}L\) | \(\omega_{ij}L\) |
| 0         | \(-2/3 + 6x^2_q\) | \(-\delta_{ij}L\) | \(-\delta_{ij}L\) | \(\omega_{ij}L\) | \(\omega_{ij}L\) |
| 0         | \(-2/3 + 6x^2_q\) | \(-\delta_{ij}L\) | \(-\delta_{ij}L\) | \(\omega_{ij}L\) | \(\omega_{ij}L\) |
| 0         | \(-2/3 + 6x^2_q\) | \(-\delta_{ij}L\) | \(-\delta_{ij}L\) | \(\omega_{ij}L\) | \(\omega_{ij}L\) |
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 \( C^0_e \) 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 (ii) to the doublet in 2 + 1 or (iii) to the singlet, and in both cases we can consider (12...) or (21...) structures and combine these possibilities with the four quark cases. Moreover, for (abab), type of structures electrophobia is enforced by a cancellation from LH mixing, while for (abba), from RH mixing. All in all, there are 2 x 2 x 4 x 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 \( g_{\alpha\gamma} = C f_\alpha f_\gamma \) the axion to photon coupling to \( \alpha = p, n, e \), including corrections from mixing effects, and by \( g_{\alpha\gamma} = \alpha/(8\pi f_\alpha) E/N - f_\gamma \) the coupling to photons. The most relevant astrophysical bounds are [13,14] as follows: (a) \( |g_{\alpha\gamma}| < 6.6 \times 10^{-11} \text{ GeV}^{-1} \) (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) \( g_{\alpha\gamma}^2 + g_{\alpha\gamma}^2 < 3.6 \times 10^{-10} \) 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, \( m_\alpha < 0.8 \text{ eV} \) [13,33,34]. However, nucleophobic axions are also pionphobic, and the main thermalization process \( \pi\pi \rightarrow \pi\pi \) 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.

**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 \( H_1 \). The numerical values of \( C^0_i \) are given in parentheses, and the corrections \( \Delta C^0_i \) can come from RH or LH mixings.

| \( E_L/N \) | \( E/N \) | \( C^0_i \) | \( \Delta C^0_i \) |
|---|---|---|---|
| M1: (i) + (i) | 2 - 6\( \delta^2 \) | 2/3 | - \( \delta^2 (2/3) \) + \( \delta_1 \) |
| M2: (ii) + (i) | 2 - 6\( \delta^2 \) | 8/3 | - \( \delta^2 (2/3) \) + \( \delta_2 \) |
| M3: (ii) + (ii) | -4 + 6\( \delta^2 \) | -4/3 | \( \delta^2 (1/3) \) - \( \delta_3 \) |
| M4: (ii) + (i) | 4 - 6\( \delta^2 \) | 14/3 | - \( \delta^2 (2/3) \) + \( \delta_4 \) |

The lines are broken at • marks, which indicate the upper bounds on \( m_\alpha \) 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 \( m_\alpha \sim 10^{-2} \text{ eV} \) 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 \( \delta_i > 0 \) in Eq. (3) to \( g_{\alpha\gamma} \)]. We obtain \( m_\alpha < 0.20 \text{ eV} \) for M1 and M2 (blue bullets), \( m_\alpha < 0.25 \text{ eV} \) for M3, and \( m_\alpha < 0.12 \text{ eV} \) for M4 (red bullets).

**Searches with helioscopes.**—Helioscopes are sensitive to \( g_{\alpha\gamma} \), 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+ [40] will be able to cover the whole interesting region up to \( m_\alpha \sim 0.2 \text{ eV} \).

**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

\[
B_{K^+ \rightarrow \pi^+ a} \approx 10^{-2} \left( \frac{m_\alpha}{0.2 \text{ eV}} \right)^2 \omega_{12}^2 \lesssim 7.3 \times 10^{-11},
\]

where \( \omega_{12}^2 = |\omega_{12}|^2 = \delta_{1L} \delta_{2L} \), \( \delta_{1h} \delta_{2h} \) for (i1,2) and \( \omega_{12}^2 = |\omega_{12}|^2 \) for (ii1), while in (ii2) the branching ratio vanishes.
(see Table 1). For models (i$_{1,2}$), CKM-like entries in $V_{d_{l,R}}$ would yield $\omega_{1,2} > 10^{-8}$ saturating the limit. This implies that NA62, which is expected to improve the limit by a factor of $\sim 70$ [42,43], can probe these models (the explorable “flavor window” corresponds to the vertical magenta line in Fig. 1). For (ii), 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 $\mu \to e\nu\nu$ [44,45]. They yield $f_a/\omega_{\mu\nu} \approx 2 \times 10^9$ GeV, which implies $m_a \lesssim 2.7 \times 10^{-5}/\omega_{\mu\nu} \text{eV}$. Recalling that $\omega_{\mu\nu} = \sqrt{\delta_\mu \delta_\mu}$, and that $\delta_\mu \sim O(1)$ is needed to cancel $C_\nu$, we need to impose $\delta_\mu \lesssim 10^{-4}$ to avoid overconstraining the large mass window. This implies $\delta_\mu \sim O(1)$, from which we can predict $B_{\mu \to e\nu} \approx 7 \times 10^{-6}(m_a/0.2 \text{ eV})^2$, 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 $\Omega_a \approx \Omega_{\text{DM}}$. In postinflationary scenarios, if $N_{\text{DW}} > 1$ [16,17], an additional contribution from topological defects can concur to saturate $\Omega_{\text{DM}}$. 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 $\Omega_a$ 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 $g_{a\nu} \sim 1.5 \times 10^{-13}$, 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 could play a role as a flavor symmetry in determining specific textures for the SM Yukawa matrices.

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