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Axion mass prediction from minimal grand unification

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We propose a minimal realization of the Peccei Quinn mechanism in a realistic SU(5) model, where the axion mass is directly connected to the grand-unification scale. By taking into account constraints from proton decay, collider searches and gauge coupling unification, we predict the axion mass: $m_a \in [4.8, 6.6]$ neV. The upper bound can be relaxed up to $m_a < 330$ neV, at the cost of tuning the flavour structure of the proton decay operators. The predicted mass window will be complementarily probed by the axion dark matter experiments ABRACADABRA and CASPER-Electric, which could provide an indirect evidence for the scale of grand unification before the observation of proton decay.

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Introduction. It is a widespread belief that the standard model (SM) of particle physics should break down at some intermediate energy between the electroweak and the Planck scale. The quantum numbers of the SM fermions, together with the apparent convergence of the SM gauge couplings at high energies, hint to a unified gauge dynamics around 10^{15} GeV. This scale is generically compatible with indirect constraints from the non-observation of proton decay, the smoking-gun signature of Grand Unified Theories (GUTs). The search for proton decay was vigorously pushed in the past decades, and has slowly reached its limits with the Super-Kamiokande (SK) observatory [1]. Planned large-volume facilities, such as Hyper-Kamiokande (HK) [2], will improve the bound on the proton lifetime by one order of magnitude in the next decade. Though fundamentally important, that translates only into a factor of two on the GUT scale.

Another well-motivated framework which points to energies in between the electroweak and the Planck scale is associated with the Peccei-Quinn (PQ) solution of the strong CP problem [3, 4], which predicts the axion as a low-energy remnant [5, 6]. The axion needs to be extremely light and decoupled, and in a certain mass range it is also a viable dark matter (DM) candidate [7–9]. The experimental program for axion searches is rapidly evolving, with many novel detection techniques and new experiments being proposed recently [10]. It is reasonable to expect that a large portion of the parameter space predicted by the QCD axion will be probed in the next decade. From an experimental point of view, however, one of the main bottlenecks of axion DM searches (e.g. those exploiting microwave cavities or nuclear magnetic resonance techniques) is the need to perform a fine scan in the axion mass in order to meet a resonance condition. Since the axion mass is not predicted by the PQ mechanism, any extra theoretical information which could pin-down precisely the axion mass would be extremely helpful for experiments.

Following recent attempts to revive PQ-GUTs in SO(10) [11] (see also [12–18]), in this Letter we revisit the more minimal option of SU(5). The simplest implementation of the axion in non-supersymmetric[1] SU(5) was proposed long ago by Wise, Georgi and Glashow (WGG) [22]. However, similarly to the original SU(5) model of Georgi and Glashow (GG) [23], the WGG model is ruled out in its minimal formulation because of gauge coupling unification and neutrino masses. An elegant and minimal way to fix both these issues in the GG model was put forth some years ago by Bajc and Senjanović [24], which add to the minimal GG field content a single Majorana fermion representation, 24_F, transforming in the adjoint of SU(5). The extra degrees of freedom have the right quantum numbers to generate neutrino masses via a hybrid Type-I+III seesaw mechanism and ensure a proper unification pattern. In particular, the main observable emerging from detailed renormalization group analyses of

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[1] The reader might wonder why we care for the fine-tuning of $|\theta_{\rm QCD}| \lesssim 10^{-10}$ and not for the electroweak-GUT hierarchy. A possible answer is that the strong CP problem is qualitatively different from the hierarchy problem, and it is conceivable that the solution of the latter does not rely on a stabilizing symmetry (an interesting example is the possibility that a light Higgs might be selected by the cosmological evolution of the universe [19–21]).

the GG+24_F model (see Refs. [24–26]) is a clean correlation between light electroweak triplet states (constrained by the Large Hadron Collider (LHC)) and the unification scale (constrained by SK).

Having in mind the possibility of narrowing the axion mass range within a minimal and realistic extension of the WGG model, we extend the latter with a 24_F in analogy to the GG+24_F case. This is actually welcome also from the point of view of the GG+24_F model, which lacks a DM candidate. Within the WGG model (or any realistic extension of it) the axion mass can be put in one-to-one correspondence with the proton decay rate, regardless of the fine details of gauge coupling unification. This allows us to extract a generic upper bound on the axion mass. Including also the detailed information from gauge coupling unification available in the WGG+24_F model, we are also able to set a lower bound on the axion mass from the non-observation of electroweak-triplet states at LHC, thus predicting the following axion mass window: $m_a \in [4.8, 6.6]$ neV, where the upper bound holds in the absence of tuning of fermionic mixing. Next, we provide the axion coupling to the SM fields and estimate the sensitivity of future axion DM experiments such as ABRACADABRA [27] and CASPEr [28, 29] in the relevant mass window.

The WGG model. Let us recall the main features of the WGG model [22]. While the fermion content is that of the original GG SU(5) [23], namely three copies of $\overline{5}_F$ and 10_F comprising the chiral SM matter fields, the scalar sector is extended to include a *complex* 24_H and *two* fundamentals, 5_H and $5'_H$. The WGG Lagrangian can be written as $\mathcal{L}_{\text{WGG}} = \mathcal{L}_{\text{kin}} + \mathcal{L}_Y - V_H$, where $\mathcal{L}_{\text{kin}}$ encodes the (gauge) kinetic terms, the Yukawa Lagrangian is schematically[2]

$$\mathcal{L}_Y = \overline{5}_F 10_F 5'^*_H + 10_F 10_F 5_H + \text{h.c.}\,, \tag{1}$$

while the scalar potential (which we do not report here entirely) contains two non-trivial invariants which are affected by global re-phasings:

$$V_H \supset 5'^\dagger_H 24^2_H 5_H + 5'^\dagger_H 5_H \,\text{Tr}\,(24^2_H) + \text{h.c.}\,. \tag{2}$$

Note that the structure of the WGG Lagrangian resembles that of the DFSZ model [30, 31]. In fact $\mathcal{L}_{\text{WGG}}$ is invariant under the global $\text{U}(1)_{\text{PQ}}$ transformation: $\overline{5}_F \to e^{-i\alpha/2}\overline{5}_F$, $10_F \to e^{-i\alpha/2}10_F$, $5_H \to e^{i\alpha}5_H$, $5'_H \to e^{-i\alpha}5'_H$ and $24_H \to e^{-i\alpha}24_H$.

We have performed the minimization of the full scalar potential in [22] and computed in turn the particle spectrum. In particular, it can be shown that the vacuum expectation value (VEV) configuration

$$\langle 24_H \rangle = V \tfrac{1}{\sqrt{30}}\,\text{diag}(2,2,2,-3,-3)\,, \tag{3}$$

breaks SU(5)×$\text{U}(1)_{\text{PQ}}$ down to the SM gauge group with a *single* order parameter V.[3] The axion, the (pseudo) Nambu-Goldstone boson of the global $\text{U}(1)_{\text{PQ}}$, is dominantly contained in the phase along the SM singlet direction of 24_H, i.e.

$$24_H \supset \langle 24_H \rangle \tfrac{1}{\sqrt{2}} e^{ia/V}\,. \tag{4}$$

A crucial point of the WGG model is that the mass of the heavy vector leptoquark $V_\mu = (3, 2, -5/6)$ mediating proton decay,

$$m_V = \sqrt{\tfrac{5}{6}}\, g_5 V\,, \tag{5}$$

(where g_5 denotes the SU(5) gauge coupling) is directly connected to the axion decay constant[4]

$$f_a = V/\hat{N}\,, \tag{6}$$

where $\hat{N}$ is the $\text{U}(1)_{\text{PQ}}$-$\text{SU}(3)_C$-$\text{SU}(3)_C$ anomaly coefficient, e.g. $\hat{N} = 6$ in the WGG model.

This implies a generic relation between the axion mass and the proton decay rate. By means of chiral effective field theory techniques, we can recast the master formula for the proton decay mode $p \to \pi^0 e^+$ in SU(5) as [33, 34]:

$$\begin{aligned}\Gamma_{p\to\pi^0 e^+} &= \frac{m_p}{16\pi f_\pi^2} A_L^2 \,|\alpha|^2 \,(1 + D + F)^2 \\ &\quad\times \left(\frac{g_5^2}{2m_V^2}\right)^2 \left[4A_{SL}^2 + A_{SR}^2\right]\,,\end{aligned} \tag{7}$$

where we have set unknown fermion mixing rotations to a unit matrix (see [34] for complete expressions). $A_L = 1.25$ encodes the renormalization from the electroweak scale to the proton mass, $m_p = 938.3$ MeV; $f_\pi = 139$ MeV, $D = 0.81$, $F = 0.44$ and $\alpha = -0.011$ GeV^3 are phenomenological parameters given by the chiral Lagrangian and the lattice. $A_{SL(R)}$ are short-distance renormalization factors from the GUT to the electroweak scale which depend on the intermediate-scale thresholds [35, 36]. Compact expressions for the latter can be found e.g. in Ref. [37]. For instance, running within the SM from 10^{15} GeV to the electroweak scale yields $A_{SL} = 2.4$ and $A_{SR} = 2.2$.

[2] Non-renormalizable operators or extra scalar representations are further required in order to correct the ratio between the masses of down quarks and charged leptons.

[3] The recent work [32], which bears some analogies with our proposal, differs crucially in the fact that the PQ symmetry is broken by an SU(5) singlet and hence the axion mass cannot be predicted.

[4] We neglect corrections depending on weak-scale VEVs. For a pedagogical introduction and practical recipes on how to compute axion properties in GUTs, see Ref. [11].

By using Eqs. (5)–(6) and the relation $m_a = 5.7\,\text{neV}\,(10^{15}\,\text{GeV}/f_a)$ [38, 39] we can re-express Eq. (7) in the following parametric form:

$$\Gamma_{p\to\pi^0 e^+} \simeq \left(1.6\times 10^{34}\ \text{yr}\right)^{-1}\left(\frac{m_a}{3.7\ \text{neV}}\right)^4\left(\frac{6}{\hat{N}}\right)^4 \times\left[0.83\left(\frac{A_{SL}}{2.4}\right)^2+0.17\left(\frac{A_{SR}}{2.2}\right)^2\right], \quad (8)$$

where we have highlighted in the first parenthesis the current proton decay bound from SK [1]. Remarkably, this translates into an upper bound for the axion mass which, although affected by the model-dependent parameter $\hat{N}$, is independent of the fine details of the unification analysis that enter only logarithmically into $A_{SL(R)}$.

Axion mass prediction in WGG+24_F. The failure of the WGG model in explaining neutrino masses and gauge coupling unification can be readily fixed by adding a single Majorana representation, 24_F, in analogy to the proposal of Ref. [24]. Here, we highlight the main differences due to the presence of the PQ symmetry. The Yukawa Lagrangian is extended by

$$\Delta\mathcal{L}_Y = \overline{5}_F 24_F 5_H + \mathrm{Tr}\, 24_F^2 24_H^* + \text{h.c.}\,. \quad (9)$$

The first term provides a Dirac Yukawa interaction for the fermion triplet and singlet fields contained in 24_F, while the second term generates a Majorana mass for the full multiplet upon SU(5) symmetry breaking. We leave implicit the presence of extra non-renormalizable operators which are needed for two reasons: *i*) to avoid a rank-one light neutrino mass matrix and *ii*) to split the mass of the 24_F sub-multiplets (for further details see [24–26]). Eq. (9) also fixes the PQ transformation of the new field: $24_F \to e^{-i\alpha/2} 24_F$; including the latter the total $\mathrm{U}(1)_{\mathrm{PQ}}$-$\mathrm{SU}(3)_C$-$\mathrm{SU}(3)_C$ anomaly yields $\hat{N}=11$.

The possibility of narrowing down the axion mass range follows directly from unification constraints. The main issue with gauge coupling unification in the SM is the early convergence of the electroweak gauge couplings, α_1 and α_2, around 10^{13} GeV, at odds with proton decay bounds. Hence, the key ingredients for a viable unification pattern are additional particles charged under $\mathrm{SU}(2)_L$ which can delay the meeting of α_1 and α_2. Such a role in the WGG+24_F model can be played by the electroweak fermion $T_F = (1,3,0)$ and scalar $T_H = (1,3,0)$ triplets contained in the $24_{F,H}$.[5] They are predicted to be at the TeV scale, so that a large enough unification scale can be achieved.

Both types of triplets, if light enough, can give interesting signatures at the LHC. The fermionic component leads to same sign di-lepton events which violate lepton number [40]. A recent CMS analysis [41] sets a 95% CL exclusion at 840 GeV, while projected limits at the High Luminosity LHC (HL-LHC) [42, 43] give $m_{T_F} \gtrsim 2$ TeV. Bosonic triplets can affect the di-photon Higgs signal strength, but the bound is milder compared to the fermionic triplet and model-dependent [44]. Here we assume a conservative $m_{T_H} \gtrsim 200$ GeV.

The complete unification pattern including also the convergence of α_3 with α_1 and α_2 requires heavier colored particles. These are the color-octet fermions and scalars contained in the $24_{F,H}$, whose masses are required to be around 10^8 GeV, well beyond the LHC energy range.

The main prediction of gauge coupling unification is hence a clean correlation between a triplet mass parameter (whose analytical form is a consequence of the α_2 beta function),

$$m_3 = \left(m_{T_F}^4 m_{T_H}\right)^{1/5}, \quad (10)$$

and the unification scale. The latter is operatively defined as the energy scale where α_1 and α_2 meet up to GUT-scale thresholds [45, 46], and it can be identified with m_V, the mass of the heavy vector leptoquark V_μ mediating proton decay. Thanks to Eqs. (5)–(6), we can trade m_V for the axion mass, which allows us to present the unification constraints in the (m_a, m_3) plane.

Following Ref. [26], we have performed a gauge coupling unification analysis including the leading NNLO corrections coming from the 2-loop matching coefficients and the 3-loop beta functions due to the fermion and scalar triplets. The extra thresholds affecting the evolution of α_1 and α_2 are fixed in such a way that the value of m_3 is maximized (cf. [26] for more details), which defines the parameter $m_3^{\max}$. The results are displayed in Fig. 1 which shows the correlation in the $(m_a, m_3^{\max})$ plane. Taking into account the present bounds from LHC (on both fermion and scalar triplets) and SK (obtained by setting $A_{SL}=2.6$ and $A_{SR}=2.4$ in Eq. (8), which follow from the unification analysis), the preferred axion mass window is

$$m_a \in [4.8, 6.6]\ \text{neV}\,. \quad (11)$$

Future projections at HL-LHC (where we represent only the sensitivity to the fermion triplet mass) and HK (10 years data taking [2]) can complementary test this scenario.

We remark that the SK bound was imposed via Eq. (8), which does not account for possible cancellations in the flavour structure of the proton decay operators. By considering different proton decay channels and accounting for flavour rotations,

[5] Compared to the GG+24_F case we have in principle extra thresholds due to fact that the 24_H is complex. However, the constraints coming from the minimization of the scalar potential imply that only one *real* triplet can be light, otherwise a colored octet scalar would be lowered to the triplet mass scale, spoiling nucleosynthesis [24].

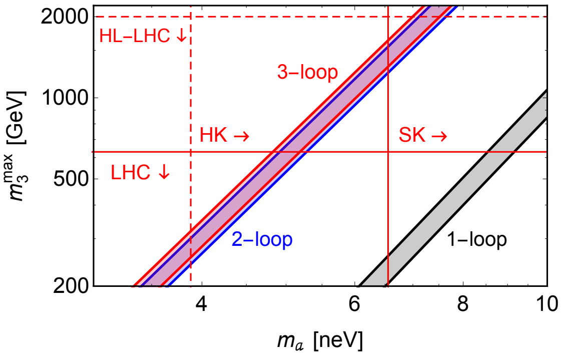

FIG. 1. Maximal triplet mass parameter as a function of the axion mass. Grey, blue and red bands denote respectively the correlation at 1, 2 and 3 loops (shaded regions encode the 1σ uncertainty on the electroweak gauge couplings). The full horizontal (vertical) red line is the current exclusion from LHC (SK), the dashed horizontal (vertical) red line is the projected exclusion from HL-LHC (HK).

one can still extract a model-independent bound on the unification scale which is about an order of magnitude smaller [47, 48]. The absolute upper bound on the axion mass is obtained by tuning to zero all the main proton decay channels, except those involving strange mesons. Using the results of Ref. [47] for the case of heavy Majorana neutrinos and updated with the latest experimental limit $\tau/\mathcal{B}(p \to K^0\mu^+) > 1.3 \times 10^{33}$ yr [49], we obtain $m_a < 330$ neV. Similarly, from the projections at HK (10 years data taking [2]) in the $p \to K^+\overline{\nu}$ channel we estimate $m_a < 160$ neV.

Sensitivity of future axion DM searches. An axion in this mass range is extremely weakly coupled to SM particles, since its couplings to e.g. photons (γ), electrons (e), protons (p), and neutrons (n) are inversely proportional to the axion decay constant,

$$\mathcal{L}_a \supset \frac{\alpha}{8\pi}\frac{C_{a\gamma}}{f_a}\, a\, F_{\mu\nu}\tilde{F}^{\mu\nu} - \frac{1}{2}\frac{C_{af}}{f_a}\,\partial_\mu a\;\overline{\Psi}_f\gamma^\mu\gamma_5\Psi_f\,. \tag{12}$$

while the coefficients C_{ax} are of order unity. In the WGG+24_F model, we find:

$$\begin{aligned} C_{a\gamma} &= \tfrac{8}{3} - 1.92(4)\,, \qquad C_{ae} = \tfrac{2}{11}\sin^2\beta\,, \\ C_{ap} &= -0.47(3) \\ &\quad + \tfrac{6}{11}[0.288\cos^2\beta - 0.146\sin^2\beta \pm 0.02]\,, \\ C_{an} &= -0.02(3) \\ &\quad + \tfrac{6}{11}[0.278\sin^2\beta - 0.135\cos^2\beta \pm 0.02]\,, \end{aligned} \tag{13}$$

where we introduced the ratio of the electroweak VEVs, $\tan\beta = \langle 5_H\rangle / \langle 5_{H'}\rangle$. This makes the GUT axion clearly invisible for purely laboratory based experiments.

However, axions in this mass range are known to be excellent DM candidates [7–9] which can be searched for in axion DM direct detection experiments. In fact, very light axion DM even tends to be overproduced and can only be reconciled with the measured amount of cold DM if the PQ symmetry remained broken during and after inflation in the early universe.[6] In this case, the relative contribution of axion DM to the energy density of the universe depends not only on the mass, but also on the initial value of the axion field a_i in units of the decay constant, $\theta_i = a_i/f_a$, inside the causally connected region which is inflated into our visible universe, cf. [39, 50]:

$$\Omega_a h^2 = 0.12\left(\frac{5.0\ \text{neV}}{m_a}\right)^{1.165}\left(\frac{\theta_{\rm i}}{1.6\times 10^{-2}}\right)^2. \tag{14}$$

Thus an axion in the neV mass range can make 100 % of DM, if the initial field value θ_i is of order 10^{-2}.[7] In this cosmological scenario, however, quantum fluctuations of a massless axion field during inflation may lead to isocurvature density fluctuations that get imprinted in the temperature fluctuations of the cosmic microwave background (CMB) [52, 53], whose amplitude is stringently constrained by observations. In the case that the 24_H stays at a broken minimum of the potential throughout inflation (e.g. for a SM-singlet inflaton), those constraints translate in an upper bound on the Hubble expansion rate during inflation [54–56]:

$$H_I < 5.7\times 10^8\ \text{GeV}\left(\frac{5.0\,\text{neV}}{m_a}\right)^{0.4175}. \tag{15}$$

Intriguingly, these isocurvature constraints can disappear completely in the case of non-minimal chaotic inflation [57–59] along one of the components of the 24_H. In this case, during inflation the 24_H is not at a minimum, Goldstone's theorem does not apply, and the lightest fluctuations orthogonal to the inflaton can have masses above H_I as long as the parameter ξ_{24_H}, describing the non-minimal coupling to the Ricci scalar, $S \supset -\int d^4x\,\sqrt{-g}\,\xi_{24_H}\,\mathrm{Tr}\,(24_H^2)R$, is larger than ~ 0.01. For ξ_{24_H} above this value, the power spectra of the isocurvature fluctuations become exponentially suppressed and the CMB bounds can be avoided. In such scenarios, one still needs to ensure that the PQ symmetry is never restored after inflation; we expect that this might be possible for small enough quartic and Yukawa couplings of the 24_H, but a dedicated analysis generalizing the non-pertubative and perturbative reheating calculations in Ref. [50] is needed.

The DM experiment ABRACADABRA [27], has very good prospects to probe the axion photon coupling, $g_{a\gamma} = \alpha\, C_{a\gamma}/(2\pi f_a)$, in the relevant mass region. This is shown in Fig. 2, from which we infer

[6] This solves at the same time the cosmological SU(5) monopole problem and the PQ domain-wall problem (the WGG+24_F model has domain-wall number 11).

[7] This value can be supported by anthropic arguments [51].

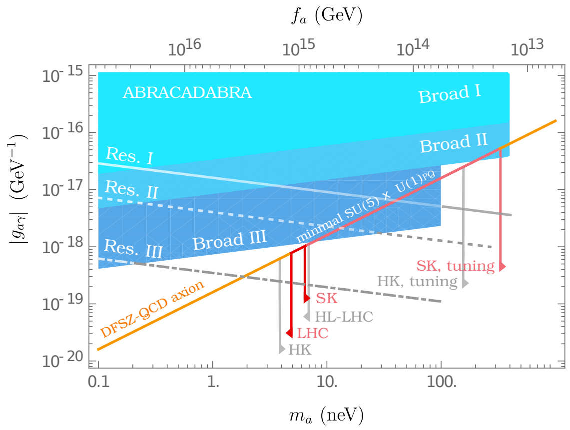

FIG. 2. Axion coupling to photons, $g_{a\gamma}$, versus axion mass m_a. The blue regions give the projected sensitivities of broadband ("Broad") and resonant ("Res.") search modes of ABRACADABRA from Ref. [27].

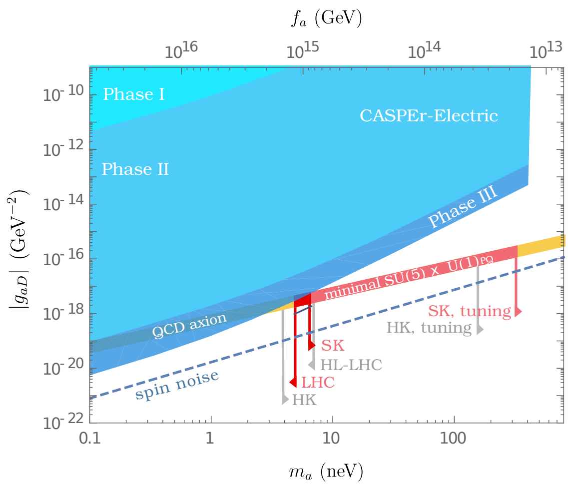

FIG. 3. Axion coupling to the nucleon EDM operator, g_{aD}, versus axion mass m_a. The blue regions give the projected sensitivities of CASPEr-Electric from Ref. [29]. The short, full blue line reflects a factor of three improvement in sensitivity for a search just concentrated on the preferred mass region.

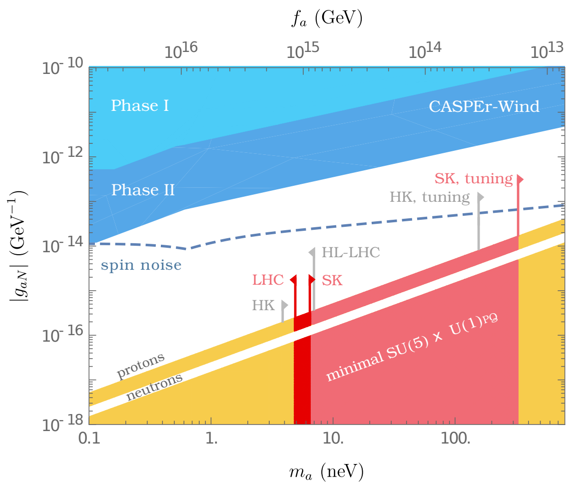

FIG. 4. Axion coupling to the nucleons, g_{aN}, versus axion mass m_a. The blue regions give the projected sensitivities of CASPEr-Wind from Ref. [29].

that the whole parameter space of the WGG+24_F model (including the tuned region) can be tested in the third phase of the broadband and resonant search modes of ABRACADABRA.

In Fig. 3, we confront our axion mass prediction with the projected sensitivity of the experiment CASPEr-Electric [28, 29], which aims to search for oscillating nucleon electric dipole moments (EDM) $d_n(t) = g_{aD} \frac{\sqrt{2\rho_{\rm DM}}}{m_a} \cos(m_a t)$ [60], where g_{aD} is the model-independent coupling of the axion to the nucleon EDM operator, $\mathcal{L}_a \supset -\frac{i}{2} g_{aD}\, a\, \overline{\Psi}_N \sigma_{\mu\nu}\gamma_5 \Psi_N F^{\mu\nu}$, and $\rho_{\rm DM} = 0.3\,{\rm GeV/cm}^3$ is the local energy density of axion DM. The QCD axion band in Fig. 3 indicates the theoretical uncertainty of the non-perturbative estimates of g_{aD}. We used the result in [61], obtained with QCD sum rules; for other evaluations see e.g. [62, 63].[8] We infer from Fig. 3, that the preferred axion mass window (11) could definitely be probed in phase III of CASPEr-Electric.[9]

On the other hand, the projected sensitivity of CASPEr-Wind [29], which exploits the axion nucleon coupling $g_{aN} = C_{aN}/(2 f_a)$ $(N = p, n)$ to search for the axion DM wind due to the movement of the Earth through the Galactic DM halo [60], misses the preferred coupling vs. mass region by two orders of magnitude or more, even in its phase II. We show this in Fig. 4, where the theoretical uncertainty of the axion band is obtained from the errors in the coefficients of Eq. (13), and from varying $\tan\beta \in [0.28, 140]$ in the perturbative unitarity domain [65].

Conclusions. In this Letter we have proposed a minimal implementation of the PQ mechanism in a realistic SU(5) model, which predicts a narrow axion mass window (cf. Eq. (11)) which can be directly tested at future axion DM experiments and indirectly probed by collider and proton decay experiments. In principle, a precise determination of m_a (via ABRACADABRA and/or CASPEr-Electric) would lead to a direct determination of the GUT scale, possibly discriminating among GUT models, and setting a target for proton decay measurements. Although we exemplified our predictions in the case of the WGG+24_F model, it would be interesting to compare axion properties in other minimal extensions of the WGG model which can simultaneously address neutrino masses and gauge coupling unification (see e.g. [66, 67]), or in realistic SO(10) models [11].

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[8] Current lattice QCD results on g_{aD} do not show a statistically significant non-zero signal [64].

[9] The sensitivity in g_{aD} improves with the scanning time as $t^{1/4}$. This amounts to a factor of three improvement (denoted by a short, full blue line in Fig. 3), if CASPEr-Electric spends all the measurement time just on the preferred mass region.

Finally, the intriguing possibility that the 24_H field could also be responsible for inflation would make the WGG+24_F model a potential candidate for a minimal and predictive GUT-SMASH [50, 68] variant aiming at a self-contained description of particle physics, from the electroweak scale to the Planck scale, and of cosmology, from inflation until today. We leave a detailed investigation of this scenario for future studies.

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