Accidental Peccei-Quinn symmetry protected to arbitrary order

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A SU(N) × SU(N) gauge theory for a scalar multiplet Y transforming in the bi-fundamental representation (N, N) preserves, for N > 4, an accidental U(1) symmetry firstly broken at operator dimension N. Two configurations are possible for the vacuum expectation value of Y, which correspond to the (maximal) little groups Hν = SU(N) × SU(N−1) and Hν = SU(N−1) × SU(N−1) × U(1). In the first case the accidental U(1) gets also broken, yielding a pseudo Nambu-Goldstone boson with mass suppression controlled by N, while in the second case a global U(1) remains unbroken. The strong CP problem is solved by coupling Y to new fermions carrying color. The first case allows for a Peccei-Quinn solution with U(1)PQ protected up to order N by the gauge symmetry. In the second case U(1) can get broken by condensates of the new strong dynamics, resulting in a composite axion. By coupling Y to fermions carrying weak isospin, models for axion-like particles can be constructed.

Introduction. In the past decades, a plethora of experimental results has firmly established QCD as the correct description of strong interaction phenomena in particle physics. However, together with a deep understanding of many fundamental issues, this beautiful theory also brings in one theoretical conundrum. The QCD gauge sector depends on two dimensionless parameters whose value is not predicted by the theory, but must be determined experimentally. The first one, α, determines the strength of the QCD interactions. Its measured value is a natural one for a dimensionless quantity (roughly speaking is of order unity, although the precise number depends on the energy scale). The second one, θ, gives the amount of CP violation in strong interactions. Theory only dictates that θ, which is an angular variable, must fall within the interval [0, 2π], and also in this case it would be natural to expect θ ∼ O(1). Instead, experimental limits on the neutron electric dipole moment yield the upper bound |θ| < 10−10, a value that is regarded as highly unnatural. This theoretical nuisance bears the name of “the strong CP problem”. QCD, however, would recover its naturalness if, for some reason, θ = 0.

An elegant mechanism to guarantee the vanishing of θ was proposed in 1977 by Peccei and Quinn (PQ) [1,2]. It relies on a U(1)PQ global symmetry, anomalous with respect to QCD, broken spontaneously by the vacuum expectation value (VEV) of a Standard Model (SM) singlet scalar field at a scale vα ≫ 100 GeV, and broken explicitly by non-perturbative QCD effects at a scale ΛQCD ∼ 100 MeV. Spontaneous breaking (SB) of a global U(1) symmetry gives a massless Nambu-Goldstone boson (NGB). However, due to the presence of a relatively tiny explicit breaking, the NGB of U(1)PQ is not exactly massless: it is a pseudo NGB, commonly referred to as the axion [3,4]. To account for |θ| < 10−10 any other source of explicit U(1)PQ breaking besides QCD must either be absent or adequately suppressed. This is difficult to achieve, especially considering that U(1)PQ, being anomalous, is not even a real symmetry. Thus, effective operators not respecting U(1)PQ are expected to arise. Even if suppressed by the Planck scale mP = 1.2 · 1019 GeV, unless their dimension is larger than d ∼ 10 they would unavoidably give |θ| > 10−10 [5–8].

In this Letter we propose a mechanism that, on the basis of first principles, can protect U(1)PQ to arbitrary accuracy. A scalar multiplet Y is assigned to the bi-fundamental representation (N, N) of the gauge group GLR(N) = SU(N) × SU(N)R. For N > 4 an accidental U(1) symmetry corresponding to phase transformations of Y is enforced at the classical level, and it only gets broken at d = N by the determinant interaction det(Y). When the scalar gauge theory is coupled to fermions carrying color, U(1) acquires a QCD anomaly. SB of GLR(N) via a VEV of Y can proceed via two patterns. In the first case U(1) also undergoes SB, acquiring all the features of a PQ symmetry: θ becomes a fundamental dynamical field with a periodic potential that drives its value to zero. In the second case a global U(1) remains perturbatively unbroken. However, condensates of the new strong gauge dynamics can break it, giving rise to a composite axion. In both cases a solution to the strong CP problem is obtained.

Accidental U(1) in GLR(N) scalar gauge theory.

For N > 4, GLR(N) gauge invariance restricts the renormalizable potential for the scalar multiplet Y to the following simple form:

\[ V_0 = \lambda \left[ T - v^2_A / 2 \right]^2 + \lambda_A A, \]

(1)

where vα is a constant with the dimension of a mass, T = Tr[YY†] and A = Tr[Mn(YY†, 2)], with Tr[Mn(M, k)] denoting the trace of the matrix of the minors of order k of M [9]. We require vA > 0.
to trigger SB and $|\lambda_A| < \frac{2N}{m} \lambda$ to ensure that the potential is bounded from below. The matrix $Y^c$ of constant background values of $\bar{Y}(x)$ can be written in its singular value decomposition as:

$$\frac{\sqrt{2}}{v_a} Y^c = U_L \tilde{Y} U_R^t = U_L \left( \tilde{\Phi} \tilde{Y} \right) U_R^t \rightarrow \Phi \tilde{Y},$$

(2)

where $\tilde{Y} = \text{diag}(y_1, y_2, \ldots, y_N)$ is diagonal with real non-negative entries normalized such that $\sum y_i^2 = 1$, $U_L, U_R$ are unitary matrices, $U_{L,R}$ are special unitary $(\det (U_{L,R}) = +1)$, $\Phi$ is a diagonal matrix of phases such that $\log \det (\Phi) = \log \det (U_L U_R^t) = i \arg \det (Y^c) = i \delta^c$, being an angular variable, ranges in the interval $[0, 2\pi]$, and physics must be invariant under the redefinition $\delta^c \rightarrow \delta^c + 2n\pi$ $(n = 1, 2 \ldots)$. The last (diagonal) form in Eq. (2) is obtained, without loss of generality, via a rigid $G_{LR}^c$ rotation. We will leave understood that $\tilde{Y}$ is always written in this basis as $Y^c = (v_a/\sqrt{2}) \Phi \tilde{Y}$.

The vacuum configurations that minimize $V_0$ are easily found \[10\]: $T$ is blind to specific $\bar{Y}$ configurations (this is because it carries a $SO(2N^2)$ accidental symmetry much larger than the gauge symmetry that allows to rotate between different configurations). Minimization of the first term then just fixes the “length” $t(\tilde{Y}) = v_a^{-1} \sqrt{2/T} = 1$. The extrema of $(A) \propto \sum_{i<j} y_i^a y_j^a$ instead depend on the structure of $\tilde{Y}$. We have two possibilities: (i) for $\lambda_A < 0$, $(A)$ is maximized at the symmetric point $y_i^a = 1/N$, $\forall i$; (ii) for $\lambda_A > 0$ the minimum occurs when $(A) = 0$, that is when all entries in $\tilde{Y}$, but one, vanish. In summary, the configurations that extremize $V_0$ are:

(i) $Y_s^c = \frac{v_a}{\sqrt{2}} \tilde{Y}_s$, $\tilde{Y}_s = \frac{1}{\sqrt{N}} \text{diag} (1, 1, \ldots, 1)$ ,

(ii) $Y_h^c = \frac{v_a}{\sqrt{2}} \tilde{Y}_h$, $\tilde{Y}_h = \text{diag} (0, 0, \ldots, 0, 1)$ .

(3)

The corresponding little groups are the two maximal subgroups of $G_{LR}^{(N)}$, $H_a = SU(N)_{L+R}$ for $Y_s^c$ and $H_h = G_{LR}^{(N-1)} \times U(1)_{L+R}$ for $Y_h^c$, where the Abelian generator corresponds to the diagonal combination of the LR Cartan generators of $G_{LR}^{(N)}$ proportional to $\lambda_{N-1}^L \approx (1, 1, \ldots, 1 - N)_{L+R}$. It is important to stress that $H_a,h$ cannot get broken further by any type of perturbative effects \[11\] or, equivalently, that neither the vanishing entries in $Y_h^c$ can be perturbatively lifted, nor the strict equality of the entries in $Y_s^c$ can be perturbatively spoiled. Some bibliographic remarks are in order: the minima of the potential for the case of global $G_{LR}^{(N)}$ (namely the SM quark flavor symmetry) were studied in \[10\] (and with the assumption of a real $\det (\bar{Y})$ previously in \[12\]). The possibility of raising perturbatively the vanishing entries in $Y_h$ to generate the SM fermion mass hierarchy was addressed in \[13\]. It was found that, in agreement with the Georgi-Pais (GP) theorem \[11\], minimization of the effective potential results in the same little groups $H_{a,h}$. Only by introducing additional reducible scalar representations a more thorough breaking, yielding $y_{i\neq j} \neq y_j$, can be obtained \[13\] \[14\].

The tree level potential $V_0$ in eq. (1) has an accidental $U(1)$ rephasing symmetry $Y \rightarrow e^{i(\alpha/N)} Y$ (under which $\delta^c \rightarrow \delta^c + \alpha$) so that the full symmetry of the classical Lagrangian is in fact $G_{LR}^{(N)} \times U(1)$. The first minimum $Y_s^c$ breaks this symmetry and yields a NGB, which, in first approximation, remains massless. However, accidental symmetries are generally not respected by operators of higher dimensions. Here it is the requirement of local gauge invariance that dictates at which order these operators can arise. A fundamental set of higher order operators can be constructed by considering the characteristic polynomial $P(\xi)$ of the matrix $YY^c$:

$$P(\xi) = \det (\xi I - YY^c) = \sum_{n=0}^{N} (-1)^n C_n \xi^{N-n},$$

(4)

where $I$ is the identity matrix, and $C_n = \text{Tr} [\text{Mtr} (YY^c)^n]$, with $C_0 = 1$, $C_1 = T$, $C_2 = A$, $\ldots$, $C_N = \det [YY^c] = |D|^2$. The solutions of $P(\xi) = 0$ are the eigenvalues of $YY^c$ and, being the eigenvalues invariant under $G_{LR}^{(N)}$, so are the coefficients $C_n$. They correspond to invariant combinations of components of $Y$ of dimension $d = 2n$ \[13\]. The determinant $D = \det Y$ is another invariant, since under $G_{LR}^{(N)}$, $D \rightarrow \det (V_L Y V_R^t) = \det Y$ \[15\]. However, while all $C_n$’s respect the $U(1)$ accidental symmetry, under $Y \rightarrow e^{i(\alpha/N)} Y$, $D(x) \rightarrow e^{i\alpha} D(x)$. Thus, $U(1)$ gets first broken at $d = N$ by:

$$V_D = \frac{k D + k^2 D^*}{m_{N-4}^2} = \frac{2\kappa D}{m_{N-4}^2} \cos[v + \delta(x)],$$

(5)

where $\kappa$ and $\varphi$ are the modulus and argument of the coupling $k$, $D = |D|$, $\delta(x) = \arg D(x)$, and the $m_{N-4}^2$ suppression stems from the assumption that $V_D$ is only generated by gravity effects. Let us see which is the fate of the NGB of case (i). The minimum of $V_D$ is obtained for $\delta(x) \equiv \delta = \pi - \varphi$, and the minimum of $V_0$ is lowered by the amount:

$$\Delta V = \frac{v_a^2}{2N} \frac{2\kappa}{(2N)_{N-2}} \left( \frac{v_a}{m_{N-4}} \right)^{N-4} .$$

(6)

Thus, in the breaking $G_{LR}^{(N)} \times U(1) \rightarrow H_a$, $N^2 - 1$ of the initial $2(N^2 - 1) + 1$ generators are left unbroken, $N^2 - 1$ are spontaneously broken with the corresponding NGB eaten by gauge bosons that acquire masses $O(v_a)$, while the NGB of the global $U(1)$ acquires a tiny mass $O(\sqrt{\Delta V/v_a})$ because of the explicit breaking Eq. (5). In case (ii) instead, a global $U(1)'$ generated by $\lambda_{N-1}^L + (N - 1)I$ is preserved by $Y_h^c$, so that at the renormalizable level $G_{LR}^{(N)} \times U(1) \rightarrow H_h \times U(1)'$. Although the higher order operator Eq. (5) breaks $U(1)'$ in interactions, since in the ground state $\langle D \rangle = 0$, there is no SB
and no NGB arises. This case, in which the symmetry of the effective Lagrangian is smaller than the symmetry of the renormalizable Lagrangian, but the resulting little group is the same, provides a non trivial test of the GP theorem [11].

**Solutions to the strong CP problem.** Solutions to the strong CP problems can be implemented by introducing fermions carrying color. Two different types of solutions are possible, depending on which vacuum is selected in the breaking of $G_{\text{LR}}^{(N)}$. Let us proceed by steps by first introducing four fermion multiplets transforming under $G_{\text{LR}}^{(N)}$ as: $Q_L \sim (N, 1), Q_R \sim (1, N), \Psi_L \sim (N, 1)$ and $\Psi_R \sim (1, N)$. Since they can be combined into real representations of $SU(N)_{\text{LR}} (Q_{L,R} \otimes \Psi_{L,R})$, there are no gauge anomalies. Gauge symmetry allows for Yukawa couplings of the form $\overline{Q}_L Y Q_R + \overline{\Psi}_L Y^\dagger \Psi_R + \text{H.c.}$. They preserve the Y rephasing symmetry if the fermions are transformed chirally with $U(1)$ charges satisfying: $\lambda Q_L - \lambda Q_R = \lambda Y$ and $\lambda \Psi_L - \lambda \Psi_R = -\lambda Y$. The opposite sign of the two charge differences (a consequence of the requirement of gauge anomalies cancellation) ensures the absence of $U(1)-SU(N)_{\text{LR}}$ mixed anomalies [16]. Let us now triplicate the fermion content, and assign $Q_{L,R}$ to the fundamental representation of color, while $\Psi_{L,R}$ ($a = 1, 2, 3$) remain color singlets. Since there is no compensating cancellation of the $Q_{L,R}$ contribution, a $U(1)$-QCD anomaly arises, and $U(1)$ acquires all the features required for a PQ symmetry.

(i) **Solution with a fundamental axion.** We choose a basis in which the SM quark masses are real while $\theta_{QCD} \neq 0$. Without loss of generality the $\Psi_{L,R}$ couplings can be taken flavor-diagonal, so that the Yukawa terms can be written as:

$$ e^{i\eta_a/N} h_0 \overline{Q}_L Y Q_R + e^{i\eta_a/N} h_a \overline{\Psi}_L Y^\dagger \Psi_R, \quad (7) $$

where $h_0, h_a$ ($a = 1, 2, 3$) are four real non-negative parameters. If $\lambda_A < 0$, the minimum $Y^\dagger = \eta_0 + \delta^c$ is selected, and all the fermions become massive, with degenerate masses within each $SU(3)_c$ and $SU(N)_{L+R}$ multiplet. We first show that the fermion masses stemming from eq. (7) can be brought into real form without inducing mixed $G_{\text{LR}}^{(N)}$ anomalies. After SB of the gauge symmetry, arg det $(M_2) = \eta_0 + \delta^c$ and arg det $(M_3) = \eta_0 - \delta^c$, that is there are four independent phases that we wish to cancel (four conditions). We can perform four chiral rotations of the fermion multiplets respectively with phases $\alpha_0, \alpha_a$, subject to a fifth condition $\sum_{a=1}^3 |\alpha_a| = 3 \alpha_0$ which avoids mixed anomalies with the $SU(N)$ gauge groups. The phase of $Y$ can be also redefined (this changes the argument of the cosine in eq. (5) by the addition of a constant term $\varphi \to \tilde{\varphi}$). All the complex phases can thus be canceled. However, the chiral rotation of $Q_{L,R}$ is anomalous with respect to $SU(3)_c$ of color, and another source of explicit $U(1)$ breaking is then introduced. Including it, the relevant potential for $\delta(x)$ acquires the form:

$$ V_\delta = \Delta V \cos[\tilde{\varphi} + \delta(x)] - f_\delta^2 m_a^2 \cos[\delta(x)], \quad (8) $$

where $\tilde{\varphi}$ is a generic constant unrelated with $\theta_{QCD}$, and we have redefined $\delta(x) + \theta_{QCD} \to \delta(x)$ so that the anomalous coupling to the gluons can be written as $\frac{a}{\pi} \delta GG$. Note that when the angular variable $\delta(x)$ is varied in the interval $[0, 2\pi)$, there is a unique minimum of the potential. Namely, independently of $N$, the number of domain walls is always one. From Eq. (5) we see that if $\alpha_c, \varphi = O(1)$, as it is natural to assume, $\delta^c < 10^{-10}$ can be ensured only if (the gravitationally induced) explicit breaking satisfies $\Delta V/(f_\delta^2 m_a^2) \lesssim 10^{-10}$. For the phenomenologically preferred interval $\Delta V \approx (f_\delta m_a^2) \lesssim 10^{-10}$, the condition can be fulfilled with $9 \leq N \leq 13$ [18].

Let us now proceed to identify the fundamental axion field and some of its properties. In the “unitary” gauge in which the rigid $G_{\text{LR}}^{(N)}$ rotation yielding $Y^\dagger = \Phi Y$ in Eq. (2) is promoted to a local one, we can write $Y(x) = \Phi(x) Y(x)$ with

$$ \Phi(x) = \text{diag} (e^{i\delta_1(x)}, \ldots, e^{i\delta_N(x)}); \quad \tilde{\gamma}_i = \sqrt{\frac{2N}{v_a}} \gamma_i. \quad (9) $$

The linear combinations of the $N$ “orbital” modes $\tilde{\gamma}_i$ corresponding to $N - 1$ non-Abelian broken generators and to the accidental $U(1)$ are:

$$ a_a(x) = 2 \text{Tr} [\gamma_i(x) \cdot T_a] \quad (10) $$

where $\gamma = (\gamma_1, \ldots, \gamma_N)$ and, for $a = 1, \ldots, N - 1, T_a$ are the $SU(N)$ Cartan generators with normalization $\text{Tr}[T_a]^2 = 1/2$, while $T_0 = (1/\sqrt{2N}) I$. Then, the canonically normalized axion field is:

$$ a_0(x) = 2 \text{Tr} [\gamma \cdot T_0] = \frac{v_a}{N} \delta(x). \quad (11) $$

Note that since the periodicity of $\delta(x)$ is $2\pi$, the periodicity of the axion field is $a_0 \to a_0 + \frac{2\pi}{N} v_a$. One might wonder whether, contrary to what stated below Eq. (5), there are $N$ domain walls corresponding to the $N$ minima $\langle a_0 \rangle + \frac{2\pi}{N} v_a$, $n = 0, \ldots, N - 1$. This is not so: all these minima are in fact gauge equivalent, in the sense that the $Z_N$ center of $SU(N)_{L+R}$ has precisely as elements $\exp(2\pi i n/N) \cdot I$, so that the cyclic values of $a_0/v_a$ can be all connected via gauge transformations. Neglecting the subdominant gravitational contributions, the mass of the axion is $m_a = N (m_a f_a)/v_a$, while the strength of its coupling to the photon via the usual term $(1/4) g_{\varphi \gamma \gamma} F_{\mu \nu} F^{\mu \nu}$ is:

$$ g_{\varphi \gamma \gamma} = -1.92 \frac{m_a}{eV} \frac{2.0}{10^{10} \text{GeV}}, \quad (12) $$

which falls within the axion window, see Fig. [4].

(ii) **Solution with a composite axion.** If in eq. (1) $\lambda_A > 0$, the minimum $Y_f^\dagger$ provides mass for just one fermion in each $N$-dimensional multiplet, $12(N - 1)$
Weyl fermions remain massless, and a global $U(1)'$ acting on these massless fermions remains unbroken. It is then tempting to try to implement a massless-quark type of solution, where the topological term is simply removed via a chiral $U(1)'$ rotation without other consequences on the Lagrangian. As regards the massless $Q$’s and $\Psi$’s, they would get confined into $F$-hadrons once $G_{LR}^{(N-1)}$ enters the strongly coupled regime at a scale $\Lambda_F \gg \Lambda_{QCD}$. However, such a scenario is not viable because matching the $U(1)'$-QCD anomaly in the low energy theory [19] requires that some composite fermions carrying color remain massless, and these are not observed. We are then led to assume that $U(1)'$ gets spontaneously broken by some color neutral condensate of $Q$ and $\Psi$. In this case the pseudo NGB of the $U(1)'$ symmetry would correspond to a composite axion [20][22] with mass and couplings suppressed as $1/\Lambda_F$. Clearly, this second scenario is more speculative, so that we will mainly focus on the first scenario.

**Phenomenology.** In the first scenario with a fundamental axion, after SB the spectrum consists of:

- $N^2 - 1$ gauge bosons with masses $O(v_a)$;
- $N$ quarks $Q$ and 3N SM singlet fermions $\Psi^a$, stable at the tree level, with $m_{Q,\Psi} \sim O(v_a/\sqrt{N})$;
- $N^2 - 1$ massless gauge bosons $F$;
- One pseudo NGB (the axion).

Tree level stability of the heavy $Q$’s and $\Psi$’s follows from the fact that they do not carry weak isospin and hypercharge $Y$, and thus cannot decay into lighter SM fermions. Different representations allowing for couplings between the $Q$’s and the SM fermions would bring the problem of canceling $[G_{LR}^{(N)}]^2 \times U(1)_Y$ gauge anomalies, a complication that we prefer to avoid. However, cosmologically stable heavy relics, and in particular long lived strongly interacting particles like the $Q$’s, represent serious issues in cosmology and astrophysics (see [23][24] for a recent discussion and relevant references). A simple way to avoid all phenomenological problems is to assume a pre-inflationary scenario (PQ symmetry broken before inflation) and $v_a > T_{\text{reheating}}$ so that, after inflation has wiped away all the heavy states, they cannot be regenerated. This holds, in particular, for the heavy $Q$ which otherwise would be copiously produced via QCD interactions. As regards the massless gauge bosons $F$, their production after inflation could proceed via gravitational effects, or via gluon-gluon scattering into a pair of $F$’s mediated by loops of heavy $Q$’s. The first mode is suppressed by powers of $m_F$ and typically very inefficient. The rate for the second process can be estimated as $(\alpha_F \alpha_a)^2 T^9/m_{Q}^8$, with $\alpha_F$ the coupling strength of the new gauge group. This reaction remains well out of equilibrium for all $T < m_Q$ and thus also this channel is typically too inefficient to produce $F$ in sizable amounts. All in all, we can conclude that in pre-inflationary scenarios no dangerous heavy relics are left in appreciably amounts.

Post-inflationary scenarios (PQ symmetry broken after inflation, and $v_a < T_{\text{reheating}}$) yield a different picture. During reheating, all the heavy states can attain equilibrium distributions. At $T \lesssim v_a$ the heavy gauge bosons with masses $O(v_a)$ will readily decay into the lighter $Q, \Psi$ fermions. Below $m_Q$, the unbroken $SU(N)_{LR}$ corresponds to an effective pure Yang-Mills theory with large $N$ that rapidly flows towards a confining regime at a scale $\Lambda_F \gg \Lambda_{QCD}$. Bound state mesons $\Pi_{Q,\Psi} \sim Q\bar{Q}\, (\Psi\bar{\Psi})$ singlets under $SU(N)_{LR}$ form, and readily decay into lighter gauge bosons $G \sim FF$ of mass $m_G \sim O(\Lambda_F)$ (the SM must be also color singlets). These ‘gaugeballs’ are easily disposed of: they can decay into two gravitons with a rate $\Gamma_{grav} \sim \Lambda_F^4/m_F^4$ [23], or into a pair of gluons via heavy quark loops with a rate $\Gamma_{qg} \sim \Lambda_F^9/m_Q^8$. If $m_q/m_F \lesssim \Lambda_F/m_Q$, decays into gluons proceed faster than decays into gravitons. However, visible decays are subject to severe constraints from big bang nucleosynthesis and from several other observations (limits on CMB spectral distortions, diffuse $\gamma$-rays background, etc.) which, taken together, suggest $\tau_G \lesssim 10^{-2} s$. Rewriting the lifetime as $\tau_G \sim 10^{-33} \left(10^9 \text{GeV}/m_Q \right) \left(m_Q/\Lambda_F \right)^9$ s, we see that safe lifetimes are obtained if the heavy quark mass and confinement scale fulfill $m_Q/\Lambda_F \lesssim 3 \times 10^9$, a condition that can be easily satisfied for $N \gtrsim 9$. However, there are other more dangerous heavy relics, like $M_{ab} \sim \Psi_a \bar{\Psi}_b$ with $(a \neq b)$ and, since at $\Lambda_F$ color is unconfined, ‘mongrel’ mesons $M_{ab} \sim Q\bar{Q}$ will also form. $M_{ab}$ decays are forbidden by $\Psi$-flavor conservation, while decays of $M_a$ are forbidden also by color conservation. The abundance of these states is basically determined by free particle annihilation before $F$-confinement, which always results in $\Omega_M \gg \Omega_{DM}$ unless their overall mass scale is brought down to values not much larger than a few TeV. This requires tiny values for the Yukawa couplings in eq. (7) and an appropriately small initial values for the $G_{LR}^{(N)}$ gauge couplings to ensure $\Lambda_F \lesssim 2$ TeV. All in all, the post-inflationary scenario, if not ruled out, is certainly strongly disfavored with respect to pre-inflationary scenarios.

**Axion like particles.** With no attempt to solve the strong CP problem, models for axion like particles (ALPs) can be constructed along the same lines. Instead than new quarks, let us introduce new colorless fermions $T$, doublets under weak-isospin and, for simplicity, with zero hypercharge. SM singlet fermions $\Psi_{LR}^a (a = 1, 2)$ are also introduced to cancel gauge anomalies. Now, in the breaking $G_{LR}^{(N)} \rightarrow \mathcal{H}_a$ the NGB of the accidental $U(1)$ only receives mass from the gravity induced determinant operator of $d = N$. However, here $N$ does not need to be particularly large so that, compared to axions,
much larger ALP masses and photon couplings are possible. For the ALP mass we obtain:

\[ m_a = \frac{N \sqrt{2} \kappa}{(2N)^{N/2}} \left( \frac{v_a}{m_P} \right)^{N/4} v_a, \tag{13} \]

and for the ALP-photon coupling we have:

\[ g_{a\gamma\gamma} = \frac{\alpha}{2\pi} \frac{E}{v_a}, \tag{14} \]

with the electromagnetic anomaly coefficient:

\[ E = 2N \mathcal{X}_{T_L} Q_{L+}^4 = \frac{N}{2}, \tag{15} \]

where \( Q_{L+} = \pm \frac{1}{2} \) are the electric charges of the components of \( T_L \), and \( \mathcal{X}_{T_R} = 0, \mathcal{X}_{T_L} = \mathcal{X}_Y = 1 \) the (assigned) PQ charges. Eqs. (13)–(15) yield

\[ g_{a\gamma\gamma} = B_N m_a^{-\frac{1}{N - 2}}, \tag{16} \]

with

\[ B_N = \frac{N \alpha}{4\pi} \left( \frac{2\kappa N^2}{(2N)^{N/2} m_P^{N/4}} \right)^{\frac{1}{N - 2}}. \tag{17} \]

The ALP-photon coupling versus \( m_a \) is plotted in Fig. 1 for different values of \( N \). The preferred regions for the axion are \([23, 24]\) as also shown.

**Conclusions.** We have put forth a new realization of the PQ solution to the strong CP problem. Our scenario might be loosely classified as a KSVZ type of axion model \([26, 27]\) since PQ charges are assigned only by non-SM particles. A new gauge group \( G_{LR}^{(N)} = SU(N)_L \times SU(N)_R \) is postulated, new quarks \( Q_{L,R} \) are assigned to fundamental representations of \( SU(N)_L \) while the PQ scalar, rather than being a single complex field, is a matrix \( Y \) transforming in the bi-fundamental representation of \( G_{LR}^{(N)} \). A PQ symmetry arises accidentally, and remains protected by the gauge symmetry from all types of explicit breaking up to dimension \( N \), which is in principle arbitrary. Within this same construction, and depending on the gauge symmetry breaking pattern, a different possibility where the axion is composite can be realized.

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**FIG. 1.** ALP-photon coupling versus mass for different choices of the gauge group \( G_{LR}^{(N)} \). The green band and the yellow area represent two regions for phenomenologically preferred axion models \([23, 24]\). The axion coupling in Eq. (13) corresponds to the thick blue line.

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[1] R. D. Peccei and H. R. Quinn, Phys. Rev. D16, 1791 (1977).
[2] R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. 38, 1440 (1977).
[3] S. Weinberg, Phys. Rev. Lett. 40, 223 (1978).
[4] F. Wilczek, Phys. Rev. Lett. 40, 270 (1978).
[5] H. M. Georgi, L. J. Hall, and M. B. Wise, Nucl. Phys. B192, 409 (1981).
[6] M. Kamionkowski and J. March-Russell, Phys. Lett. B282, 137 (1992), arXiv:hep-th/9202003 [hep-th].
[7] R. Holman, S. D. Hsu, T. W. Kephart, E. W. Kolb, R. Watkins, and L. M. Widrow, Phys. Lett. B282, 132 (1992), arXiv:hep-ph/9203206 [hep-ph].
[8] S. M. Barr and D. Seckel, Phys. Rev. D46, 539 (1992).
[9] The reader might be more familiar with the double trace \( \text{Tr}[Y Y^\dagger Y Y^\dagger] \) as a second invariant.
[10] E. Nardi, Phys. Rev. D84, 036008 (2011), arXiv:1105.1770 [hep-ph].
[11] H. Georgi and A. Pais, Phys. Rev. D16, 3520 (1977).
[12] N. Cabibbo and L. Maiani, in: *Evolution of particle physics*, edited by M. Conversi, Vol. 50 (Academic Press, New York, 1970) pp. 68-72 (App. I).
[13] J. R. Espinosa, C. S. Fong, and E. Nardi, JHEP 02, 137 (2013), arXiv:1211.6428 [hep-ph].
[14] C. S. Fong and E. Nardi, Phys. Rev. D89, 036008 (2014), arXiv:1307.4412 [hep-ph].
[15] All other invariants can be expressed in terms of the set \( \{ C_1(x), C_2(x), \ldots, C_{N-1}(x), P \} \), see [13].
[16] This point is important for the PQ solution: it ensures that the VEV of the axion will not end up canceling the \( SU(N)_L \) \( \theta \)-terms rather than \( \theta_{QCD} \).
[17] P. Sikivie, Phys. Rev. Lett. 48, 1156 (1982).
For $N > 10$, QCD asymptotic freedom is lost beyond the heavy quarks' threshold. However, Landau poles in $\alpha_s$ remain trans-Planckian for $N \lesssim 26$.

G. ’t Hooft, in Recent Developments in Gauge Theories. Proceedings, Nato Advanced Study Institute, Cargese, France, August 26 - September 8, 1979, Vol. 59, edited by G. ’t Hooft, C. Itzykson, A. Jaffe, H. Lehmann, P. K. Mitter, I. M. Singer, and R. Stora (1980) p. 135.

J. E. Kim, Phys. Rev. D31, 1733 (1985)

L. Randall, Phys. Lett. B284, 77 (1992)

M. Redi and R. Sato, JHEP 05, 104 (2016) arXiv:1602.05427 [hep-ph]

L. Di Luzio, F. Mescia, and E. Nardi, Phys. Rev. Lett. 118, 031801 (2017) arXiv:1610.07593 [hep-ph]

L. Di Luzio, F. Mescia, and E. Nardi, In preparation.

A. Soni and Y. Zhang, Phys. Rev. D93, 115025 (2016) arXiv:1602.00714 [hep-ph]

J. E. Kim, Phys. Rev. Lett. 43, 103 (1979)

M. A. Shifman, A. I. Vainshtein, and V. I. Zakharov, Nucl. Phys. B166, 493 (1980)