Accidental Peccei-Quinn Symmetry Protected to Arbitrary Order

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A $SU(N)_L \times SU(N)_R$ gauge theory for a scalar multiplet Y transforming in the bifundamental representation $(N, \bar{N})$ preserves, for $N > 4$, an accidental $U(1)$ symmetry first broken at operator dimension N. A vacuum expectation value for Y can break the symmetry to $\mathcal{H}_s = SU(N)_{L+R}$ or to $\mathcal{H}_h = SU(N-1)_L \times SU(N-1)_R \times U(1)_{L+R}$. In the first case the accidental $U(1)$ gets also broken, yielding a pseudo-Nambu-Goldstone boson with mass suppression controlled by N. 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 $U(1)_{\rm PQ}$ protected by the gauge symmetry up to order N. 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 only weak isospin, models for axionlike particles can be constructed.

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Introduction.—QCD has been firmly established as the correct description of strong interaction phenomena in particle physics. However, 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_s, determines the strength of the QCD interactions. Its measured value is a natural one for a dimensionless quantity (roughly speaking it is of order unity). 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 $[0, 2\pi)$, and also in this case it would be natural to expect $\theta \sim O(1)$. Instead, experimental limits on the neutron electric dipole moment yield the upper bound $|\theta| < 10^{-10}$, 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, $\theta = 0$. An elegant mechanism to guarantee the vanishing of θ was proposed in 1977 by Peccei and Quinn (PQ) [1,2]. It relies on a $U(1)_{\rm PQ}$ global symmetry, anomalous with respect to QCD, spontaneously broken by the vacuum expectation value (VEV) of a standard model (SM) singlet scalar field at a scale $v_a \gg 100$ GeV, and broken explicitly by nonperturbative QCD effects at a scale $\Lambda_{\rm QCD} \sim 100$ 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 $U(1)_{\rm PQ}$ is not exactly massless: it is a pseudo-NGB, commonly referred to as the axion [3,4]. To account for $|\theta| < 10^{-10}$ any other source of explicit $U(1)_{\rm PQ}$ breaking besides QCD must either be absent or adequately suppressed. This is difficult to achieve, especially considering that $U(1)_{\rm PQ}$, being anomalous, is not even a real symmetry. Effective operators not respecting $U(1)_{\rm PQ}$ are then expected to arise and, even if suppressed by the Planck scale $m_P = 1.2 \times 10^{19}$ GeV, unless their dimension is larger than $d \sim 10$ would unavoidably give $|\theta| > 10^{-10}$ [5–8].

In this Letter we propose a mechanism that, on the basis of first principles, can protect $U(1)_{\rm PQ}$ to arbitrary accuracy. A scalar multiplet Y is assigned to the bifundamental representation $(N, \bar{N})$ of the gauge group $G_{LR}^{(N)} = SU(N)_L \times SU(N)_R$. For $N > 4$ an accidental $U(1)$ "rephasing" symmetry is enforced at the classical level, and it only gets broken at $d = N$ by $\det(Y)$. When the scalar gauge theory is coupled to fermions carrying color, $U(1)$ acquires a QCD anomaly. SB of $G_{LR}^{(N)}$ 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. 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 $G_{LR}^{(N)}$ scalar gauge theory.—For $N > 4$, $G_{LR}^{(N)}$ gauge invariance restricts the renormalizable potential for Y to the simple form

$$V_0 = \lambda[T - v_a^2/2]^2 + \lambda_A A, \tag{1}$$

where $T = {\rm Tr}[YY^\dagger]$ and $A = {\rm Tr}[{\rm Mnr}(YY^\dagger, 2)]$ with ${\rm Tr}[{\rm Mnr}(M, k)]$ denoting the trace of the matrix of the minors of order k of M [9]. We require $v_a^2 > 0$ to trigger SB and $|\lambda_A| < (2N/N-1)\lambda$ to ensure a potential bounded from below. The matrix Y^c of constant background values of $Y(x)$ can be written in its singular value decomposition as

$$\frac{\sqrt{2}}{v_a} Y^c = \mathcal{U}_L \hat{Y} \mathcal{U}_R^\dagger = U_L (\hat{\Phi}\, \hat{Y}) U_R^\dagger \to \hat{\Phi}\, \hat{Y}, \tag{2}$$

where $\hat{Y} = \mathrm{diag}(y_1, y_2, \ldots, y_N)$ is diagonal with real non-negative entries normalized such that $\sum_i y_i^2 = 1$, $\mathcal{U}_{L,R}$ are unitary matrices, $U_{L,R}$ are special unitary [$\det(U_{L,R}) = +1$], and $\hat{\Phi}$ is a diagonal matrix of phases such that $\log\,\det(\hat{\Phi}) = \log\,\det(\mathcal{U}_L \mathcal{U}_R^\dagger) = i\,\arg\,\det(Y^c) \equiv i\delta^c$. δ^c, being an angular variable, ranges in the interval $[0, 2\pi)$, and the physics must be invariant under the redefiniton $\delta^c \to \delta^c + 2n\pi$ ($n = 1, 2 \ldots$). The last (diagonal) form in Eq. (2) is obtained, without loss of generality, via a rigid $G_{LR}^{(N)}$ rotation.

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

$$\text{(i)}\; Y_s^c = \frac{v_a}{\sqrt{2}}\hat{\Phi}\hat{Y}_s, \qquad \hat{Y}_s = \frac{1}{\sqrt{N}}\mathrm{diag}(1, 1, \ldots, 1),$$
$$\text{(ii)}\; Y_h^c = \frac{v_a}{\sqrt{2}}\hat{\Phi}\hat{Y}_h, \qquad \hat{Y}_h = \mathrm{diag}(0, \ldots, 0, 1). \qquad (3)$$

The corresponding little groups are the two maximal subgroups of $G_{LR}^{(N)}$: $\mathcal{H}_s = SU(N)_{L+R}$ for Y_s^c and $\mathcal{H}_h = G_{LR}^{(N-1)} \times U(1)_{L+R}$ for Y_h^c, where $U(1)_{L+R}$ corresponds to the diagonal combination of the two generators of $G_{LR}^{(N)}$ proportional to $\lambda_{N^2-1}^{L,R} = \mathrm{diag}(1, 1, \ldots, 1-N)_{L,R}$. It is important to stress that $\mathcal{H}_{s,h}$ cannot get broken further by any type of perturbative effects [11] or, equivalently, that neither can the vanishing entries in Y_h^c be lifted nor can the equality of the entries in Y_s^c be spoiled. Some bibliographic remarks are in order: the minima of the potential for the case of global $G_{LR}^{(3)}$ (namely, the SM quark flavor symmetry) were studied in Ref. [10] [and with the assumption of a real det (Y) previously in Ref. [12]]. The possibility of raising perturbatively the vanishing entries in Y_h was addressed in Ref. [13], and it was found that minimization of the one-loop effective potential results in the same little groups $\mathcal{H}_{s,h}$. A more thorough breaking, yielding $y_{i\neq j} \neq y_j$, can be in fact obtained only by introducing additional reducible scalar representations [13,14].

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

$$\mathcal{P}(\xi) = \det(\xi I - YY^\dagger) = \sum_{n=0}^{N}(-1)^n C_n \xi^{N-n}, \qquad (4)$$

where I is the identity matrix, and $C_n = \mathrm{Tr}[\mathrm{Mnr}(YY^\dagger, n)]$ with $C_0 = 1$, $C_1 = T$, $C_2 = A$, $C_N = \det[YY^\dagger] \equiv |\mathcal{D}|^2$. The solutions of $\mathcal{P}(\xi) = 0$ are the eigenvalues of $YY^\dagger$ 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 $\mathcal{D} = \det Y$ is another invariant, since under $G_{LR}^{(N)}$, $\mathcal{D} \to \det(V_L Y V_R^\dagger) = \det Y$ [15]. However, while all C_n's respect the $U(1)$ accidental symmetry, under $Y \to e^{i(\alpha/N)}Y$, $\mathcal{D}(x) \to e^{i\alpha}\mathcal{D}(x)$. Thus, $U(1)$ gets first broken at $d = N$ by

$$V_D = \frac{k\mathcal{D} + k^*\mathcal{D}^*}{m_P^{N-4}} = \frac{2\kappa D}{m_P^{N-4}}\cos[\varphi + \delta(x)], \qquad (5)$$

where κ and φ are the modulus and argument of the coupling k, $D = |\mathcal{D}|$, $\delta(x) = \arg\mathcal{D}(x)$, and the m_P^{N-4} suppression stems from the assumption that V_D is generated by gravity effects. In case (i), the minimum of V_D is obtained for $\langle\delta(x)\rangle \equiv \delta^c = \pi - \varphi$, lowering the minimum of V_0 by the amount

$$\Delta V = v_a^4 \frac{2\kappa}{(2N)^{N/2}}\left(\frac{v_a}{m_P}\right)^{N-4}. \qquad (6)$$

Thus, in the breaking $G_{LR}^{(N)} \times U(1) \to \mathcal{H}_s$, of the initial $2(N^2-1)+1$ generators N^2-1 are left unbroken, N^2-1 are spontaneously broken, while, because of the explicit breaking V_D in Eq. (5), the NGB of the spontaneously broken global $U(1)$ acquires a tiny mass $O(\sqrt{\Delta V}/v_a)$. In case (ii) instead, a global $U(1)'$ generated by $\lambda_{N^2-1}^{L+R} + (N-1)I$ is preserved by Y_h^c, so that at the renormalizable level $G_{LR}^{(N)} \times U(1) \to \mathcal{H}_h \times U(1)'$. Although V_D breaks $U(1)'$ in interactions, since $\langle D\rangle = 0$ there is no SB, and no NGB arises.

Solutions to the strong CP problem.—Solutions to the strong CP problems can be implemented by introducing fermions carrying color. Let us proceed by first introducing four fermion multiplets transforming under $G_{LR}^{(N)}$ as $Q_L \sim (N, 1)$, $Q_R \sim (1, N)$, $\Psi_L \sim (\bar{N}, 1)$, $\Psi_R \sim (1, \bar{N})$. Since they can be combined into real representations of $SU(N)_{L,R}$ $(Q_{L,R} \oplus \Psi_{L,R})$, there are no gauge anomalies. Gauge symmetry allows for Yukawa couplings of the form $\bar{Q}_L Y Q_R + \bar{\Psi}_L Y^\dagger \Psi_R + \mathrm{H.c.}$, which preserve the Y rephasing symmetry if the fermions are transformed chirally with

$U(1)$ charges satisfying $\mathcal{X}_{Q_L} - \mathcal{X}_{Q_R} = \mathcal{X}_{\Psi_R} - \mathcal{X}_{\Psi_L} = \mathcal{X}_Y$. The absence of $U(1) - SU(N)_{L,R}$ mixed anomalies is ensured by the opposite sign of the two charge differences [16]. Let us now triplicate the fermion content, and assign $Q_{L,R}$ to the fundamental representation of color, while $\Psi^a_{L,R}$ $(a = 1, 2, 3)$ remain color singlets. Since there is no compensating cancellation of the $Q_{L,R}$ contribution, a $U(1)$-QCD anomaly arises.

(i) Solution with a fundamental axion.—We choose a basis in which the SM quark masses are real while $\theta_{\rm QCD} \neq 0$ and, without loss of generality, we take the $\Psi_{L,R}$ couplings flavor diagonal:

$$e^{i(\eta_0/N)} h_0 \bar{Q}_L Y Q_R + e^{i(\eta_a/N)} h_a \bar{\Psi}^a_L Y^\dagger \Psi^a_R, \qquad (7)$$

where h_0, h_a $(a = 1, 2, 3)$ are four real non-negative parameters. If $\lambda_A < 0$, the minimum Y^c_s is selected, and all the fermions become massive, with degenerate masses within each multiplet that, as we will now show, can be brought into real form without inducing mixed $G^{(N)}_{LR}$ anomalies. After SB, $\arg \det(M_Q) = \eta_0 + \delta^c$ and $\arg \det(M_a) = \eta_a - \delta^c$; 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 α_0, α_a, subject to a fifth condition $\sum_{a=1}^3 \alpha_a = 3\alpha_0$ to avoid mixed anomalies. The phase of Y can be also redefined [this changes the argument of the cosine in Eq. (5) as $\varphi \to \tilde{\varphi}$], and all the complex phases can thus be canceled. However, chiral rotations of $Q_{L,R}$ are anomalous with respect to $SU(3)_c$, and another source of explicit $U(1)$ breaking is then introduced. After including it, the relevant potential for $\delta(x)$ acquires the form

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

where $\tilde{\varphi}$ is unrelated with $\theta_{\rm QCD}$, and we have redefined $\delta(x) + \theta_{\rm QCD} \to \delta(x)$ so that the anomalous coupling to the gluons reads $(\alpha_s/4\pi)\delta G\tilde{G}$. Note that when the angular variable $\delta(x)$ is varied in the interval $[0, 2\pi)$, there is a unique minimum of the potential so that, independently of N, the number of domain walls [17] is always 1. From Eq. (8) we see that if $\kappa, \tilde{\varphi} \sim O(1)$, as it is natural to assume, $\delta^c < 10^{-10}$ can be ensured only if the explicit breaking satisfies $\Delta V/(f_\pi^2 m_\pi^2) \lesssim 10^{-10}$. For the phenomenologically preferred interval $10^9 \lesssim v_a/N \lesssim 10^{12}$ GeV this can be fulfilled for $9 \leq N \leq 13$ [18].

Let us now proceed to identify the axion field. In the unitary gauge in which the rigid $G^{(N)}_{LR}$ rotation yielding $Y^c \sim \hat{\Phi}\hat{Y}$ in Eq. (2) is promoted to local, we can write $Y(x) = \hat{\Phi}(x)\hat{Y}(x)$ with

$$\hat{\Phi}(x) = {\rm diag}(e^{i\hat{\gamma}_1(x)}, \ldots, e^{i\hat{\gamma}_N(x)}); \quad \hat{\gamma}_i = \frac{\sqrt{2N}}{v_a}\gamma_i. \qquad (9)$$

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

$$a_a(x) = 2{\rm Tr}[\vec{\gamma}(x) \cdot T_a], \qquad (10)$$

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

$$a_0(x) = 2\,{\rm Tr}[\vec{\gamma} \cdot T_0] = \frac{v_a}{N}\delta(x). \qquad (11)$$

Note that since the periodicity of $\delta(x)$ is 2π, the periodicity of the axion is $a_0 \to a_0 + (2\pi/N)v_a$. Then, one might wonder whether there are N domain walls corresponding to the N minima $\langle a_0 \rangle + (2\pi n/N)v_a$, $(n = 0, \ldots, N - 1)$. This is not so because all these minima are gauge equivalent, in the sense that the Z_N center of $SU(N)_{L+R}$ has precisely as elements $\exp(i2\pi n/N)I$, so that the cyclic values of a_0/v_a are all connected by gauge transformations. Neglecting the subdominant gravitational contributions, the axion mass is $m_a = N(m_\pi f_\pi)/v_a$, while the strength of the axion-photon coupling $(1/4)g_{a\gamma\gamma}F_{\mu\nu}\tilde{F}^{\mu\nu}$ is

$$g_{a\gamma\gamma} = -1.92\frac{m_a}{\rm eV}\frac{2.0}{10^{10}\ {\rm GeV}}, \qquad (12)$$

which falls within the axion window in Fig. 1.

(ii) Solution with a composite axion.— For $\lambda_A > 0$ the VEV Y^c_h 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 them remains

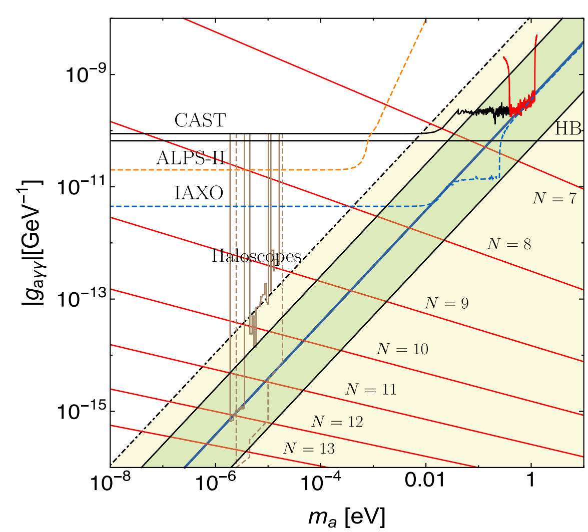

FIG. 1. ALP-photon coupling versus mass for different gauge groups $G^{(N)}_{LR}$. The green band and the yellow area represent two preferred regions for axion models [19,20]. The blue line depicts the axion coupling in Eq. (12).

unbroken. While this hints to the possibility of a massless-quark type of solution where θ is simply removed via a chiral $U(1)'$ rotation, such a scenario is not viable. This is because, although the massless Q's and Ψ's would get confined into F hadrons once, at $\Lambda_F \gg \Lambda_{\rm QCD}$, $G_{LR}^{(N-1)}$ enters the strongly coupled regime, matching the $U(1)'$-QCD anomaly in the low energy theory [21] requires that some composite fermions carrying color remain massless, and these are not observed. We are then led to assume SB of $U(1)'$ by some color neutral condensate of Q and Ψ. Then, the pseudo-NGB of the $U(1)'$ symmetry would correspond to a composite axion [22–24] with mass and couplings suppressed as $1/\Lambda_F$. Clearly, this second scenario is more speculative, so we will mainly focus on the first one.

Phenomenology.—In the first scenario, after SB the spectrum consists of (i). N^2-1 gauge bosons with masses $O(v_a)$, (ii). N quarks Q and $3N$ SM singlet fermions Ψ^a, stable at the tree level, with $m_{Q,\Psi} \sim O(v_a/\sqrt{N})$, (iii). N^2-1 massless gauge bosons F, and (iv). one pseudo-NGB (the axion). Tree level stability of the heavy Q's and Ψ's follows from the fact that they do not carry weak isospin and hypercharge $\mathcal{Y}$, and thus cannot decay into SM fermions. However, cosmologically stable heavy relics, and in particular long-lived strongly interacting particles, represent serious issues in cosmology and astrophysics (see Refs. [19,20] for a recent discussion and relevant references). A simple way to avoid all phenomenological problems is to assume a preinflationary scenario (PQ symmetry broken before inflation) and $v_a > T_{\rm reheating}$ so that, after inflation has wiped away all the heavy states, they cannot be regenerated. After inflation, the massless gauge bosons F could be produced via gravitational effects, or via gluon-gluon to FF scatterings. The first mode is suppressed by powers of m_P and typically very inefficient. The rate for the second process can be estimated as $(\alpha_F\alpha_s)^2T^9/m_Q^8$ with α_F the coupling strength of the new gauge group. This reaction remains well out of equilibrium for all $T < m_Q$ and is also too inefficient to produce F in sizable amounts. We can then conclude that preinflationary scenarios do not leave dangerous heavy relics in appreciably amounts.

Postinflationary scenarios (PQ symmetry broken after inflation and $v_a < T_{\rm reheating}$) yield a different picture. During reheating, all the heavy states attain equilibrium distributions. At $T \lesssim v_a$ the massive gauge bosons readily decay into the lighter Q, Ψ fermions. Below m_Q, the unbroken $SU(N)_{L+R}$ corresponds to a pure Yang-Mills theory with large N that rapidly flows towards a confining regime at $\Lambda_F \gg \Lambda_{\rm QCD}$. Bound state mesons $\Pi_{Q(\Psi)} \sim Q\bar{Q}(\Psi\bar{\Psi})$ singlets under $SU(N)_{L+R}$ form, and readily decay into lighter "gaugeballs" $G \sim FF$ of mass $m_G \sim O(\Lambda_F)$. Gaugeballs can decay invisibly into two gravitons [25] or visibly into a pair of gluons with a rate $\Gamma_{gg} \sim \Lambda_F^9/m_Q^8$. Visible decays are generally dominant, and lifetimes sufficiently short ($\tau_G \lesssim 10^{-2}$ s) to evade constraints from big bang nucleosynthesis are ensured if $m_Q/\Lambda_F \lesssim 3\times 10^3$. However, there are other more dangerous relics, like $\mathcal{M}_{ab} \sim \Psi_a\bar{\Psi}_b$ with $(a \neq b)$ and, since at Λ_F color is unconfined, "mongrel" mesons $\mathcal{M}_a \sim Q\Psi_a$ will also form. $\mathcal{M}_{ab}$ decays are forbidden by Ψ-flavor conservation, and $\mathcal{M}_a$ decays are forbidden also by color conservation. The abundance of these states is basically determined by free particle annihilation before Λ_F confinement, which always results in $\Omega_{\mathcal{M}} \gg \Omega_{\rm DM}$ unless the relevant mass scale is not much larger than a few TeVs. This requires tiny values for the Yukawa couplings in Eq. (7) and an appropriately small initial value of α_F to ensure that also $\Lambda_F \lesssim$ a few TeVs. All in all, the postinflationary scenario, if not ruled out, is certainly strongly disfavored with respect to preinflationary scenarios.

Axionlike particles.—With no attempt to solve the strong CP problem, models for axionlike particles (ALPs) can be constructed along these same lines. Instead of new quarks, let us introduce new colorless fermions $\mathcal{T}$, doublets under weak isospin and with zero hypercharge. SM singlet fermions $\Psi_{L,R}^{1,2}$ are again needed to cancel gauge anomalies. Now, in the breaking $G_{LR}^{(N)} \to \mathcal{H}_s$ the NGB of the accidental $U(1)$ only receives mass from the 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 couplings to the photon are possible. We obtain

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

$$g_{a\gamma\gamma} = \frac{\alpha}{2\pi}\frac{E}{v_a} = \frac{\alpha}{4\pi}\frac{N}{v_a}, \tag{14}$$

where in the second equation we have used the electromagnetic anomaly coefficient $E = 2N\mathcal{X}_{\mathcal{T}_L}\mathcal{Q}_\pm^2 = N/2$ with $\mathcal{Q}_\pm = \pm\frac{1}{2}$ the electric charges of the components of $\mathcal{T}_L$, and $\mathcal{X}_{\mathcal{T}_R} = 0$, $\mathcal{X}_{\mathcal{T}_L} = \mathcal{X}_Y = 1$ the (assigned) PQ charges. Equations (13) and (14) yield

$$g_{a\gamma\gamma} = \frac{\sqrt{2N}\alpha}{8\pi m_P}\left(\frac{\kappa N m_P^2}{m_a^2}\right)^{\frac{1}{N-2}}. \tag{15}$$

Figure 1 gives the ALP-photon coupling versus m_a for different values of N, together with the preferred regions for axion models [19,20].

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 Kim-Shifman-Vainshtein-Zakharov type of axion models [26,27] since PQ charges are carried only by non-SM particles. A new gauge group $G_{LR}^{(N)} = SU(N)_L \times SU(N)_R$ is postulated, and new quarks $Q_{L,R}$ are assigned to fundamental representations of $SU(N)_{L,R}$ while the PQ scalar, rather than being a single complex field, is a matrix Y transforming in the bifundamental 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. Depending on the gauge symmetry breaking pattern, a different possibility where the axion is composite can be also realized.

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