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The physical origin of the strong CP problem in QCD, rooted in the structures of the vacuum of the standard model, is reviewed. The chiral solution to this problem, with its accompanying axion, is explained and various characteristics of axions are detailed. Although visible axion models are excluded experimentally, models with very light and very weakly coupled axions are still tenable. The astrophysical and cosmological implications of such axion models are discussed, along with ongoing experimental attempts to detect such, so called, invisible axions.

THE $U(1)_A$ PROBLEM

Quantum Chromodynamics (QCD) describes the strong interactions of hadrons in terms of the interactions of their quark and gluon constituents. The QCD Lagrangian

$$\mathcal{L}_{\text{QCD}} = - \sum_f \bar{q}_f \left(\gamma^\mu \frac{1}{i} D_\mu + m_f \right) q_f - \frac{1}{4} G_a^{\mu\nu} G_{a\mu\nu} \quad (1)$$

for f flavors of quarks has a large global symmetry in the limit when $m_f \rightarrow 0$: $G = U(f)_R \times U(f)_L$. This symmetry corresponds to the freedom of arbitrary chiral rotations of the f flavor of quarks into each other. Because $m_u, m_d \ll \Lambda_{\text{QCD}}$ —with Λ_{QCD} being the dynamical scale of the theory—in practice only a chiral $U(2)_R \times U(2)_L$ symmetry is actually a very good approximate global symmetry of the strong interactions. This symmetry, however, is not manifest in the spectrum of hadrons. It is instructive to understand why this is so.

The $U(1)_V$ subgroup of the $U(2)_R \times U(2)_L$ symmetry, corresponding to vectorial ($V = R+L$) baryon number, is an exact symmetry of QCD, irrespective of the value of the quark masses. In the limit of $m_u = m_d$ a further $SU(2)_V$ subgroup would be an exact (isospin) symmetry of QCD. Because the u - and d -quark masses are so light compared to Λ_{QCD} , even with $m_u \neq m_d$ one has an approximately degenerate multiplet of hadrons corresponding to this symmetry (e.g. the π -triplet and the doublet of neutrons and protons). However, the corresponding axial symmetries in $U(2)_R \times U(2)_L$ ($A = R - L$) are not seen in the spectrum. For instance, there are no parity doublets degenerate with the neutron and proton. This phenomena is understood because the $SU(2)_A$ symmetry is not preserved by the QCD vacuum. Indeed, as is appropriate for a spontaneously broken global symmetry, there is instead in the QCD hadron spectrum an approximate triplet of Nambu-Goldstone bosons—the pions—whose mass vanishes in the limit that $m_u, m_d \rightarrow 0$. Remarkably, however, there is no corresponding pseudoscalar state with vanishing mass, in the same limit, corresponding to the Nambu-Goldstone boson of a $U(1)_A$ symmetry.

The nature of the approximate $SU(2)_V \times SU(2)_A \times U(1)_V$ symmetry of the strong interactions was actually understood before the advent of QCD [1]. Although the anomalous features of the $U(1)_A$ symmetry were also pinpointed around the time of the development of QCD as the theory of the strong interactions [2], the resolution of the $U(1)_A$ problem required a better understanding of the QCD vacuum. In short, the reason why there are no approximate Nambu-Goldstone bosons associated with this Abelian chiral symmetry is that, as a result of chiral anomaly [3], this is really **not** a quantum symmetry of QCD. In turn, however, this more complete understanding opens up another problem—the strong CP problem [4]. Namely, the richer vacuum of QCD, combined with the violation of CP in the weak interactions, allows for the presence of an effective interaction

$$\mathcal{L}_{\text{strong CP}} = \bar{\theta} \frac{\alpha_s}{8\pi} G_{a\mu\nu} \tilde{G}_a^{\mu\nu} \quad (2)$$

which leads to an enormous neutron electric dipole moment [$d_n \sim e\bar{\theta}(m_q/M_N^2)$], unless the parameter $\bar{\theta}$ is tiny ($\bar{\theta} \leq 10^{-9}$).

In QCD, the dynamical formation of quark condensates, $\langle \bar{u}u \rangle = \langle \bar{d}d \rangle \neq 0$, breaks the $U(2)_A$ global symmetry. As a result, in the limit of vanishing quark masses, the pion triplet are the Nambu Goldstone bosons associated with the spontaneously broken $SU(2)_A$ current $J_{5i}^\mu = \bar{q} \gamma^\mu \gamma_5 \frac{\tau_i}{2} q$, where q is the (u, d) doublet of quark fields. The existence of these massless states produces a $q^2 = 0$ pole in the Green's functions of the $SU(2)_A$ currents with pseudoscalar quark densities. One has

$$\begin{aligned}
M_{ij}^\mu(q) &= \int d^4x e^{-iqx} \langle T(J_{5i}^\mu(x), \bar{q}\gamma_5\tau_j q) \rangle \\
&\sim \frac{\langle \bar{q}q \rangle \delta_{ij} q^\mu}{q^2} + \text{non singular terms} .
\end{aligned} \tag{3}$$

Because the $\langle \bar{q}q \rangle$ condensates also break the $U(1)_A$ global symmetry, naively one expects also a similar q^2 singularity associated with the $U(1)_A$ current, connected to a further Nambu-Goldstone boson—the η in this 2-flavor discussion. However, such an approximate Nambu-Goldstone excitation cannot really exist, since if it did, it would lead (in the limit of finite but small quark masses) to a state essentially degenerate with the pions [2]—in contradiction with experiment.

The situation is a bit more complicated because the $U(1)_A$ current $J_5^\mu = \bar{q}\gamma^\mu\gamma_5\frac{1}{2}q$ has a chiral anomaly [3]

$$\partial_\mu J_5^\mu = \frac{\alpha_s}{4\pi} G_a^{\mu\nu} \tilde{G}_{a\mu\nu} . \tag{4}$$

However, since the RHS of the above equation is a total divergence [5]

$$\begin{aligned}
G_a^{\mu\nu} \tilde{G}_{a\mu\nu} &= \partial_\mu K^\mu ; \\
K^\mu &= \epsilon^{\mu\alpha\beta\gamma} A_{a\alpha} \left[G_{a\beta\gamma} - \frac{g_s}{3} f_{abc} A_{b\beta} A_{c\gamma} \right] ,
\end{aligned} \tag{5}$$

one can construct a conserved chiral current

$$\tilde{J}_5^\mu = J_5^\mu - \frac{\alpha_s}{4\pi} K^\mu . \tag{6}$$

As a result, in the massless quark limit, if $\langle \bar{q}q \rangle \neq 0$, one can show that the Green's functions of the \tilde{J}_5^μ current have a $q^2 = 0$ singularity analogous to that of the chiral $SU(2)_A$ currents:

$$\begin{aligned}
\tilde{M}^\mu(q) &= \int d^4x e^{-iqx} \langle T(\tilde{J}_5^\mu(x), \bar{q}\gamma_5 q) \rangle \\
&\sim \frac{\langle \bar{q}q \rangle q^\mu}{q^2} + \text{non-singular terms} .
\end{aligned} \tag{7}$$

The $U(1)_A$ problem, then, is why such a singularity is not manifested physically in the spectrum of hadrons?

THE QCD VACUUM

The resolution of the $U(1)_A$ problem came from a better understanding of the structure of the QCD vacuum. For gauge fields one can take the vacuum state to be the state where the vector potential is either zero or is in a gauge configuration equivalent to zero. A proper study of these configurations reveals that the structure of the vacuum state for non-Abelian gauge theories (like QCD) is richer than expected. Following Callan, Dashen, and Gross [6] it is convenient to study these theories in the temporal gauge $A_a^0 = 0$ and concentrate on an $SU(2)$ subgroup. With this gauge choice, the spatial gauge fields are time independent and under a gauge transformation transform as

$$\frac{\tau_a}{2} A_a^i(\vec{r}) \equiv A^i(\vec{r}) \rightarrow \Omega(\vec{r}) A^i(\vec{r}) \Omega(\vec{r})^{-1} + \frac{i}{g_s} \Omega(\vec{r}) \nabla^i \Omega^{-1}(\vec{r}) . \tag{8}$$

Thus, in this gauge, vacuum configurations corresponds to a vanishing vector potential, or to $\frac{i}{g_s} \Omega(\vec{r}) \nabla^i \Omega^{-1}(\vec{r})$.

The rich vacuum structure of QCD results from the requirement that the gauge transformation matrices $\Omega(\vec{r})$ go to unity at spatial infinity. Such a requirement maps the physical space onto the group space and this $S_3 \rightarrow S_3$ map splits $\Omega(\vec{r})$ into different homotopy classes $\{\Omega_n(\vec{r})\}$, characterized by an integer winding number n detailing how precisely $\Omega(\vec{r})$ behaves as $\vec{r} \rightarrow \infty$:

$$\Omega_n(\vec{r}) \xrightarrow{\vec{r} \rightarrow \infty} e^{2\pi i n} . \tag{9}$$

The winding number n is related to the Jacobian of the $S_3 \rightarrow S_3$ transformation

$$n = \frac{ig_s^3}{24\pi^2} \int d^3r \text{Tr} \epsilon_{ijk} A_n^i(\vec{r}) A_n^j(\vec{r}) A_n^k(\vec{r}) , \tag{10}$$

with $A_n^i(\vec{r})$ being the pure gauge field corresponding to the gauge transformation matrix $\Omega_n(\vec{r})$.

Because one can construct the n -gauge transformation matrix $\Omega_n(\vec{r})$ by compounding n times $\Omega_1(\vec{r})$, it follows that the vacuum state corresponding to $A_n(\vec{r})$ is not really gauge invariant. Indeed, the action of the gauge transformation matrix Ω_1 on an n -vacuum state gives an $(n + 1)$ -vacuum state

$$\Omega_1|n\rangle = |n + 1\rangle . \quad (11)$$

Nevertheless, one can construct a gauge-invariant vacuum state—the, so called, θ -vacuum—by superposing these n -vacua [6]

$$|\theta\rangle = \sum_n e^{-in\theta}|n\rangle . \quad (12)$$

Indeed, it is easy to check that $\Omega_1|\theta\rangle = e^{i\theta}|\theta\rangle$.

Because of this more complex vacuum structure, the vacuum functional for the theory splits into distinct sectors.

$${}_+\langle\theta|\theta\rangle_- = \sum_{m,n} e^{im\theta} e^{-in\theta} {}_+\langle m|n\rangle_- \quad (13)$$

$$= \sum_\nu e^{i\nu\theta} \left[\sum_n {}_+\langle n + \nu|n\rangle_- \right] . \quad (14)$$

This superposition of transition amplitudes from $t = -\infty$ to $t = +\infty$ with fixed n -vacuum difference ν , weighted by different phases $e^{i\nu\theta}$ introduces an additional intrinsic source of CP violation. By using Eqs. (5) and (10), one can show [4] that the integer ν has a gauge invariant meaning

$$\nu = \frac{\alpha_s}{8\pi} \int d^4x G_{a\mu\nu} \tilde{G}_a^{\mu\nu} . \quad (15)$$

Thus, writing for the vacuum functional the usual path integral formula

$${}_+\langle\theta|\theta\rangle_- = \sum_\nu \int \delta A^\mu e^{iS_{\text{eff}}[A]} \delta \left[\nu - \frac{\alpha_s}{8\pi} \int d^4x G_a^{\mu\nu} \tilde{G}_{a\mu\nu} \right] , \quad (16)$$

one can re-interpret the θ term in terms of an effective action

$$S_{\text{eff}} = S[A] + \theta \frac{\alpha_s}{8\pi} \int d^4x G_a^{\mu\nu} \tilde{G}_{a\mu\nu} . \quad (17)$$

This additional term violates P and CP, since it corresponds to an $\vec{E}_a \cdot \vec{B}_a$ interaction of the color fields.

Perturbation theory is connected to the $\nu = 0$ sectors where the $G\tilde{G}$ term vanishes. The effects of the $\nu \neq 0$ sectors thus necessarily are nonperturbative. These contributions are naturally selected by the connection of $G\tilde{G}$ to the chiral anomaly. For n_f flavors, the axial current J_5^μ has an anomaly

$$\partial_\mu J_5^\mu = n_f \frac{\alpha_s}{8\pi} G_{a\mu\nu} \tilde{G}_a^{\mu\nu} \quad (18)$$

Thus a chirality change ΔQ_5 is related to ν as:

$$\Delta Q_5 = \int d^4x \partial_\mu J_5^\mu = n_f \frac{\alpha_s}{8\pi} \int d^4x G_a^{\mu\nu} \tilde{G}_{a\mu\nu} = n_f \nu \quad (19)$$

The solution to the $U(1)_A$ problem is connected to the chirality breakdown in QCD in the $\nu \neq 0$ sectors. If these sectors are included—as they should—then $U(1)_A$ is **never** a symmetry.

For $n_f = 2$, the conserved current discussed earlier \tilde{J}_5^μ is not gauge invariant for transformations with non-trivial winding number. One can show [4] that the associated chiral charge

$$\tilde{Q}_5 = \int d^3x \tilde{J}_5^0 \quad (20)$$

transforms as

$$\Omega_1 \tilde{Q}_5 \Omega_1^{-1} = \tilde{Q}_5 + 2 \quad (21)$$

when one worries about $\nu \neq 0$ sectors. 't Hooft showed that it is this lack of gauge invariance which is crucial to resolve the $U(1)_A$ problem [7]. Although Green's functions of the \tilde{J}_5^μ currents contain $1/q^2$ singularities, these singularities are cancelled for the corresponding Green's functions of the gauge invariant currents J_5^μ . The Green's function for this current with a pseudoscalar density obeys the equation

$$M^\mu(q) = \tilde{M}^\mu(q) + \frac{\alpha_s}{4\pi} \int d^4x e^{-iqx} \langle T(K^\mu(x), \bar{q}\gamma_5 q) \rangle \quad (22)$$

Although the first term above is singular at $q^2 = 0$, when one considers the effects of $\nu \neq 0$ sectors, so is the second term. These singularities **cancel** for the Green's function of the physical current J_5^μ [7]. Because of Eq. (19), the $U(1)_A$ symmetry is really not a symmetry at all. Therefore, not surprisingly, there are really no physical Nambu Goldstone bosons associated with this pseudo symmetry!

THE STRONG CP PROBLEM

The effective interaction (17) is not the only new source of CP violation arising from the more complex structure of the QCD vacuum. It is augmented by an analogous term coming from the electroweak sector of the theory. The origin of this additional interaction is as follows. In general, the mass matrix of quarks which emerges from the spontaneous breakdown of the electroweak gauge symmetry is neither Hermitian nor diagonal:

$$\mathcal{L}_{\text{mass}} = -\bar{q}_{R_i} M_{ij} q_{L_j} - \bar{q}_{L_i} (M^\dagger)_{ij} q_{R_j} . \quad (23)$$

This matrix can be diagonalized by separate unitary transformations of the chiral quark fields. These transformations encompass a chiral $U(1)_A$ transformation

$$q_R \rightarrow e^{i\alpha/2} q_R ; \quad q_L \rightarrow e^{-i\alpha/2} q_L \quad (24)$$

with $\alpha = \frac{1}{n_f} \text{Arg det } M$. Such chiral transformations, in effect, change the vacuum angle [8]. That this is so follows from Eq. (21). Using this equation (replacing the factor of 2 by the number of flavors n_f), it is easy to see that a chiral $U(1)_A$ rotation on the θ vacuum shifts the vacuum angle by $\text{Arg det } M$ [4]:

$$e^{i\alpha\tilde{Q}_5} |\theta\rangle = |\theta + n_f\alpha\rangle = |\theta + \text{Arg det } M\rangle . \quad (25)$$

Hence, in the full theory, the effective CP-violating interaction ensuing from the more complex structure of the QCD vacuum is that given in Eq. (2), with

$$\bar{\theta} = \theta + \text{Arg det } M . \quad (26)$$

The strong CP problem is really why the combination of QCD and electroweak parameters which make up $\bar{\theta}$ should be so small. As we remarked upon earlier, the effective interaction (2) gives a direct contribution to the electric dipole moment of the neutron and the strong experimental bound on d_n requires that $\bar{\theta} \leq 10^{-9}$ [9]. In principle, because $\bar{\theta}$ is a free parameter in the theory, any value of $\bar{\theta}$ is equally likely. However, one would like to obtain an understanding from some underlying physics of why this number is so small.

THE CHIRAL SOLUTION TO THE STRONG CP PROBLEM

Almost twenty years ago, Helen Quinn and I [10] suggested a dynamical solution to the strong CP problem. What we postulated was that the **full** Lagrangian of the standard model was invariant under an additional global chiral $U(1)$ symmetry. If this $U(1)_{PQ}$ symmetry existed, and it were exact, the strong CP problem would be trivially solved since θ could be set to zero through such a chiral transformation (cf. Eq. (25)). However, physically such a symmetry cannot be exact. Nevertheless, what Quinn and I [10] showed was that, even if $U(1)_{PQ}$ is spontaneously broken, the parameter $\bar{\theta}$ is dynamically driven to zero. However, in this case there is an associated pseudo Nambu-Goldstone boson in the theory, the axion [11].¹

¹The axion is not massless because the chiral $U(1)_{PQ}$ symmetry is anomalous. As a result, the axion gets a mass of order $m_a \sim \Lambda_{\text{QCD}}^2/f$, where f is the scale associated with the breakdown of the $U(1)_{PQ}$ symmetry.

It is useful to understand schematically how a global $U(1)_{\text{PQ}}$ symmetry solves the strong CP problem [4]. Basically, what happens is that by incorporating this symmetry in the theory one replaces the static CP violating parameter $\bar{\theta}$ by the dynamical (CP conserving) interactions of the axion field $a(x)$. Because the axion field is the Nambu-Goldstone boson associated with the spontaneously broken $U(1)_{\text{PQ}}$ symmetry, this field translates under a $U(1)_{\text{PQ}}$ transformation. If α is the phase parameter of this transformation and f is the scale associated with the breakdown of the symmetry, one has

$$a(x) \xrightarrow{U(1)_{\text{PQ}}} a(x) + \alpha f . \quad (27)$$

It follows, therefore, that if the effective Lagrangian describing the full theory is to be $U(1)_{\text{PQ}}$ invariant, the axion field must only enter derivatively coupled. This statement would be exactly true if the $U(1)_{\text{PQ}}$ symmetry were not anomalous. Because of the chiral anomaly, however, \mathcal{L}_{eff} must also have a term in which the axion field couples directly to the gluon density $G\tilde{G}$, so as to guarantee that J_{PQ}^μ has the correct QCD chiral anomaly.

The above considerations fix the form of the effective standard model Lagrangian, if it is augmented by an extra $U(1)_{\text{PQ}}$ global symmetry as we suggested [10]. Focusing only on the extra terms involving the axion field one has

$$\begin{aligned} \mathcal{L}_{\text{SM}}^{\text{eff}} = & \mathcal{L}_{\text{SM}} + \bar{\theta} \frac{\alpha_s}{8\pi} G_a^{\mu\nu} \tilde{G}_{a\mu\nu} - \frac{1}{2} \partial_\mu a \partial^\mu a \\ & + \mathcal{L}_{\text{int.}}[\partial^\mu a/f; \psi] + \frac{a}{f} \xi \frac{\alpha_s}{8\pi} G_a^{\mu\nu} \tilde{G}_{a\mu\nu} . \end{aligned} \quad (28)$$

In the above ψ stands for any field in the theory and ξ is a model-dependent parameter associated with the chiral anomaly of the $U(1)_{\text{PQ}}$ current

$$\partial_\mu J_{\text{PQ}}^\mu = \xi \frac{\alpha_s}{8\pi} G_a^{\mu\nu} \tilde{G}_{a\mu\nu} . \quad (29)$$

The presence of the last term in (28) provides an effective potential for the axion field. Thus its vacuum expectation is no longer arbitrary. Indeed the minimum of this potential determines the axion field VEV $\langle a \rangle$. One has

$$\left\langle \frac{\partial V_{\text{eff}}}{\partial a} \right\rangle = -\frac{\xi}{f} \frac{\alpha_s}{8\pi} \langle G_a^{\mu\nu} \tilde{G}_{a\mu\nu} \rangle \Big|_{\langle a \rangle} = 0 . \quad (30)$$

What Quinn and I showed [10] is that the periodicity of the pseudoscalar density expectation value $\langle G\tilde{G} \rangle$ in the relevant θ -parameter, $\bar{\theta} + \frac{\langle a \rangle}{f} \xi$, forces the axion VEV to take the value

$$\langle a \rangle = -\frac{f}{\xi} \bar{\theta} . \quad (31)$$

This solves the strong CP problem, since $\mathcal{L}_{\text{SM}}^{\text{eff}}$, when expressed in terms of the physical axion field, $a_{\text{phys}} = a - \langle a \rangle$, no longer contains the CP violating $\bar{\theta}G\tilde{G}$ term. Furthermore, expanding V_{eff} at its minimum, one sees that the axion itself gets a mass:

$$m_a^2 = \left\langle \frac{\partial^2 V_{\text{eff}}}{\partial a^2} \right\rangle = -\frac{\xi}{f} \frac{\alpha_s}{8\pi} \frac{\partial}{\partial a} \langle G_a^{\mu\nu} \tilde{G}_{a\mu\nu} \rangle \Big|_{\langle a \rangle} . \quad (32)$$

Therefore, the standard model with an additional $U(1)_{\text{PQ}}$ symmetry no longer has a dangerous CP violating interaction. Instead, it contains additional interactions of a massive axion field both with matter and gluons characterized by a scale f :

$$\begin{aligned} \mathcal{L}_{\text{SM}}^{\text{eff}} = & \mathcal{L}_{\text{SM}} + \mathcal{L}_{\text{int.}}[\partial^\mu a_{\text{phys.}}/f; \psi] - \frac{1}{2} \partial^\mu a_{\text{phys.}} \partial_\mu a_{\text{phys.}} \\ & - \frac{1}{2} m_a^2 a_{\text{phys.}}^2 + \frac{a_{\text{phys.}}}{f} \xi \frac{\alpha_s}{8\pi} G_a^{\mu\nu} \tilde{G}_{a\mu\nu} . \end{aligned} \quad (33)$$

In practical applications, it is more convenient to use the freedom of $U(1)_{\text{PQ}}$ transformations to replace the last term in Eq. (33) in favor of effective interactions of axions with the light pseudoscalar mesons (π and η) of QCD [12]. Although the properties of axions are model-dependent, this dependence can be isolated in a few numerical coefficients related to the axion mass, its coupling to two photons and the mixing of axions with the neutral pion and the eta. In terms of the pion decay constant $f_\pi \simeq 92$ MeV and the $SU(2) \times U(1)$ VEV $v \simeq 250$ GeV, it proves convenient to define a “standard” axion mass parameter

$$m_a^{\text{st}} = \frac{m_\pi f_\pi}{v} \frac{\sqrt{m_u m_d}}{(m_u + m_d)} \simeq 25 \text{ KeV} . \quad (34)$$

Then, for all models, one can characterize the mass of the axion and its π and η couplings by

$$m_a = \lambda_m m_a^{\text{st}} \left(\frac{v}{f} \right) \quad (35)$$

and

$$\xi_{a\pi} = \lambda_3 \frac{f_\pi}{f} ; \quad \xi_{a\eta} = \lambda_o \frac{f_\pi}{f} , \quad (36)$$

with λ_m , λ_3 , and λ_o model parameters of $O(1)$.² Similarly, one can characterize the effective couplings of axions to two photons by the interaction

$$\mathcal{L}_{a\gamma\gamma} = \frac{\alpha}{4\pi} K_{a\gamma\gamma} \frac{a_{\text{phys.}}}{f} F^{\mu\nu} \tilde{F}_{\mu\nu} , \quad (37)$$

where $K_{a\gamma\gamma}$ is again a model-dependent parameter of $O(1)$. It is clear that if $f_\pi \ll f$, the axion is both very light and very weakly coupled!

It is informative to sketch the derivation of the above formulas in the original Peccei-Quinn model [10]. Here the $U(1)_{\text{PQ}}$ symmetry is introduced in the standard model by having two Higgs doublets:

$$\mathcal{L}_{\text{Yukawa}} = \Gamma_{ij}^u \bar{Q}_{L_i} \Phi_1 u_{R_j} + \Gamma_{ij}^d \bar{Q}_{L_i} \Phi_2 d_{R_j} + \text{h.c.} \quad (38)$$

The presence of Φ_1 and Φ_2 allows $\mathcal{L}_{\text{Yukawa}}$ to be invariant under independent rotations of u_{R} and d_{R} —the desired chiral $U(1)_{\text{PQ}}$ transformation. The axion in this model is the common phase field of Φ_1 and Φ_2 orthogonal to weak hypercharge. If $x = v_2/v_1$ is the ratio of the Higgs VEV and $f = \sqrt{v_1^2 + v_2^2} = (\sqrt{2}G_F) = v \simeq 250$ GeV, then it is easy to isolate the axion content in Φ_1 and Φ_2 as

$$\Phi_1 = \frac{v_1}{\sqrt{2}} \begin{bmatrix} 1 \\ 0 \end{bmatrix} e^{ixa/f} ; \quad \Phi_2 = \frac{v_2}{\sqrt{2}} \begin{bmatrix} 0 \\ 1 \end{bmatrix} e^{ia/xf} . \quad (39)$$

Using Eq. (27), to guarantee invariance of $\mathcal{L}_{\text{Yukawa}}$, u_{R} and d_{R} transform under a $U(1)_{\text{PQ}}$ transformation as

$$u_{\text{R}} \rightarrow e^{-i\alpha x} u_{\text{R}} ; \quad d_{\text{R}} \rightarrow e^{-i\alpha/x} d_{\text{R}} . \quad (40)$$

Whence, it is easy to see that the current associated with the $U(1)_{\text{PQ}}$ symmetry is:³

$$\begin{aligned} J_{\text{PQ}}^\mu &= -f \partial^\mu a + x \sum_{i=1}^{N_g} \bar{u}_{\text{R}_i} \gamma^\mu u_{\text{R}_i} + \frac{1}{x} \sum_{i=1}^{N_g} \bar{d}_{\text{R}_i} \gamma^\mu d_{\text{R}_i} \\ &\quad + \frac{1}{x} \sum_{i=1}^{N_g} \bar{\ell}_{\text{R}_i} \gamma^\mu \ell_{\text{R}_i} . \end{aligned} \quad (41)$$

²In certain models, it is possible for λ_3 or λ_o to vanish identically.

³Here I assumed that the charged leptons ℓ_{R} have a Yukawa coupling involving Φ_2 .

This $U(1)_{\text{PQ}}$ current has both a QCD and an electromagnetic anomaly

$$\partial_\mu J_{\text{PQ}}^\mu = \xi \frac{\alpha_s}{8\pi} G_{a\mu\nu} \tilde{G}_a^{\mu\nu} + \xi_\gamma \frac{\alpha}{4\pi} F_{\mu\nu} \tilde{F}^{\mu\nu} , \quad (42)$$

where [4]

$$\xi = N_g \left(x + \frac{1}{x} \right) ; \quad \xi_\gamma = N_g \frac{4}{3} \left(x + \frac{1}{x} \right) . \quad (43)$$

The QCD anomaly renders a bit more difficult the calculation of the effective interactions of axions with light hadrons [13]. Perhaps the simplest way to proceed is by using an effective Lagrangian technique [12] to describe the interactions of pions and the η , both among themselves and with axions. Since π and η are the Nambu-Goldstone bosons associated with the approximate global $U(2)_{\text{R}} \times U(2)_{\text{L}}$ symmetry of QCD, their interactions are described by an effective chiral Lagrangian

$$\mathcal{L}_{\text{chiral}} = -\frac{1}{4} f_\pi^2 \text{Tr} \partial_\mu \Sigma^\dagger \partial^\mu \Sigma \quad (44)$$

with

$$\Sigma = \exp \frac{i}{f_\pi} [\vec{\tau} \cdot \vec{\pi} + \eta] . \quad (45)$$

This Lagrangian must be augmented by a mass breaking term that mimics the Yukawa interactions (38) and thus involves the axion field, along with an axion kinetic energy term

$$\mathcal{L}_{\text{axion}} = -\frac{1}{2} f_\pi^2 m_\pi^2 \text{Tr}[\Sigma A M + M^\dagger A^\dagger \Sigma^\dagger] - \frac{1}{2} \partial_\mu a \partial^\mu a , \quad (46)$$

where

$$A = \begin{bmatrix} e^{-ixa/f} & 0 \\ 0 & e^{-ia/xf} \end{bmatrix} ; \quad M = \begin{bmatrix} \frac{m_u}{m_u+m_d} & 0 \\ 0 & \frac{m_d}{m_u+m_d} \end{bmatrix} . \quad (47)$$

The presence of A guarantees the $U(1)_{\text{PQ}}$ invariance of $\mathcal{L}_{\text{axion}}$ since under $U(1)_{\text{PQ}}$, to mimic the quark transformations,

$$\Sigma \rightarrow \Sigma \begin{bmatrix} e^{ix\alpha} & 0 \\ 0 & e^{i\alpha/x} \end{bmatrix} . \quad (48)$$

The anomaly interactions, which break the above $U(2) \times U(2) \times U(1)_{\text{PQ}}$ symmetry through the coupling of gluons to axions and the η , serve to give an effective mass term to the field combination which couples to $G\tilde{G}$. One has:⁴

$$\mathcal{L}_{\text{anomaly}} = -\frac{1}{2} m^2 \left[\eta + \frac{f_\pi}{f} \frac{(N_g - 1)}{2} \left(x + \frac{1}{x} \right) a \right]^2 , \quad (49)$$

with $m^2 \simeq m_\eta^2 \gg m_\pi^2$.

The quadratic terms in $\mathcal{L}_{\text{axion}}$ along with $\mathcal{L}_{\text{anomaly}}$, when diagonalized, allows one immediately to compute the axion mass in the model, as well as the axion mixings with π^0 and η

$$a \simeq a_{\text{phys.}} - \xi_{a\pi} \pi_{\text{phys.}}^0 - \xi_{a\eta} \eta_{\text{phys.}} . \quad (50)$$

A simple calculation yields for the Peccei-Quinn model [10]:

⁴The factor of $(N_g - 1)$ in $\mathcal{L}_{\text{anomaly}}$ enters since here one must include only the contribution of the heavy quarks in the axion-gluon interactions, as the light quark interactions are explicitly accounted for in $\mathcal{L}_{\text{axion}}$.

$$\begin{aligned}
\lambda_m &= N_g \left(x + \frac{1}{x} \right) ; \\
\lambda_3 &= \frac{1}{2} \left[\left(x - \frac{1}{x} \right) - N_g \left(x + \frac{1}{x} \right) \frac{(m_d - m_u)}{(m_d + m_u)} \right] ; \\
\lambda_o &= \frac{1}{2} (1 - N_g) \left(x + \frac{1}{x} \right) .
\end{aligned} \tag{51}$$

Using these results, one can deduce readily the effective axion to two-photon coupling. The electromagnetic anomaly of the π^o and η fields

$$\mathcal{L}_{\pi,\eta} = \frac{\alpha}{4\pi} \left[\frac{\pi^o}{f_\pi} + \frac{5}{3} \frac{\eta}{f_\pi} \right] F_{\mu\nu} \tilde{F}^{\mu\nu} , \tag{52}$$

through the mixing with axions, gives an effective $a\gamma\gamma$ coupling coming from the light quark sector. To this coupling, one must add the heavy quark contribution. This is given by Eq. (42), except that for the parameter ξ_γ one should replace N_g by $N_g - 1$. Adding these contributions together yields finally, for the model,

$$K_{a\gamma\gamma} = N_g \left(x + \frac{1}{x} \right) \frac{m_u}{m_u + m_d} . \tag{53}$$

THE DEMISE OF VISIBLE AXION MODELS

By changing the detailed way in which the Higgs fields Φ_1 and Φ_2 couple to the quarks and leptons, one can obtain a variety of axion models [14]. All of these variant models, as well as the original Peccei-Quinn model, have been ruled out experimentally. I will not discuss in detail here the different evidence against these weak scale ($f = v$) axions, as this is reviewed in depth in Ref. [4]. I will, however, give a few examples to give a flavor both of the models and of the experiments used to rule these axions out.

If $f = v$, then the resulting axions are quite light, since $m_a = 25\lambda_m$ KeV. Unless $\lambda_m \gg 1$, then $m_a < 2m_e$ and these axions are also very long lived, since they can only decay via the $a \rightarrow 2\gamma$ process. These light, long lived, weak scale axions are ruled out by the non-observation of the process $K^+ \rightarrow \pi^+ a$. One predicts [12]

$$BR(K^+ \rightarrow \pi^+ a) \simeq 3 \times 10^{-5} \lambda_o^2 \tag{54}$$

while, experimentally, from KEK there is a bound [15]:

$$BR(K^+ \rightarrow \pi^+ \text{ nothing}) < 3.8 \times 10^{-8} . \tag{55}$$

One can avoid this bound in models where $\lambda_o = 0$.⁵ However, experiments looking for axion production via nuclear de-excitation in reactors are sensitive to both λ_o and λ_3 . The absence of a tell-tale $a \rightarrow 2\gamma$ signal downstream, although less stringent than the $K^+ \rightarrow \pi^+ a$ process, then serves to rule out this variant [16].

Weak scale axions can be short-lived if $m_a > 2m_e$. For this to happen, λ_m must be large and this necessitates either x or x^{-1} to be large. This possibility runs into difficulty with experiments looking for axions in quarkonium decays, $Q\bar{Q} \rightarrow a\gamma$. If the $U(1)_{\text{PQ}}$ assignments of the Higgs fields are as in the original Peccei-Quinn model [10], then the rate for $\psi \rightarrow a\gamma$ is proportional to x^2 and that for $\Upsilon \rightarrow a\gamma$ is proportional to x^{-2} . The present bounds on these processes [10] are only consistent with values of $x \sim O(1)$. Thus short-lived axion models to be viable must have variant quark couplings, [14] such that both the $\psi \rightarrow a\gamma$ and $\Upsilon \rightarrow a\gamma$ rates are either proportional to x^2 or x^{-2} .

These variant models are ruled out by a combination of other experiments. [12] The process $\pi^+ \rightarrow e^+ e^- e^+ \nu_e$ observed at SIN [18] with a branching ratio of $O(10^{-9})$ bounds the parameter λ_3 in these models, since the decay chain $\pi^+ \rightarrow a e^+ \nu_e$ followed by $a \rightarrow e^+ e^-$ would contribute to the signal. One finds [12] from these considerations that $\lambda_3 \leq 0.2$. On the other hand, the $I = 0$ to $I = 0$ M1 decay of the 3.58 MeV excited state of ^{10}B to the ground state [19] provides a bound on the isoscalar mixing parameter λ_o [12], $\lambda_o \leq 2$. Although there are models in which

⁵This requires that, effectively, only the first generation of quarks ($N_g = 1$) feel the $U(1)_{\text{PQ}}$ symmetry.

one can satisfy either one, or the other, of these two constraints, one cannot satisfy both because these parameters are themselves constrained by the relation [12]

$$|\lambda_3 - \lambda_o| \simeq \frac{m_a}{m_\pi} \frac{v}{f_\pi} \sqrt{\frac{m_u}{m_d}} \geq 15 . \quad (56)$$

INVISIBLE AXION MODELS

Although it was a sensible assumption to suppose that the $U(1)_{\text{PQ}}$ breaking scale f was the same as the weak scale v , this is not necessary. The dynamical adjustment of the strong CP angle $\bar{\theta} \rightarrow 0$ works for any scale f . If $f \gg v$, the resulting axions are very light ($m_a \sim 1/f$), very weakly coupled (coupling $\sim 1/f$) and very long lived ($\tau(a \rightarrow 2\gamma) \sim f^5$). Thus these axions are, apparently, invisible.

Because $f \gg v$ by assumption, in invisible axion models the $U(1)_{\text{PQ}}$ symmetry must be broken by an $SU(2) \times U(1)$ singlet VEV. Hence these models all introduce an $SU(2) \times U(1)$ singlet complex scalar field σ . The invisible axion is then, essentially, the phase of σ . Hence, concentrating again only on the axion degrees of freedom, one can write

$$\sigma = \frac{f}{\sqrt{2}} e^{ia/f} . \quad (57)$$

Broadly speaking, one can classify invisible axion models into two types depending on whether or not they have direct couplings to leptons. The, so called, KSVZ axions [20] are hadronic axions with only induced coupling to leptons. The, so called, DFSZ axions [21], on the other hand arise in models where axions naturally couple to leptons already at tree level. I describe below these two types of invisible axions in some more detail.

KSVZ Axions

In these models one assumes that the ordinary quarks and leptons are PQ singlets. The $SU(2) \times U(1)$ singlet field σ , however, interacts with some new heavy quarks X which carry $U(1)_{\text{PQ}}$ charge [20] via the Yukawa interaction:

$$\mathcal{L}_{\text{KSVZ}} = -h\bar{X}_L\sigma X_R - h^*\bar{X}_R\sigma^\dagger X_L . \quad (58)$$

The interactions of the KSVZ axions with ordinary quarks arises as a result of the chiral anomaly which induces a coupling

$$\mathcal{L}_{\text{anomaly}} = \frac{a}{f} \left[\frac{\alpha_s}{8\pi} G_a^{\mu\nu} \tilde{G}_{a\mu\nu} + 3q_X^2 \frac{\alpha_s}{4\pi} F_{\mu\nu} \tilde{F}^{\mu\nu} \right] , \quad (59)$$

where q_X is the electric charge of the heavy quarks X . Calculations analogous to the ones discussed earlier give the following axion parameters for the KSVZ model [4]

$$\lambda_m = 1 ; \quad \lambda_3 = -\frac{1}{2} \frac{(m_d - m_u)}{(m_d + m_u)} ; \quad \lambda_o = -\frac{1}{2} \quad (60)$$

and

$$K_{a\gamma\gamma} = 3q_X^2 - \frac{4m_d + m_u}{3(m_d + m_u)} . \quad (61)$$

DFSZ Axions

In this class of models [21], both the quarks and leptons carry $U(1)_{\text{PQ}}$ and hence one again needs two Higgs fields Φ_1 and Φ_2 . However, the quarks and leptons feel the effects of the axions only through the interactions that the field σ has with Φ_1 and Φ_2 in the Higgs potential. This interaction, conventionally, occurs through the term

$$\mathcal{L}_{\text{axion}} = \kappa \Phi_1^T C \Phi_2 (\sigma^\dagger)^2 + \text{h.c.} , \quad (62)$$

which serves to fix the PQ properties of σ relative to Φ_1 and Φ_2 .

The contribution of the axion field a in Φ_1 and Φ_2 is readily isolated [4]. Defining, as before, $v_1^2 + v_2^2 = v^2$, one has

$$\begin{aligned}\Phi_1 &= \frac{v_1}{\sqrt{2}} \exp\left(i\frac{2v_2^2 a}{v^2 f}\right) \begin{bmatrix} 1 \\ 0 \end{bmatrix} \\ &\equiv \frac{v_1}{\sqrt{2}} \exp\left(i\frac{X_1}{f} a\right) \begin{bmatrix} 1 \\ 0 \end{bmatrix}\end{aligned}\quad (63)$$

$$\begin{aligned}\Phi_2 &= \frac{v_2}{\sqrt{2}} \exp\left(i\frac{2v_1^2 a}{v^2 f}\right) \begin{bmatrix} 0 \\ 1 \end{bmatrix} \\ &\equiv \frac{v_2}{\sqrt{2}} \exp\left(i\frac{X_2 a}{f}\right) \begin{bmatrix} 0 \\ 1 \end{bmatrix},\end{aligned}\quad (64)$$

where $X_1 = 2v_2^2/v^2$; $X_2 = 2v_1^2/v^2$. These formulas correspond to what was done earlier (cf. Eq. (39)) with the replacements: $x \leftrightarrow X_1$; $x^{-1} \leftrightarrow X_2$. Using instead of f the scale

$$\tilde{f} = f/2N_g \quad (65)$$

in the formula for the axion mass and mixing parameters (Eqs. (35)-(37)) gives the same expression for the axion mass in the DFSZ and KSVZ models. Adopting this convention, the axion parameters for the DFSZ case are

$$\begin{aligned}\lambda &= 1; \quad \lambda_3 = \frac{1}{2} \left[\frac{X_1 - X_2}{2N_g} - \frac{(m_d - m_u)}{(m_d + m_u)} \right]; \\ \lambda_o &= \frac{1 - N_g}{2N_g}\end{aligned}\quad (66)$$

and

$$K_{a\gamma\gamma} = \frac{4}{3} - \frac{4m_d + m_u}{3(m_d + m_u)}. \quad (67)$$

Note that the DFSZ axion has a coupling to electrons given by

$$\mathcal{L}_{\text{aee}}^{\text{DFSZ}} = -i\frac{X_2}{2N_g} \frac{m_e}{f} a \bar{e} \gamma_5 e. \quad (68)$$

ASTROPHYSICAL AND COSMOLOGICAL CONSTRAINTS ON INVISIBLE AXIONS

Invisible axion models are constrained substantially by astrophysical and cosmological considerations, which restrict the allowed range for $f(\tilde{f})$ —or equivalently the range for the axion mass

$$m_a \simeq 6.3 \left[\frac{10^6 \text{ GeV}}{f} \right] \text{ eV}. \quad (69)$$

The astrophysical bounds on m_a arise, essentially, because axion emission removes energy from stars altering their evolution. Because the axion couplings to matter are inversely proportional to f , for large values of f (and therefore small axion masses) axions are less effective at cooling stars. Hence astrophysics provides a lower bound for f or, equivalently, an upper bound on m_a . If axions are sufficiently light, their emission eventually is irrelevant for the evolution of stars.

The bounds on invisible axions one derives from stellar evolution considerations are less restrictive for KSVZ axions. It turns out that the most effective cooling process is Compton axion production [22] $\gamma e \rightarrow ea$, which is proportional to the coupling of axions to electrons. This process is absent for KSVZ axions, since $g_{\text{eea}}^{\text{KSVZ}} = 0$. However, one can obtain a bound on these axions since they can remove energy from stars via the Primakoff process $\gamma Z, e \rightarrow Z, ea$ involving the axion to 2-photon coupling $K_{a\gamma\gamma}$.

Schematically, the energy lost by a star from the Compton process is given by

$$Q = \frac{1}{\rho_{\text{star}}} \int dn_e dn_\gamma |v| \sigma E_a, \quad (70)$$

where ρ_{star} is the stellar density and one integrates over the axion energy over the density of the initial states. For $m_a \ll T_{\text{star}} \sim O(\text{KeV})$, the only dependence on m_a are through the dependence of the interaction rate on f : $|v|\sigma \sim f^{-2} \sim m_a^2$. Bounds are then obtained by requiring that $Q \leq Q_{\text{nucl}}$ —the rate of nuclear energy generation ($Q_{\text{nucl}} \sim 10^2$ ergs/g sec).

From a detailed and careful study of how axion emission would affect stellar evolution, Raffelt [22] gives the following bounds on the two classes of invisible axions

$$(m_a)_{\text{DFSZ}} \leq \frac{10^{-2}}{X_2} \text{eV} ; \quad (m_a)_{\text{KSVZ}} \leq \frac{0.27}{K_{a\gamma\gamma}} \text{eV} . \quad (71)$$

Stronger bounds than these can be derived from the observation of neutrinos from SN 1987a. The bounds arise because if the axion luminosity is comparable to the neutrino luminosity ($\sim 10^{53}$ ergs/sec) during the core collapse, then the neutrino signal would be altered. It turns out that the dominant process for axion production during the collapse is axion bremsstrahlung off nucleons ($N + N \rightarrow N + N + a$). As a result the SN 1987a bounds one obtains are quite similar for KSVZ and DFSZ axions. The results that Turner [23] gives from his study of this issue are

$$(m_a)_{\text{DFSZ}} \leq \frac{1.7 \times 10^{-3}}{\xi_{\text{brem}}^2} \text{eV} ;$$

$$(m_a)_{\text{KSVZ}} \leq 8.4 \times 10^{-4} \text{eV} \quad (72)$$

where

$$\xi_{\text{brem}}^2 = 1.44 + \frac{1}{2}(X_1 - X_2 - 1.55)^2 . \quad (73)$$

Cosmology, on the other hand, provides an upper bound for f (\tilde{f}) or, equivalently, a lower bound for the axion mass. This bound was derived by a number of authors [24] and its origin is easy to understand. When the Universe goes through the $U(1)_{\text{PQ}}$ phase transition, at temperatures of order $T \sim f$, the axion field acquires a vacuum expectation. At this stage the color anomaly is not effective, so the axion is a Nambu-Goldstone boson and $\langle a_{\text{phys.}} \rangle \sim f$. As the Universe cools to a temperature $T^* \sim \Lambda_{\text{QCD}}$, the axion gets a mass of order $m_a \sim \Lambda_{\text{QCD}}^2/f$ and the axion VEV is driven to zero dynamically $\langle a_{\text{phys.}} \rangle \rightarrow 0$, corresponding to $\bar{\theta} = 0$. The relaxation of $\langle a_{\text{phys.}} \rangle$ to this value is oscillatory and this coherent oscillation of the zero-momentum component of the axion field contributes to the Universe's energy density. The larger f is, the larger the axion contribution to the energy density of the Universe is. Asking that this contribution not exceed the Universe's closure density then gives an upper bound on f :

A little more quantitatively [4], one can examine the equation of motion for $\langle a_{\text{phys.}} \rangle$ in the expanding Universe in the approximation that only the axion mass term is relevant in the axion potential:

$$\frac{d^2 \langle a_{\text{phys.}} \rangle}{dt^2} + 3 \frac{\dot{R}(t)}{R(t)} \frac{d \langle a_{\text{phys.}} \rangle}{dt} + m_a^2(t) \langle a_{\text{phys.}} \rangle = 0 , \quad (74)$$

where $R(t)$ is the cosmic scale parameter and $\dot{R}(t)/R(t) = H(t)$ is the Hubble constant. If $H(t) \gg m_a(t)$ then there are no oscillations, while in the reverse limit the oscillations are sinusoidal. Oscillations start at T^* , when

$$m_a(T^*) \sim H(T^*) \sim \frac{\Lambda_{\text{QCD}}^2}{M_{\text{Planck}}} . \quad (75)$$

At the start of oscillations, the energy density in the axion field is of order $\rho_a(T^*) \sim m_a^2(T^*) f^2$. If one assumes that $m_a(t)$ is slowly varying then one can show that [24] $\rho_a(t) \sim m_a(t)/R^3(t)$. Thus the contribution of axion oscillations to the Universe energy density today is of order

$$\rho_a = \rho_a(T^*) \left[\frac{m_a}{m_a(T^*)} \right] \left[\frac{R^3(T^*)}{R^3} \right] \sim \frac{\Lambda_{\text{QCD}}^3 T^3}{m_a M_{\text{Planck}}} , \quad (76)$$

where $T \sim 3^\circ K$ is the temperature of the Universe today. Requiring that ρ_a be less than the closure density of the Universe provides the lower bound on m_a . Amazingly, the above order of magnitude formula gives the same bounds on m_a that the more careful calculations give [4]

$$m_a \geq (10^{-5} - 10^{-6}) \text{eV} . \quad (77)$$

The cosmological and astrophysical constraints on invisible axions allow only a rather narrow window for the axion mass, roughly

$$10^{-6} \text{ eV} \leq m_a \leq 10^{-3} \text{ eV} . \quad (78)$$

If the axion has a mass near the lower limit, then axions play a very important role in the Universe and could be (part of) the Universe's cold dark matter. The cosmological bounds are subject to caveats, so it is useful to explore all masses below 10^{-3} eV. For instance, decay of axionic strings [25] may increase the contribution of axions to the present Universe's energy density by about a factor of 10^3 —thereby pushing the axion mass lower bound to the astrophysical limit. On the other hand, Linde [26] has argued that in inflationary models, it is quite possible that $\langle a_{\text{phys}}(f) \rangle \sim cf$, with $c \ll 1$. This would decrease the lower bound on the axion mass by a factor of c^2 .

These uncertainties notwithstanding, it is very worthwhile trying to detect a possible signal of invisible axions. Indeed, experiments are presently under way to try to detect such axions on the assumption that they constitute the dark matter component of our galactic halo. Under this assumption one can estimate both the presumed axion energy density, $\rho_a^{\text{halo}} \sim 5 \times 10^{-25} \text{ g cm}^{-3} \sim 300 \text{ MeV cm}^{-3}$, and their velocity, $v_a \simeq 10^{-3}$ —the virial velocity in the galaxy. The basic idea for these experiments was suggested more than a decade ago by Sikivie [27]. It makes use of the axion to 2-photon coupling to resonantly convert in a laboratory magnetic field the halo axions into photons, which can then be detected in a cavity.

The interaction of halo axions with a constant magnetic field B_o in a resonant cavity, as a result of Eq. (37), produces an electric field with frequency $\omega = m_a$. The generated electromagnetic energy can be detected in the cavity. When the cavity is tuned to the axion frequency, one should observe a narrow line on top of the noise spectrum. On resonance, one can write for the axion to photon conversion power the expression [28]

$$P = \left[\frac{\rho_a^{\text{halo}}}{m_a} \right] [V B_o^2] \left[\frac{\alpha}{\pi f} K_{a\gamma\gamma} \right]^2 C_{\text{over}} Q_{\text{eff}} . \quad (79)$$

Here the first factor is just the number of axions per unit volume; the second details the magnetic energy; the third is the axion coupling strength ($g_{a\gamma\gamma}^2$); the fourth is an overlap factor of $O(1)$ which depends in detail on which mode is excited in the cavity. Finally the last factor is the effective Q for the experiment, which is the minimum of the cavity's Q and that produced by the spread in the axion frequency due to its velocity. Typically both numbers are of $O(10^6)$.

Halo axions produce microwave photons ($1 \text{ GHz} = 4 \times 10^{-6} \text{ eV}$). Two pilot experiments, done at Brookhaven [29] and the University of Florida [30] had limited “magnetic energy”—typically $B_o^2 V \sim 0.5 \text{ (Tesla)}^2 m^3$ —and relatively noisy amplifiers. These experiments set limits for $g_{a\gamma\gamma}^2$ about a factor of 100-1000 away from theoretical expectations. There are now, however, two second generation experiments which should be sensitive at, or near, the theoretical limits.

The first of these experiments, which has just started taking data, is being carried out at Livermore National Lab [31]. It has a very large magnetic volume ($B_o V^2 \sim 12 \text{ (Tesla)}^2 m^3$) and state of the art amplifiers to reduce the noise level⁶. The second experiment, to be carried out in Kyoto University [32], uses a rather moderate magnetic volume ($B_o V^2 \sim 0.2 \text{ (Tesla)}^2 m^3$). However, it employs an extremely clever technique for counting the produced photons using Rydberg atoms, which is extremely sensitive. I show in the Figure the expectations of these experiments, along with the limits obtained by the pilot experiments. It is clear that these experiments have the capability to answer the question of the existence of axions in the mass range that they can probe. Results are expected in the next few years. So, stay tuned!

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⁶Even so the power produced by halo axions is tiny. At 1 GHz one expects $P_{\text{axion}} \sim 4 \times 10^{-20}$ Watts!

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