89. Axions and Other Similar Particles

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89.1 Introduction

In this section, we list coupling-strength and mass limits for light axion-like particles (ALPs) and QCD axions [1–4]. Such bosons may arise from the spontaneous breaking of a global U(1) Peccei-Quinn (PQ) symmetry, resulting in a massless Nambu-Goldstone (NG) boson. If there is a small explicit symmetry breaking, either already in the Lagrangian or due to quantum effects such as anomalies, the boson acquires a mass and is called a pseudo-NG boson. Axions may also arise in extra dimension constructions as the zero-modes of higher-dimensional gauge fields compactified on internal manifolds; in this case, the absence of a local contribution to the axion mass is due to the higher-dimensional gauge symmetry [5, 6].

A common feature of these light bosons a is that their coupling to Standard-Model particles is suppressed by the energy scale that characterizes the symmetry breaking, *i.e.*, the decay constant f_a. For example, in models where the axion arises as a Goldstone boson, the interaction Lagrangian at energy scales below f_a is

$$\mathcal{L} = f_a^{-1} J^\mu \partial_\mu a \,, \tag{89.1}$$

where J^μ is the Noether current of the spontaneously broken global symmetry. If f_a is very large, these new particles interact very weakly. Detecting them would provide a window to physics far beyond what can be probed at accelerators.

The QCD axions is of particular interest because it provides a compelling scheme to preserve CP-symmetry in QCD. Moreover, the cold dark matter (CDM) of the Universe may well consist of axions and they are searched for in dedicated experiments with a realistic chance of discovery.

Originally it was assumed that the PQ scale f_a was related to the electroweak symmetry-breaking scale $v_{\rm EW} = (\sqrt{2}G_{\rm F})^{-1/2} = 246.22$ GeV. However, the associated "standard" and "variant" axions were quickly excluded—we refer to the Listings for detailed limits. Here we focus on "invisible axions" with $f_a \gg v_{\rm EW}$ as the main possibility.

Axions have a characteristic two-photon vertex, inherited from their mixing with π^0 and η. This coupling allows for the main search strategy based on axion-photon conversion in external magnetic fields [7], an effect that also can be of astrophysical interest. While for axions the product "$a\gamma\gamma$ interaction strength $\times$ mass" is essentially fixed by the corresponding π^0 properties, one may consider a more general class of ALPs where the two parameters (coupling and mass) are independent. A number of experiments explore this more general ALP parameter space. ALPs populating the latter are predicted to arise generically, in addition to the axion, in low-energy effective field theories emerging from string theory [5, 6, 8–14].

89.2 Theory

89.2.1 *Peccei-Quinn mechanism and axions*

The QCD Lagrangian includes a CP-violating term $\mathcal{L}_\Theta = -\bar\Theta\,(\alpha_s/8\pi)\,G^{\mu\nu a}\tilde{G}^a_{\mu\nu}$, where $-\pi \le \bar\Theta \le +\pi$ is the effective Θ parameter after diagonalizing quark masses, $G^a_{\mu\nu}$ is the color field strength tensor, and $\tilde{G}^{a,\mu\nu} \equiv \epsilon^{\mu\nu\lambda\rho}G^a_{\lambda\rho}/2$, with $\varepsilon^{0123} = 1$, its dual. This term induces an electric dipole moment (EDM) of the neutron of strength $d_n = C_{\rm EDM}\, e\, \bar\Theta$. The coefficient $C_{\rm EDM}$ is calculated using QCD sum rules as $C_{\rm EDM} = 2.4(1.0) \times 10^{-16}$ cm [15], while lattice QCD computations find $C_{\rm EDM} = 1.48(14)(31) \times 10^{-16}$ cm [16] including statistical and systematic uncertainties. Experimental upper bounds on the latter, $|d_n| < 1.8 \times 10^{-26}\, e\,$cm [17, 18], imply $|\bar\Theta| \lesssim 10^{-10}$ even though

$\bar{\Theta} = \mathcal{O}(1)$ is otherwise completely satisfactory. Note that the proton has a similar EDM as the neutron, but it is less strongly constrained at present; this could change in the future [19]. The proton EDM may be written as $d_p = C^p_{\rm EDM}\, e\, \bar{\Theta}$, with $C^p_{\rm EDM} = -3.8(11)(8)$ [16]. The opposite signs of the neutron and proton EDMs may be understood from the opposite signs of their magnetic moments. The spontaneously broken global PQ symmetry $\mathrm{U(1)_{PQ}}$ was introduced to solve the "strong CP problem" [1, 2] of small Θ , the axion being the pseudo-NG boson of $\mathrm{U(1)_{PQ}}$ [3, 4]. This symmetry is broken due to the axion's anomalous triangle coupling to gluons,

$$\mathcal{L} = \left(\frac{a}{f_a} - \bar{\Theta}\right) \frac{\alpha_s}{8\pi}\, G^{\mu\nu a} \tilde{G}^a_{\mu\nu} \,, \tag{89.2}$$

where a is the axion field and f_a the axion decay constant. Color anomaly factors have been absorbed in the normalization of f_a which is defined by this Lagrangian. Thus normalized, f_a is the quantity that enters all low-energy phenomena [20]. Non-perturbative topological fluctuations of the gluon fields in QCD induce a potential for a whose minimum is at $a = \bar{\Theta}\, f_a$, thereby canceling the $\bar{\Theta}$ term in the QCD Lagrangian and thus restoring CP symmetry.

The resulting axion mass, in units of the PQ scale f_a, is identical to the square root of the topological susceptibility in QCD, $m_a f_a = \sqrt{\chi}$. The latter can be evaluated further [21, 22], exploiting the chiral limit (masses of up and down quarks much smaller than the scale of QCD), yielding $m_a f_a = \sqrt{\chi} \approx f_\pi m_\pi$, where $m_\pi \approx 135$ MeV and $f_\pi \approx 92$ MeV . In more detail one finds, by including $\mathcal{O}(\alpha)$ QED corrections and next-to-next-to-leading order (NNLO) corrections in chiral perturbation theory [23],

$$m_a = 5.691(51) \left(\frac{10^9\,\mathrm{GeV}}{f_a}\right) \mathrm{meV}\,. \tag{89.3}$$

A direct calculation of the topological susceptibility via QCD lattice simulations finds almost the same central value, albeit with an about five times larger error bar [24].

The QCD axion acquires a variety couplings to low-energy Standard Model observables, as we discuss in detail in this review. The most model-independent couplings are those of the axion to the neutron and proton EDM operators. While the axion removes the static neutron and proton EDMs, displacements of the axion field a from its vacuum expectation value induce field-dependent EDMs, which are characterized through the operator

$$\mathcal{L}_{aN\gamma} = -\frac{i}{2} g_{aN\gamma}\, a\, \overline{\Psi}_N \sigma_{\mu\nu}\gamma_5 \Psi_N F^{\mu\nu} \,, \tag{89.4}$$

with $N = n, p$ for neutrons or protons, respectively. The coupling $g_{an\gamma}$ is related to $C_{\rm EDM}$ by $g_{an\gamma} = eC_{\rm EDM}/f_a$ and similarly for $g_{ap\gamma}$. Direct probes of the axion-EDM operators are difficult, though by studying the thermal production of axions in the early universe one can broadly constrain, using Planck and Baryon Acoustic Oscillation (BAO) data, $m_a \lesssim 0.3$ eV [25]. In the presence of an oscillatory relic axion DM field, this operator induces an oscillating neutron EDM; this effect forms the basis of the CASPEr experiment [26, 27] that is discussed more later in this review. Non-linearly realized discrete-symmetries may be used to suppress the QCD axion mass below the prediction in Eq. (89.3), in which case a variety of cosmological and terrestrial constraints on the operator in Eq. (89.4) are also relevant [28–35]

Axions with $f_a \gg v_{\rm EW}$ evade almost all current experimental limits. Given that the axion interacts with ordinary matter through higher-dimensional operators, it requires an ultraviolet (UV) completion at an energy scale of order f_a. Broadly speaking, there are two distinct classes of UV scenarios that give rise to axions in the infrared (IR) low-energy effective field theory (EFT) (see [36] for a review of UV axion models): (i) 4d field theory models that realize the axion as a Goldstone mode of a spontaneously broken but also anomalous abelian $U(1)_{\rm PQ}$ symmetry, (ii) extra-dimensional models where the axion arises as the zero-mode of a higher-dimensional gauge field. Within the first category, one generic class of models is that of "hadronic axions," where new heavy quarks carry $\mathrm{U}(1)_{\rm PQ}$ charges, leaving ordinary quarks and leptons without tree-level axion couplings. The archetype is the KSVZ model [37, 38], where in addition the heavy new quarks are electrically neutral. Another generic class of field theory axion models requires at least two Higgs doublets and ordinary quarks and leptons carry PQ charges, the archetype being the DFSZ model [39, 40]; in this case the quarks and leptons may have tree-level axion couplings. All of these models contain at least one electroweak singlet scalar that acquires a vacuum expectation value and thereby breaks the PQ symmetry. The KSVZ and DFSZ models are frequently used as benchmark examples, but other field theory models exist where both heavy quarks and Higgs doublets carry PQ charges [36].

The field theory axion constructions, like KSVZ, DFSZ, and related variants, suffer from the so-called PQ quality problem [41–45]. The issue is that global symmetries are not expected to be respected in quantum gravity (see, *e.g.*, [46]); in particular, PQ-violating operators at the Planck scale spoil the quality of the $U(1)_{\rm PQ}$ symmetry, generating a neutron EDM in conflict with the experimental constraints. On the other hand, the axion quality is naturally exponentially better in the second class of UV completions mentioned in the previous paragraph, when the axion arises as the zero mode of an extra dimensional gauge field. In this case, the axion potential is protected by the gauge symmetry of the higher-dimensional theory, instead of a global symmetry. (More technically, the axion quality in extra dimensional constructions is best understood through the language of generalized global symmetries [47–49].) Contributions to the axion potential must be non-local and are thus naturally exponentially suppressed. Extra dimensional axions arise naturally in string theory constructions, as string theory has a plethora of higher-form gauge fields and lives most naturally in 10 or 11 dimensions, which requires some of those dimensions to be compactified. The presence of multiple distinct and non-trivial cycles in the compact manifold and higher-form gauge fields leads to a plethora of axions, which is known as the axiverse [6, 10, 12]—one of these axions can be the QCD axion. While field theory axions can have any $f_a \lesssim M_{\rm pl}$ and thus any mass larger than roughly 10^{-12} eV, axions in string theory most naturally have $m_a \lesssim 10^{-8}$ eV [6, 50], though higher QCD axion masses are also possible in these constructions at the cost of fine tuning [51, 52].

89.2.2 *Model-dependent axion couplings*

Although the generic axion interactions scale approximately with f_π/f_a from the corresponding π^0 couplings, there are non-negligible model-dependent factors and uncertainties. The axion's two-photon interaction plays a key role for many searches,

$$\mathcal{L}_{a\gamma\gamma} = -\frac{g_{a\gamma\gamma}}{4}\, a\, F_{\mu\nu}\tilde{F}^{\mu\nu} = g_{a\gamma\gamma}\, a\, \mathbf{E}\cdot\mathbf{B}\,, \tag{89.5}$$

where F is the electromagnetic field-strength tensor and $\tilde{F}^{\mu\nu} \equiv \epsilon^{\mu\nu\lambda\rho}F_{\lambda\rho}/2$, with $\varepsilon^{0123} = 1$, its dual. The coupling constant is [53]

$$g_{a\gamma\gamma} = \frac{\alpha}{2\pi f_a}\left(\frac{E}{N} - 1.92(4)\right) = \left(0.203(3)\frac{E}{N} - 0.39(1)\right)\frac{m_a}{\mathrm{GeV}^2}\,, \tag{89.6}$$

where E and N are the electromagnetic and color anomalies of the axial current associated with the axion, and α is the low-energy fine structure constant. In grand unified models, and notably for DFSZ [39, 40], $E/N = 8/3$, whereas for KSVZ [37, 38] $E/N = 0$ if the electric charge of the new heavy quark is taken to vanish. In general, a broad range of E/N values is possible [36, 54–56], as indicated by the diagonal yellow band in Fig. 89.1, whose width is roughly determined by the boundary values $E/N = 44/3$ and $E/N = 5/3$, respectively. However, this band still does not exhaust all the possibilities. In fact, there exist classes of QCD axion models whose photon couplings populate the entire still-allowed region outside the yellow band in Fig. 89.1 [57–62], motivating axion search efforts over a wide range of masses and couplings.

The two-photon decay width is

$$\Gamma_{a\to\gamma\gamma} = \frac{g^2_{a\gamma\gamma} m^3_a}{64\,\pi} \approx 1.1 \times 10^{-24}\ \mathrm{s}^{-1} \left(\frac{m_a}{\mathrm{eV}}\right)^5 . \tag{89.7}$$

The second expression uses Eq. (89.6) with $E/N = 0$. Axions decay faster than the age of the Universe if $m_a \gtrsim 20$ eV.

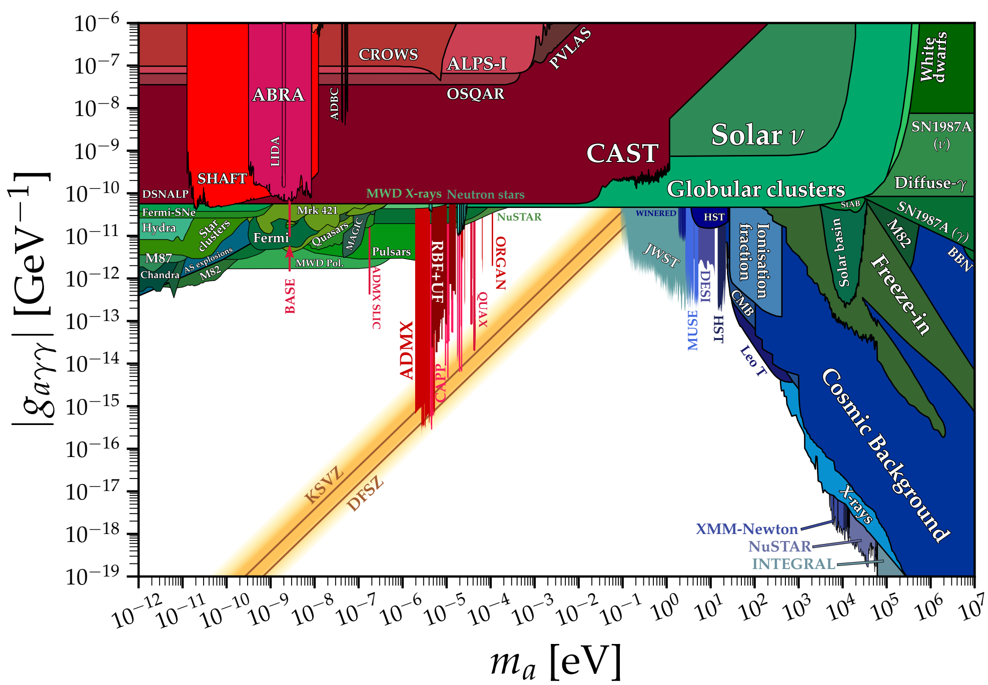

Figure 89.1: Exclusion plot for the ALP-photon coupling. Figure from [63], contains exclusions from refs. [64–182]. Constraints in red are from terrestrial experiments. Haloscope searches assume all local dark matter is axions. Those shown in green are astrophysical but make no assumption about axions being DM. See text for details regarding the diagonal yellow band expectation for the QCD axion and assumptions for blue high mass axion astrophysical limits.

The interaction with fermions f has derivative form and is invariant under a shift $a \to a + a_0$ as behooves a NG boson,

$$\mathcal{L}_{aff} = \frac{C_f}{2f_a}\,\partial_\mu a\,\bar{\Psi}_f \gamma^\mu \gamma_5 \Psi_f \,. \tag{89.8}$$

Here, Ψ_f is the fermion field, m_f its mass, and C_f a model-dependent coefficient. The dimensionless combination $g_{aff} \equiv C_f m_f / f_a$ plays the role of a Yukawa coupling and $\alpha_{aff} \equiv g^2_{aff}/4\pi$ of a "fine-structure constant." The often-used pseudoscalar form $\mathcal{L}_{aff} = -\mathrm{i}\,(C_f m_f/f_a)\,a\,\bar{\Psi}_f \gamma_5 \Psi_f$ need not be equivalent to the appropriate derivative structure, for example when two NG bosons are attached to one fermion line as in axion emission by nucleon bremsstrahlung [183–185].

In hadronic axion models, C_e vanishes at tree level, but is then generated radiatively at one loop [186, 187],

$$C_e \simeq \frac{3\alpha^2}{4\pi^2}\left[\frac{E}{N}\log\left(\frac{f_a}{m_e}\right) - 1.92\log\left(\frac{\Lambda_\chi}{m_e}\right)\right], \tag{89.9}$$

where $\Lambda_\chi \simeq 1$ GeV is the chiral symmetry breaking scale. In the DFSZ model [39, 40], the tree-level coupling coefficient to electrons is [186]

$$C_e = \frac{\sin^2\beta}{3}, \tag{89.10}$$

where $\tan\beta$ is the ratio of the vacuum expectation values of the two Higgs doublets giving masses to the up- and down-type quarks, respectively: $\tan\beta = v_u/v_d$. As a cautionary note, many of the original axion references, such as [186], defined $\tan\beta = v_d/v_u$, but the choice $\tan\beta = v_u/v_d$ adopted here and in more modern axion works is a reflection of the choice taken in the two-Higgs-doublet literature [188]. The resulting prediction for the axion-electron coupling, as a function of the axion mass, for the KSVZ axion ($E/N = 0$), hadronic axions, with $5/3 \leq E/N \leq 44/3$, and the DFSZ axion is displayed in Fig. 89.2. For the DFSZ range we have taken into account the constraint $0.28 \lesssim \tan\beta \lesssim 140$ [189] arising from the requirement of perturbative unitarity of the Yukawa couplings of Standard Model fermions.

For nucleons, $C_{p,n}$ have been determined as [53]

$$\begin{aligned}C_p &= -0.47(3) + 0.88(3)C_u - 0.39(2)C_d - 0.038(5)C_s\\ &\qquad - 0.012(5)C_c - 0.009(2)C_b - 0.0035(4)C_t\,,\\ C_n &= -0.02(3) + 0.88(3)C_d - 0.39(2)C_u - 0.038(5)C_s\\ &\qquad - 0.012(5)C_c - 0.009(2)C_b - 0.0035(4)C_t\,,\end{aligned} \tag{89.11}$$

in terms of the corresponding model-dependent quark couplings C_q, $q = u, d, s, c, b, t$. For hadronic axions, the latter vanish at tree-level, which means that C_n is expected to be much smaller than C_p. In the DFSZ model, $C_u = C_c = C_t = \frac{1}{3}\cos^2\beta$ and $C_d = C_s = C_b = \frac{1}{3}\sin^2\beta$, and C_p and C_n, as functions of β,

$$\begin{aligned}C_p &= -0.435\sin^2\beta + (-0.182 \pm 0.025)\,,\\ C_n &= 0.414\sin^2\beta + (-0.160 \pm 0.025)\,.\end{aligned} \tag{89.12}$$

The resulting prediction for the axion-neutron coupling, as a function of the axion mass, for the KSVZ axion and the DFSZ axion is displayed in Fig. 89.3. Note, however, that these couplings receive in-medium corrections, which may be important when considering compact stars [190].

The axion-pion interaction is given by the Lagrangian [187, 191]

$$\mathcal{L}_{a\pi} = \frac{C_{a\pi}}{2f_\pi f_a}\,\partial_\mu a\left(\pi^0\pi^+\partial_\mu\pi^- + \pi^0\pi^-\partial_\mu\pi^+ - 2\pi^+\pi^-\partial_\mu\pi^0\right), \tag{89.13}$$

where $C_{a\pi} = 2(C_p - C_n)/(3g_A)$ and $g_A = 1.2723(23)$ is the nucleon axial charge.

Additionally, there are axion-pion-nucleon interactions of the form

$$\mathcal{L}_{a\pi N} = \frac{C_{a\pi N}}{2f_\pi f_a}\partial_\mu a\left(i\pi^+\bar{p}\gamma^\mu n - i\pi^-\bar{n}\gamma^\mu p\right), \tag{89.14}$$

where $C_{a\pi N} = (C_{ap} - C_{an})/(\sqrt{2}g_A)$. Axion-kaon interactions are also present.

89.3 Laboratory Searches

89.3.1 *Light shining through walls*

Searching for "invisible axions" is extremely challenging due to their extraordinarily feeble coupling to normal matter and radiation. Currently, the most promising approaches rely on the axion-two-photon interaction, allowing for axion-photon conversion in external electric or magnetic fields [7]. For the Coulomb field of a charged particle, the conversion can be viewed as a scattering process, $\gamma + Ze \leftrightarrow Ze + a$, called the Primakoff effect [192]. In the other extreme of a macroscopic field, usually a large-scale B-field, the momentum transfer is small, the interaction is coherent over a large distance, and the conversion can be viewed as an axion-photon oscillation phenomenon in analogy to neutrino flavor oscillations [193].

Photons propagating through a transverse magnetic field, with incident $\mathbf{E}_\gamma$ and magnetic field $\mathbf{B}$ parallel, may convert into axions. For $m_a^2 L/2\omega \ll 2\pi$, where L is the length of the B field region and ω the photon energy, the resultant axion beam is coherent with the incident photon beam and the conversion probability is $\Pi \sim (1/4)(g_{a\gamma\gamma} BL)^2$. One such realization uses a laser beam propagating down the bore of a superconducting dipole magnet (like, say, the bending magnets in high-energy accelerators). If another magnet is in line with the first, but shielded by an optical barrier, then photons may be regenerated from the pure axion beam [194, 195]. The overall probability is $P(\gamma \to A \to \gamma) = \Pi^2$.

The first such Light-Shining-through-Walls (LSW) experiment was performed by the BFRT (Brookhaven-Fermilab-Rochester-Trieste) collaboration. It utilized two magnets of length L = 4.4 m and B = 3.7 T and found $|g_{a\gamma\gamma}| < 6.7 \times 10^{-7}$ GeV^{-1} at 95% CL for $m_a < 1$ meV [196, 197]. More recently, several such experiments were performed (see Listings) [99, 102, 198–202]. The current best limit, $|g_{a\gamma\gamma}| < 3.5 \times 10^{-8}$ GeV^{-1} at 95% CL for $m_a \lesssim 0.3$ meV (see Fig. 89.1), has been achieved by the OSQAR (Optical Search for QED Vacuum Birefringence, Axions, and Photon Regeneration) experiment, which exploited two 9 T LHC dipole magnets and an 18.5 W continuous wave laser emitting at the wavelength of 532 nm [102]. The ALPS I (Any Light Particle Search I) experiment achieved a similar sensitivity [99], see Fig. 89.1. Some of these experiments have also reported limits for scalar bosons where the photon $\mathbf{E}_\gamma$ must be chosen perpendicular to the magnetic field $\mathbf{B}$.

The concept of resonantly enhanced photon regeneration may open unexplored regions of coupling strength [203–205]. In this scheme, both the production and detection magnets are within Fabry-Perot optical cavities and actively locked in frequency. The $\gamma \to a \to \gamma$ rate is enhanced by a factor $\mathcal{F}\mathcal{F}'/\pi^2$ relative to a single-pass experiment, where $\mathcal{F}$ and $\mathcal{F}'$ are the finesses of the two cavities. The resonant enhancement could be of order $10^{(10-12)}$, improving the $g_{a\gamma\gamma}$ sensitivity by $10^{(2.5-3)}$. The experiment ALPS II (Any Light Particle Search II) is based on this concept and aims at an improvement of the current laboratory bound on $g_{a\gamma\gamma}$ by a factor $\sim 10^3$ [206]. ALPS II began operations in 2023, the regeneration cavity has been characterized [207], and a result from the first data taking campaign is pending.

Resonantly enhanced photon regeneration has already been exploited in experiments searching for "radiowaves shining through a shielding" [208–211]. For $m_a \lesssim 10^{-5}$ eV, the upper bound on $g_{a\gamma\gamma}$ established by the CROWS (CERN Resonant Weakly Interacting sub-eV Particle Search) experiment [101] is slightly less stringent than the one set by OSQAR, see Fig. 89.1.

89.3.2 *Long-range forces*

New bosons would mediate new long-range forces, which are severely constrained by "fifth force" experiments [212]. Those looking for new mass-spin couplings provide significant constraints on pseudoscalar bosons [213–218], see for example in Fig. 89.3 the limit on the axion-neutron coupling [219] from torsion balance tests of the gravitational inverse square law [220]. Presently, the most restrictive limits are obtained from combining long-range force measurements with stellar cooling arguments [221, 222]. For the moment, any of these limits are far from realistic values expected for the QCD axion. Still, these efforts provide constraints on more general low-mass bosons.

In Ref. [223], a method was proposed that can extend the search for axion-mediated spin-dependent forces by several orders of magnitude. By combining techniques used in nuclear magnetic resonance and short-distance tests of gravity, this method appears to be sensitive to the QCD axion in the μeV – meV mass range, independent of the cosmic axion abundance, if axions have a CP-violating interaction with nuclei as large as the current experimental bound on the electric dipole moment of the neutron allows. (Note that the CP-violation already present in the Standard Model is sufficient for generating an axion-mediated force, though in this case the signal appears too small to be detectable for the QCD axion in the foreseeable future [223].) The ARIADNE (Axion Resonant InterAction DetectioN Experiment) is under development and employs this approach to search for axion-mediated spin-dependent short-range interactions between a hyper-polarized ^{3}He sample and an unpolarized tungsten source mass [224]. The method relies on superconducting magnetic shielding to screen the sample from ordinary magnetic field noise. Experimental tests to demonstrate the requirements of ARIADNE, including characterization of the magnetic field backgrounds, are under way [225, 226].

89.4 Axions from Astrophysical Sources

89.4.1 *Stellar energy-loss limits*

Axions are produced in hot astrophysical plasmas, and can thus transport energy out of stars. The coupling strength of these particles with normal matter and radiation is bounded by the constraint that stellar lifetimes or energy-loss rates are not in conflict with observation [227, 228].

We begin this discussion with our Sun and concentrate first on the axion-photon coupling. In the Sun axions are produced by the Primakoff process $\gamma + Ze \to Ze + a$. Integrating this process over a standard solar model yields the axion luminosity [229]

$$L_a = g_{10}^2 \times 1.85 \times 10^{-3}\, L_\odot\,, \tag{89.15}$$

where $g_{10} = |g_{a\gamma\gamma}| \times 10^{10}$ GeV. The maximum of the spectrum is at 3.0 keV, the average at 4.2 keV, and the number flux at Earth is $g_{10}^2 \times 3.75 \times 10^{11}$ cm^{-2} s^{-1}. The solar photon luminosity is fixed, so energy losses due to the Primakoff process require enhanced nuclear energy production and thus enhanced neutrino fluxes. The all-flavor measurements by SNO (Sudbury Neutrino Observatory), together with a standard solar model, imply $L_a \lesssim 0.10\, L_\odot$, corresponding to $g_{10} \lesssim 7$ [230], mildly superseding a similar limit from helioseismology (sound speed, surface helium and convective radius) [231]. In Ref. [148], this limit was improved to $g_{10} < 4.1$ (at $3\,\sigma$), see Fig. 89.1, exploiting a new statistical analysis that combined helioseismology and solar neutrino observations, including theoretical and observational errors, and accounting for tensions between input parameters of solar models, in particular the solar element abundances. Going beyond the axion-photon coupling, Ref. [230] considered also a non-zero axion-electron coupling and obtained the bound on the latter displayed in Fig. 89.2.

A more restrictive limit derives from globular-cluster (GC) stars that allow for detailed tests of stellar-evolution theory. The stars on the horizontal branch (HB) in the color-magnitude diagram have reached helium burning with a core-averaged energy release of about 80 erg g^{-1} s^{-1}, compared to Primakoff axion losses of $g_{10}^2\, 30$ erg g^{-1} s^{-1}. The accelerated consumption of helium reduces the HB lifetime by about $80/(80 + 30\, g_{10}^2)$. In contrast, the red giant (RG) branch stars are not strongly affected by Primakoff losses, since unlike for the HB stars the plasma frequency in the RG stellar cores is below the temperature. Thus, Primakoff production reduces the R-parameter,

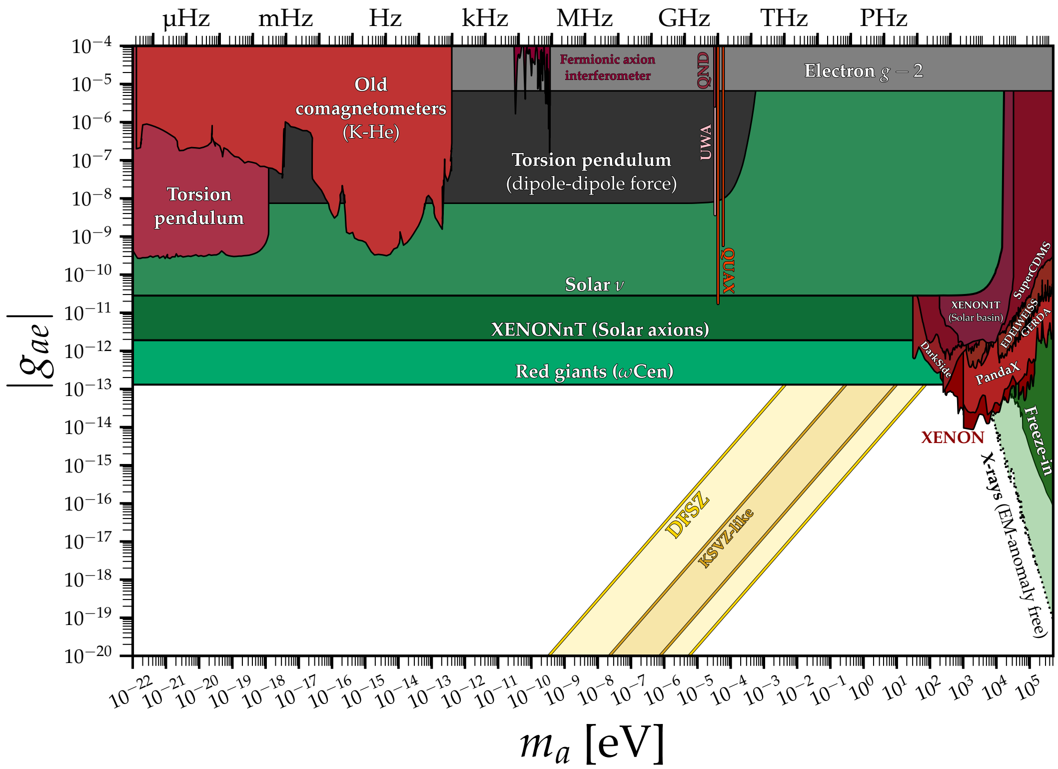

Figure 89.2: Exclusion plot for ALP-electron coupling as described in the text. Figure from [63], contains data from refs. [169, 230, 232–243]. The hadronic QCD axion-electron coupling arises at one loop and is given in Eq. (89.9); varying E/N between 5/3 and 44/3 gives the indicated hadronic axion band. The KSVZ axion has $E/N = 0$, while the DFSZ axion-electron coupling has the tree-level coupling given in Eq. (89.10) that depends on the unknown but bounded value of $\tan\beta$ (see text for details).

which is the number of observed HB branch stars to RG branch stars in a GC. Measurements of the R-parameter in a large sample of 39 Galactic GCs give a weak indication of non-standard losses which may be accounted by Primakoff-like axion emission, if the photon coupling is in the range $|g_{a\gamma\gamma}| = (2.9 \pm 1.8) \times 10^{-11}$ GeV^{-1} [123, 244]. Still, the upper bound found in this analysis,

$$|g_{a\gamma\gamma}| < 6.6 \times 10^{-11}\,\text{GeV}^{-1}\ (95\%\,\text{CL}), \tag{89.16}$$

represents one of the strongest limits on $g_{a\gamma\gamma}$ for a wide mass range. More recently, the R_2-parameter has emerged as a potentially more sensitive probe of axions; the R_2-parameter is defined as the ratio of the asymptotic giant branch (AGB) stars to the HB stars in a GC. AGB stars have more efficient Primakoff production than HB stars, meaning that R_2 is expected to decrease for increasing $|g_{a\gamma\gamma}|$. Ref. [124] used an R_2 measurement based off of 48 GCs as observed by the Hubble Space Telescope [245] to constrain

$$|g_{a\gamma\gamma}| < 4.7 \times 10^{-11}\,\text{GeV}^{-1}\ (95\%\,\text{CL})\,. \tag{89.17}$$

The constraint above on $g_{a\gamma\gamma}$ may be translated to $f_a > 4.8 \times 10^7$ GeV $(m_a < 0.12$ eV$)$, using $E/N = 0$ as in the KSVZ model, or to $f_a > 1.8 \times 10^7$ GeV $(m_a < 0.32$ eV$)$, for the DFSZ axion model, with $E/N = 8/3$, see Fig. 89.1.

If axions couple directly to electrons, the dominant emission processes are atomic axio-recombination and axio-deexcitation, axio-bremsstrahlung in electron-ion or electron-electron collisions, and Compton scattering [246]. Stars in the RG branch of the color-magnitude diagram of GCs are particularly sensitive to these processes: they would lead to an extension of the RG branch to larger brightness. In fact, the tip of the RG branch (TRGB) – the brightest point of the RG branch – provides the currently most sensitive method to test the axion coupling to electrons. The strongest bounds on it are derived from analyses of the TRGB in several GCs [247] and in the Galactic GC ω Centauri [239]. The two analyses lead to very similar results,

$$|g_{aee}| < 1.48 \times 10^{-13} \text{ (95\% CL) and } |g_{aee}| < 1.3 \times 10^{-13} \text{ (95\% CL)}, \tag{89.18}$$

respectively, see Fig. 89.2. Reference [247] finds also a small hint for extra cooling, corresponding to $|g_{aee}| = 0.60^{+0..32}_{-0.58} \times 10^{-13}$, while Ref. [239] does not find any evidence for exotic cooling.

Bremsstrahlung is also efficient in white dwarfs (WDs), where the Primakoff and Compton processes are suppressed by the large plasma frequency. A comparison of the predicted and observed luminosity function of WDs can be used to put limits on $|g_{aee}|$ [248, 249]. An analysis based on detailed WD cooling treatment and data on the WD luminosity function (WDLF) of the Galactic disk found that electron couplings above $|g_{aee}| \gtrsim 3 \times 10^{-13}$ are disfavoured [250]. Lower couplings cannot be discarded from the current knowledge of the WDLF of the Galactic disk [250–252]. These probes may be improved by the Rubin Observatory Large Synoptic Survey Telescope (LSST), which is expected to increase the sample of WDs in the Galactic halo to hundreds of thousands [253]. For pulsationally unstable WDs (ZZ Ceti stars), the period decrease $\dot{P}/P$ is a measure of the cooling speed. The corresponding observations of a handful pulsating WDs imply additional cooling that can be interpreted in terms of axion energy losses [254–258], though with axion-electron coupling values likely incompatible with other constraints [259].

Analogous constraints derive from neutron star (NS) cooling. Broadly speaking, there are two classes of NS cooling searches for axions. One approach uses the proto-NS formed after SN 1987A to constrain the axion luminosity produced during the first seconds of the proto-NS when the neutrino luminosity is measured with terrestrial experiments. The second class of NS cooling constraints consider older NSs, typically much older than ∼100 yr, and constrains the axion luminosity using the surface temperature as measured with photons versus the age of NS as measured through *e.g.* kinematic considerations. We begin with a discussion of the limits from SN 1987A.

Numerical simulations reveal that the energy-loss rate due to axions in a nuclear medium at the density 3×10^{14} g cm^{-3} and temperature 30 MeV should not exceed about 1×10^{19} erg g^{-1} s^{-1} to be consistent with the neutrino signal observed from SN 1987A [274]. The energy-loss rate from nucleon bremsstrahlung, $N + N \to N + N + a$, is $(C_N/2f_a)^2(T^4/\pi^2 m_N)\,F$. Here F is a numerical factor that represents an integral over the dynamical spin-density structure function because axions couple to the nucleon spin. For realistic conditions, even after considerable effort, one is limited to a heuristic estimate leading to $F \approx 1$ [228]. The SN 1987A limits are of particular interest for hadronic axions where the bounds on $|g_{aee}|$ are moot. Using a proton fraction of 0.3, $g_{ann} = 0$, $F = 1$, and $T = 30$ MeV, one finds $f_a > 4 \times 10^8$ GeV and $m_a < 16$ meV [228]. A more detailed numerical calculation [275] with state of the art SN models, again assuming $g_{ann} = 0$, found that a coupling larger than $|g_{app}| \gtrsim 6 \times 10^{-10}$, would shorten significantly the timescale of the neutrino emission. This result is, not surprisingly, rather close to the estimate in Ref. [228]. Improving the calculation of axion emission via nucleon-nucleon bremsstrahlung beyond the basic one-pion exchange approximation appears to loosen the bound [276,277]. The latter analysis finds a reduction of the axion emissivity by an order of magnitude if one takes into account the non-vanishing mass of the exchanged pion, the contribution from two-pion exchange, effective in-medium nucleon masses

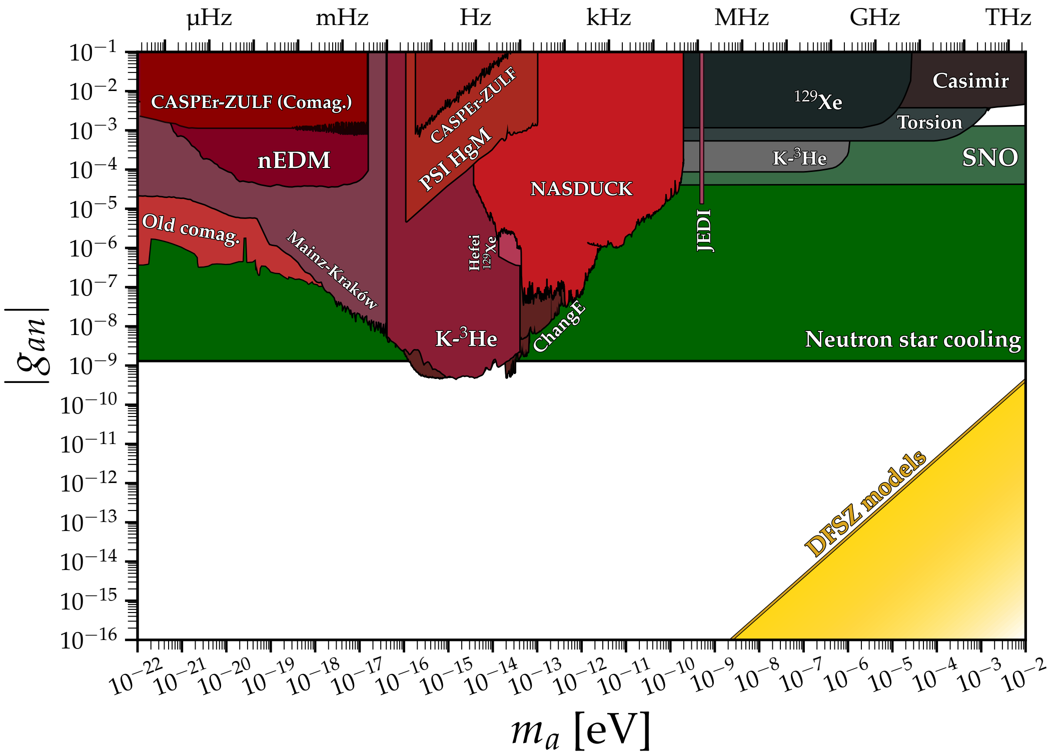

Figure 89.3: Exclusion plot for ALP-neutron coupling as described in the text. Figure from [63], includes data from refs. [31, 33, 219, 260–273]. The hadronic axion model prediction is given in Eq. (89.11) with vanishing quark couplings, while the DFSZ model prediction depends on $\tan\beta$ as is found in Eq. (89.12), giving the shaded yellow region above. Note that for a fine-tuned value of $\tan\beta$ g_{an} can be taken to zero. On the other hand, the neutron star cooling constraints [269] also probe the axion-proton coupling g_{ap} at a comparable level (not shown), and both g_{an} and g_{ap} cannot simultaneously be taken to zero in the DFSZ model.

and multiple nucleon scattering, leading to the result [277]

$$g_{ann}^2 + 0.61\, g_{app}^2 + 0.53\, g_{ann}\, g_{app} < 8.26 \times 10^{-19}\,, \tag{89.19}$$

which is consistent with the original estimate in [228]. This analysis, however, still neglects another very efficient mechanism for axion production in a SN: the pion production process, $\pi^- + p \to a + n$, also induced by the axion coupling to nucleons. Pion production was recognized long ago as being competitive with nucleon bremsstrahlung [278–280]. However, only recently it was shown that the pion abundance in the early stages of a SN is much larger than previously expected [281]. It was shown in Ref. [282] that pion induced processes dominate over bremsstrahlung and so strengthen the bound on the axion-nucleon coupling. On the other hand, given the modeling uncertainties and the sparse data from SN 1987A, the constraint on the axion-nucleon couplings from SN 1987A should be considered more as indicative than as a sharp bound [275]. Furthermore, note that if axions interact sufficiently strongly they are trapped. Only about three orders of magnitude in g_{aNN}, where $N = n, p$, or m_a are excluded. For even larger couplings, the axion flux would have been negligible, yet it would have triggered additional events in the detectors, excluding a further range [283].

Further bounds on the axion-nucleon coupling can be derived from observations of the cooling of older NSs. Recently, strong constraints were obtained by modeling the cooling of the "Magnificent

Seven" (M7) NSs in the presence of axions [269]. Axion emission would accelerate NS cooling. Thus, if the age and temperature of a NS are independently known, then by modeling axion-induced cooling in conjunction with the standard NS cooling mechanisms, in particular cooling via neutrino emission from the bulk of the NS and thermal emission from the NS surface, one is able to constrain the axion's coupling strength. Ref. [269] used a sample of four of the M7 along with PSR J0659 to constrain $m_a < 16$ meV for the KSVZ QCD axion at 95% CL and $m_a \lesssim 5$ meV ($m_a \lesssim 30$ meV) in the DFSZ scenario for asymptotically large (small) $\tan\beta$. For axions that couple exclusively to neutrons (protons) Ref. [269] constrains $|g_{ann}| < 1.3 \times 10^{-9}$ ($|g_{app}| < 1.5 \times 10^{-9}$), as illustrated in Fig. 89.3. The five NSs used in that work are unique in that their ages, all in the range $0.1 - 1$ Myr, are known by kinematic measurements (*e.g.*, tracing back to the SN remnant), while their surface temperatures are measured in the X-ray band. Beyond statistical uncertainties on these measurements, there are a number of confounding systematic uncertainties that must be accounted for in NS cooling analyses, including uncertainties related to the core-surface temperature relation, which involves the chemical composition of the NS envelope. The NS equation of state and the superfluid critical temperatures of the dense nuclear matter are also sources of uncertainty. Note that axions are predominantly produced by nuclear bremsstrahlung [284, 285] except for temperatures near the superfluid critical temperature, when Cooper pair breaking and formation (PBF) processes may efficiently produce axions [286, 287]. Given the large densities it is crucial to account for medium dependent couplings when computing the axion luminosities [288, 289].

The young NS in the supernova remnant Cassiopeia A (Cas A) was previously thought to strongly constrain axions, at a level similar to the NS cooling bound described above, due to PBF axion production [290]. This is because it was observed that Cas A was cooling rapidly, which may be explained by neutrino cooling with neutron superfluidity and proton superconductivity [291,292]. However, it is possible that the excessive cooling of Cas A is at least in part instrumental in origin [293]. Fully accounting for uncertainties on both the NS properties and the data, constraints from young NSs such as Cas A and the NS in the supernova remnant HESS J1731-347 [294] appear to be sub-dominant compared to the limits from SN 1987A and the M7 [269].

Finally, let us note that if $m_a \sim$ meV, SNe would lose a large fraction of their energy as axions. This would lead to a diffuse SN axion background in the Universe with an energy density comparable to the extra-galactic background light [295]. However, there is no apparent way of detecting it or the axion burst from the next nearby SN. On the other hand, neutrino detectors such as IceCube, Super-Kamiokande or a future mega-ton water Cherenkov detector will probe exactly the mass region of interest by measuring the neutrino pulse duration of the next Galactic SN [275].

89.4.2 *Searches for solar axions and ALPs*

Instead of using stellar energy losses to derive axion limits, one can also search directly for these fluxes, notably from the Sun. The main focus has been on ALPs with a two-photon vertex. They are produced by the Primakoff process with a flux given by Eq. (89.15) and an average energy of 4.2 keV, and can be detected at Earth with the reverse process in a macroscopic B-field ("axion helioscope") [7]. In order to extend the sensitivity in mass towards larger values, one can endow the photon with an effective mass in a gas, $\omega_{\rm plas}$, and tune the mass to match the axion's dispersion relation ($\omega_{\rm plas} = m_a$) [296].

An early implementation of these ideas used a conventional dipole magnet, with a conversion volume of variable-pressure gas with a xenon proportional chamber as x-ray detector [297]. The conversion magnet was fixed in orientation and collected data for about 1000 s/day. Axions were excluded for $|g_{a\gamma\gamma}| < 3.6 \times 10^{-9}$ GeV^{-1} for $m_a < 0.03$ eV, and $|g_{a\gamma\gamma}| < 7.7 \times 10^{-9}$ GeV^{-1} for $0.03 < m_a < 0.11$ eV at 95% CL.

Later, the Tokyo axion helioscope used a superconducting magnet on a tracking mount, viewing the Sun continuously. They reported $|g_{a\gamma\gamma}| < 6 \times 10^{-10}$ GeV^{-1} for $m_a < 0.3$ eV [298, 299]. This experiment was recommissioned and a similar limit for masses around 1 eV was reported [300].

The most recent helioscope CAST (CERN Axion Solar Telescope) uses a decommissioned LHC dipole magnet on a tracking mount. The hardware includes grazing-incidence x-ray optics with solid-state x-ray detectors, as well as x-ray Micromegas position-sensitive gaseous detectors. Exploiting an IAXO (see below) pathfinder system, CAST has established the limit

$$|g_{a\gamma\gamma}| < 5.8 \times 10^{-11}\ \text{GeV}^{-1} \quad (95\%\ \text{CL}), \tag{89.20}$$

for $m_a < 0.02$ eV [100, 182]. To cover larger masses, the magnet bores are filled with a gas at varying pressure. The runs with ^{4}He cover masses up to about 0.4 eV [301], providing the high-mass component of the CAST limits shown in Fig. 89.1. To cover yet larger masses, ^{3}He was used to achieve a larger pressure at cryogenic temperatures. Limits up to 1.17 eV allowed CAST to "cross the axion line" for the KSVZ model [302–304], see Fig. 89.1.

Sensitivity to significantly smaller values of $g_{a\gamma\gamma}$ can be achieved with a next-generation axion helioscope with a much larger magnetic-field cross section. Realistic design options for this "International Axion Observatory" (IAXO) have been studied in some detail [305], as well as its physics potential [306]. Such a next-generation axion helioscope may also push the sensitivity in the product of couplings to photons and to electrons, $g_{a\gamma\gamma}g_{aee}$, into a range beyond stellar energy-loss limits and test the hypothesis that WD, RG, and HB cooling is dominated by axion emission [307, 308]. As a first step towards IAXO, an intermediate experimental stage called BabyIAXO is currently under preparation at DESY [309].

Direct detection experiments searching for dark matter (DM) consisting of weakly interacting massive particles have also the capability to search for solar axions and ALPs. For low masses, $m_a \lesssim 100\,\text{eV}$, the XENONnT experiment [310] has provided the most stringent bound among those experiments on the axion-electron coupling constant,

$$|g_{aee}| < 1.9 \times 10^{-12} \quad (90\%\ \text{CL}), \tag{89.21}$$

see Fig. 89.2, by exploiting the axio-electric effect in liquid xenon. Limits at approximately the same value have been set by PandaX-4T [311] and LZ [312]. However, this technique has not reached the sensitivity of energy-loss considerations in stars.

89.4.3 *Conversion of astrophysical axion fluxes*

Large-scale B fields exist in astrophysics that can induce axion-photon oscillations. In practical cases, the strength of B is much smaller than in the laboratory, whereas the conversion region L is much larger. Therefore, while the product BL can be large, realistic sensitivities are usually restricted to very low-mass particles, far away from the "QCD axion band" in a plot like Fig. 89.1.

One example of this is SN 1987A, which would have emitted a burst of ALPs due to the Primakoff production in its core. It was pointed out shortly after the SN that the ALPs could have partially converted into γ-rays in the Galactic B-field. The lack of a gamma-ray signal in the GRS instrument of the SMM satellite in coincidence with the observation of the neutrinos emitted from SN 1987A therefore provides a strong bound on their coupling to photons [313, 314]. This bound has been revisited and the underlying physics has been brought to the current state-of-the-art, as far as modeling of the supernova and the Milky-Way along with stellar magnetic field are concerned, resulting in the limit [152, 153] $|g_{a\gamma\gamma}| < 5.2 \times 10^{-12}$ GeV^{-1} for $m_a \lesssim 5 \times 10^{-10}$ eV. More recently, it has been understood that the axions could convert to gamma-rays in the stellar magnetic fields of the progenitor star, extending the SN1987A limits to higher masses and potentially allowing gamma-ray observations of future SN to probe the QCD axion [315].

Reference [120] reports no evidence of a γ-ray burst in observations of extragalactic SNe with the Fermi Large Area Telescope (LAT). Under the assumption that at least one SN was contained within the LAT field of view, the authors derive an upper bound on the axion photon coupling which is about a factor of 5 weaker than the one from SN 1987A, see Fig. 89.1.

The cumulative emission of ALPs from all past core-collapse SNe would lead to a diffuse gamma-ray flux peaked at energies $\sim 50-100\,\mathrm{MeV}$ which can convert in the Galactic magnetic field into photons. Using Fermi-LAT measurements of the diffuse γ-ray flux, Ref. [316] obtains a conservative bound on the photon coupling, $|g_{a\gamma\gamma}| < 5\times10^{-10}\ \mathrm{GeV}^{-1}$, for $m_a \lesssim 10^{-11}\,\mathrm{eV}$, see see Fig. 89.1, which can be decreased by nearly three orders of magnitude, if an ALP-nucleon coupling of maximal phenomenologically allowed strength, $|g_{aNN}| \sim 10^{-9}$, is allowed.

Hot, young stars, such as Wolf-Rayet stars, efficiently produce ALPs with energies $\sim 10-100\,\mathrm{keV}$ via the Primakoff effect. Large numbers of those stars are hosted by the Galactic Quintuplet and Westerlund 1 super star clusters (SSCs). The non-observation of hard X-rays originating from axion-photon conversion in the Galactic magnetic field in archival NuSTAR (Nuclear Spectroscopic Telescope Array) data from these SSCs leads to a bound on the axion-photon coupling [160], $|g_{a\gamma\gamma}| < 3.6\times10^{-12}\,\mathrm{GeV}^{-1}$, for $m_a \lesssim 5\times10^{-11}\,\mathrm{eV}$, see Fig. 89.1. A somewhat weaker bound in the same mass range was established exploiting NuSTAR data on Betelgeuse [106]. Stronger constraints at low axion masses are found with NuSTAR data by considering axion production and conversion (or decay) in nearby galaxies, such as M82 and M87, accounting for production in all of the stellar members of the galaxies [317–319] (see Fig. 89.1).

A hard X-ray excess in data from the nearby M7 isolated NSs [320] may be explained by ALPs produced in the cores of those stars and converted in the surrounding magnetic fields, with $|g_{aNN}g_{a\gamma\gamma}| \in (2\times10^{-21}, 10^{-18})\,\mathrm{GeV}^{-1}$, for $m_a \lesssim 10^{-5}\,\mathrm{eV}$ [321]. The non-observation of an X-ray excess from the magnetic WD RE J0317-853 [322] by Chandra yields the constraint $|g_{aee}g_{a\gamma\gamma}| \lesssim 1.3\times10^{-25}\,\mathrm{GeV}^{-1}$, for $m_a \ll 10^{-5}\,\mathrm{eV}$ [129], which provides a non-trivial constraint on the ratio g_{aee}/g_{aNN} for axion models explaining the M7 X-ray excess.

In addition to being produced in NS cores, axions may also be produced in NS magnetospheres. In particular, it has recently been shown that axions could be copiously produced in the polar-cap regions of NS magnetospheres due to time-varying, un-screened electric fields [146]. Those axions could then convert to photons in the outer regions of the NS magnetosphere, leading to observable, broadband radio flux. The absence of such radio flux sets competitive constraints on ALPs, as indicated by the "pulsars" constraint in Fig. 89.1.

89.4.4 *Conversion of astrophysical photon fluxes*

Magnetically induced oscillations between photons and ALPs can modify the photon fluxes from distant sources in various ways, featuring (i) frequency-dependent dimming, (ii) modified polarization, and (iii) avoiding absorption by propagation in the form of axions. For example, dimming of SNe Ia could influence the interpretation in terms of cosmic acceleration [323], although it has become clear that photon-ALP conversion could only be a subdominant effect [324].

Polarization measurements are particularly powerful probes of photon-to-axion conversion since only photons polarized parallel to astrophysical magnetic fields may convert into axions, meaning that in the presence of large-scale ordered magnetic fields axions tend to leave distinct polarization signatures in passing radiation. Searches for linear polarization from radio galaxies (see, *e.g.*, Ref. [325]) and quasars [326] may approach in sensitivity $|g_{a\gamma\gamma}| \sim \text{few}\times10^{-13}\ \mathrm{GeV}^{-1}$, albeit with uncertainties related to the underlying modeling assumptions. Recently, robust limits at the level $|g_{a\gamma\gamma}| \le 1.7\times10^{-12}\ \mathrm{GeV}^{-1}$ at 95% CL have been derived from optical polarization measurements of nearby magnetic WDs [130, 327] (see also [328]), see Fig. 89.1. The magnetic fields of magnetic WDs may be well-measured through the Zeeman splitting of bound-bound transitions in the WD atmospheres.

It appears that the Universe could be too transparent to TeV γ-rays that should be absorbed by pair production on the extra-galactic background light [110, 329–332]. The situation is not conclusive at present [333–336], but the possibility of axions playing a role has been suggested [337–339]. On the other hand, much of this parameter space is excluded by the previously-discussed WD polarization probes. The region in ALP parameter space, $g_{a\gamma\gamma} \sim 10^{-12} - 10^{-10}$ GeV^{-1} for $m_a \lesssim 10^{-7}$ eV [340], required to explain the anomalous TeV transparency of the Universe, could be conceivably probed by the next generation of laboratory experiments (ALPS II) and helioscopes (IAXO) mentioned above.

This parameter region can also be probed by searching for an irregular behavior of the gamma ray spectrum of distant active galactic nuclei (AGN), expected to arise from photon-ALP mixing in a limited energy range. In these type of studies, the uncertainty in the magnetic field around the source needs to be taken into account. This typically leads to a range of limits on the ALP-photon coupling that depend on the modeling assumptions. The H.E.S.S. collaboration has set a limit of $|g_{a\gamma\gamma}| \lesssim 2.1\times10^{-11}$ GeV^{-1}, for 1.5×10^{-8} eV $\lesssim m_a \lesssim 6.0\times10^{-8}$ eV, from the non-observation of an irregular behavior of the spectrum of the AGN PKS 2155-304 [341]. The Fermi-LAT collaboration has put an even more stringent limit on the ALP-photon coupling [119] from observations of the gamma ray spectrum of NGC 1275, the central galaxy of the Perseus cluster, see Fig. 89.1. A similar analysis has been carried out in Ref. [342], using Fermi-LAT data of PKS 2155-304, and in Ref. [132], using ARGO-YBJ and Fermi-LAT data of Mrk 421, see Fig. 89.1. However, these constraints were obtained under certain theoretical assumptions about magnetic fields surrounding the sources, not confirmed yet by direct astronomical observations in these particular directions; this introduces large systematic uncertainties in the reported constraints [343].

Evidence for spectral irregularities has been reported in Galactic sources, such as pulsars and supernova remnants, and has been interpreted as hints for ALPs [344, 345] (see also discussion in Ref. [346] and references therein). However, the inferred ALP parameters, $|g_{a\gamma\gamma}| \sim 10^{-10}$ GeV^{-1}, $m_a \sim$ neV, are in tension with the CAST helioscope bounds. Reference [346] updated the analysis of the pulsar signal region including astrophysical nuisance parameters and correctly deriving confidence intervals on ALP parameters by means of Monte Carlo simulations. The tension with the CAST bounds can be evaded if environmental effects in matter, which would suppress the ALP production in dense astrophysical plasmas like the solar interior, are invoked. If this explanation is correct, the claimed ALP signal would be in reach of the next-generation LSW experiment ALPS II. Other ways to make the CAST bound compatible with photon-ALP conversions in the low-density Galactic medium are invoking photon-ALP-dark photon oscillations [347] or the existence of a number of 'hidden' ALPs [348].

At smaller masses, $m_a \lesssim 10^{-12}$ eV, galaxy clusters become highly efficient at interconverting ALPs and photons at X-ray energies. Constraints on spectral irregularities in the spectra of luminous X-ray sources (Hydra A, M87, NGC 1275, NGC 3862, Seyfert galaxy 2E3140; taken by Chandra and XMM-Newton) located in or behind galaxy clusters then lead to stringent upper limits on the ALP-photon coupling [111–113, 349–351] (for the limits exploiting spectra from Hydra and M87 from Refs. [111] and [112], respectively, see Fig. 89.1). Reference [113] recently performed the most sensitive x-ray searches for ALPs to date by employing Chandra's High-Energy Transmission Gratings that allow for an unsurpassed spectral resolution. New observations of the AGN NGC 1275 then led to the bound $|g_{a\gamma\gamma}| < 8 \times 10^{-13}$ GeV^{-1} at 99.7% C.L. for ultra-light ALPs, as illustrated in Fig. 89.1. On the other hand, magnetic field modeling uncertainties in the clusters may not be fully accounted for in these analyses [343].

89.4.5 *Finite density effects for the QCD axion*

If the QCD axion sector has a discrete, $\mathbb{Z}_{\mathcal{N}}$, shift symmetry, its potential is much shallower, and its mass, in units of its decay constant, $m_a f_a \simeq (3\pi)^{-1/4}\,\mathcal{N}^{3/4}\,2^{-\mathcal{N}/2}\,m_\pi f_\pi$, is smaller than the one of the canonical QCD axion by an exponential factor $\propto 2^{-\mathcal{N}/2}$ [59,60]. This opens up the parameter space over which one should search for a QCD axion towards the left of the yellow canonical QCD axion band in Fig. 89.1, Fig. 89.2, and Fig. 89.3. Novel bounds apply to this exceptionally light QCD axion due to finite density effects [352–354]. In fact, in dense stellar media, the minimum of the axion potential may be shifted to π. This has a number of phenomenological consequences that span from the modification of nuclear processes in stars due to $\theta = \mathcal{O}(\pi)$ to modifications in the orbital decay of binary systems (and subsequently in the emitted gravitational waves). Recently strong constraints have been set using the modification of the WD mass-radius relation due to such axions [355].

In fact, $\theta \sim 1$ in the solar core would lead to an increased proton-neutron mass difference, which would prohibit the neutrino line corresponding to the ^{7}Be-^{7}Li mass difference observed by Borexino [356].

Moreover, the fact that the position of the minimum of the axion potential depends on the nuclear density of the medium may also source a long-range force between dense stars [352]. This new long-range force sourced by the axion can be constrained by the measurement of the orbital decay of double pulsar or NS-pulsar binaries [352]. The gravitational wave signal of NS-NS mergers or black hole (BH) - NS mergers would also be modified by these axionic long-range forces [352,353]. There is also a corresponding bound [354] exploiting the gravitational waves observation from the binary NS inspiral GW170817 detected by LIGO (Laser Interferometer Gravitational-Wave Observatory) and Virgo [357].

89.4.6 *Superradiance of black holes*

Light bosonic fields such as axions or ALPs can affect the dynamics and gravitational wave emission of rapidly rotating astrophysical BHs through the superradiance mechanism. When the boson's Compton wavelength is of order of the size of the BH's ergoregion, they form gravitational bound states around the BH. Their occupation number grows exponentially by extracting energy and angular momentum from the black hole, forming a coherent axion or ALP bound state emitting gravitational waves. When accretion cannot replenish the spin of the BH, superradiance dominates the BH spin evolution; this is true for both supermassive and stellar mass BHs. The existence of destabilizing light bosonic fields thus leads to exclusion regions in the mass versus spin plot of rotating BHs. Stellar BH spin measurements exploiting well-studied binaries and two independent techniques exclude a mass range – up to self-interaction caveats discussed below – 6×10^{-13} eV $< m_a < 2\times10^{-11}$ eV at $2\,\sigma$, which for the axion excludes 3×10^{17} GeV $< f_a < 1\times10^{19}$ GeV [10,358,359], see Fig. 89.4. These bounds apply when gravitational interactions dominate over the axion self-interaction, which is largely true for the QCD axion in this mass range. On the other hand, Refs. [360,361] considered superradiance bounds for string theory ALPs with a quartic self-coupling (of either sign) given by $\lambda = m_a^2/f_a^2$, using the bosenova model of Ref. [359], and found that for $m_a \sim 10^{-12}$ eV the superradiant cloud, and thus the constraints, are destroyed for $f_a \lesssim 10^{11} - 10^{12}$ GeV, depending precisely on the mass and the source.

Reference [362] used a more advanced model for the effect of self-interactions on superradiant evolution, and a different (frequentist) statistical methodology, leading to more conservative bounds from stellar mass BHs (region near 10^{-12} eV). In particular, Ref. [362] found that QCD axions are only excluded for $m_a \lesssim 6\times10^{-12}$ eV instead of the 2×10^{-11} eV, as claimed in [360, 361]. Reference [362] also vetoed supermassive BHs (SMBHs) with spins measured at low significance, and thus found no ALP bounds in the SMBH region (near 10^{-18} eV); Ref. [360,361,363], in contrast, included SMBHs in a (quasi-)Bayesian analysis with a flat prior on the spin and found non-trivial constraints. In summary, it appears likely that the lower end of the QCD axion mass range is excluded by superradiance, with the requirement $m_a \gtrsim 10^{-11}$ eV, but the exact location of this limit is still debated. It is also too early to claim that ultra-light ALP parameter space is excluded by SMBHs, though this may change with future work.

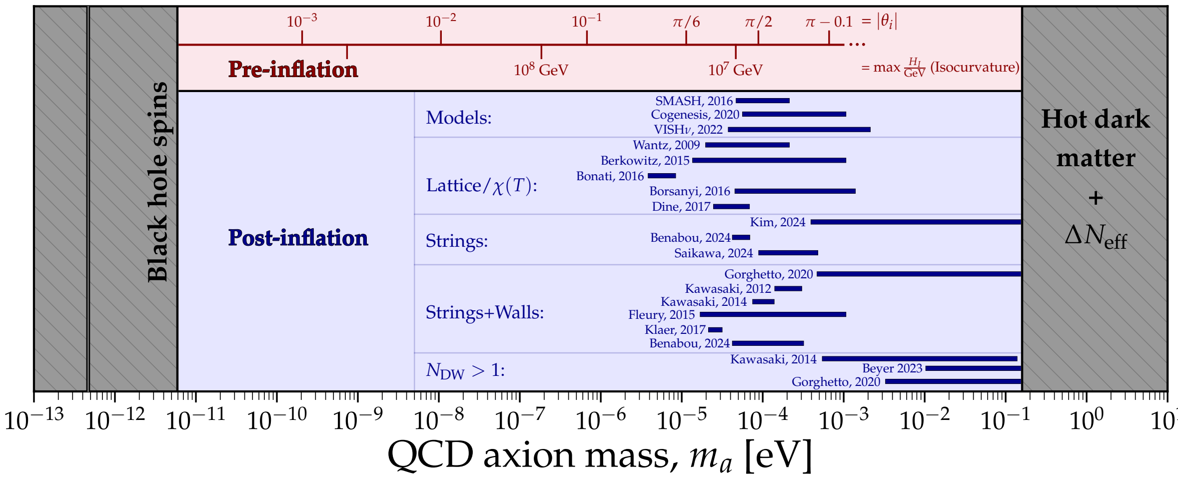

Figure 89.4: QCD axion mass predictions for obtaining the correct DM abundance assuming the PQ symmetry is and is not restored after inflation and, in the case the PQ symmetry is restored after inflation, depending on the domain wall number N. Figure from [63], includes data from refs. [24, 364–378]. The lower-end of the possible mass range is bounded by BH superradiance constraints and the theoretical requirement that f_a be below the Planck scale, while the upper end is bounded by the somewhat model-dependent NS cooling constraints (not shown), which exclude QCD axion masses above roughly ∼15-30 meV, depending on the UV completion. At even higher masses there are strong constraints from the production of DM in the early universe, which behaves as hot DM and contributes to ΔN_{eff}.

Long lasting, monochromatic gravitational wave signals, which can be distinguished from ordinary astrophysical sources by their clustering in a narrow frequency range, are expected to be produced by axions or ALPs annihilating to gravitons. Gravitational waves could also be sourced by axions/ALPs transitioning between energy levels of the effective gravitational atom. Accordingly, the gravitational wave detector Advanced LIGO should be sensitive to the axion in the $m_a \lesssim 10^{-10}$ eV region. LIGO measurements of BH spins in binary merger events could also provide statistical evidence for the presence of an axion [379, 380]. Similar signatures could arise for supermassive and intermediate mass BHs for particles with masses $\lesssim 10^{-15}$ eV. Gravitational waves from such sources could be detected at lower-frequency observatories such as LISA (Laser Interferometer Space Antenna). First attempts to use LIGO/Virgo data to search for monochromatic signatures of axion superradiance were recently made in [381, 382].

89.5 Cosmic Axions

89.5.1 *Cosmic axion populations*

There are two distinct populations of cosmic relic axion populations: a non-thermal one behaving as CDM, and a thermal one comprising a hot DM (HDM) component, in analogy to massive neutrinos. For $m_a \lesssim 0.01$ eV, thermal axions are dominantly produced by processes involving quarks and gluons [383, 384], while, for larger masses, the dominant thermalization process is $\pi + \pi \leftrightarrow \pi + a$ [187, 385, 386]. A number of evaluations exploiting cosmological precision data have found restrictive constraints on a possible HDM fraction that translate into an upper bound on the axion mass, $m_a \lesssim 1$ eV, if a leading order (LO) axion-pion chiral effective field theory (EFT) analysis of the axion-pion thermalization rate is used [387–390]. Recently, it was found that the next leading order (NLO) contribution exceeds the LO contributions for masses below $m_a \lesssim 1.2\,(0.12/C_{a\pi})\,\mathrm{eV}$ [391]. Therefore, in order to assess the reach in sensitivity of future cosmological data sets, one has to improve the EFT description or to compute axion-pion scattering via lattice QCD techniques. Alternatively, recently Ref. [392] used pion-pion scattering data to overcome the EFT description issues and found the constraint $m_a \lesssim 0.24$ eV, see Fig. 89.4.

For $m_a \gtrsim 10$ eV, axions may decay faster than the age of the Universe (see Eq. (89.7)), depending on the decay constant, which would remove the cosmic axion population while injecting photons. If the axions were in thermal equilibrium in the early Universe, which depends on the inflationary reheat temperature and assumptions about the pre-BBN cosmology [393, 394], this excess radiation provides additional limits on the axion- or ALP-photon coupling up to very large masses [393, 394]. Such photons could contribute to the extragalactic background light (EBL), could ionize primordial hydrogen and thus contribute to the optical depth after recombination, or could spoil the agreement of big bang nucleosynthesis (BBN) with observations. The corresponding up-to-date bounds from Refs. [167, 168] are displayed in Fig. 89.1. It has recently been shown that there is an *irreducible axion background* [169] that arises from axion freeze-in before BBN; decays of that axion population provide strong constraints on axions with masses above roughly 100 eV (see Fig. 89.1). For high enough decay constants, high-mass axions ($m_a \gtrsim 10$ eV) can account for the DM abundance through the misalignment mechanism, described in more detail below, though additional entropy dilution or fine-tuning is needed to avoid overproducing the DM abundance [395] (see also [396] for heavy axion production through lepton-flavor-violating decays). The decay of such high-mass axions to photon pairs could leave narrow spectral features in *e.g.* optical or X-ray backgrounds [127, 139–141, 162, 164, 397, 398] (see Fig. 89.1).

The main cosmological interest in axions at present derives from their possible role as CDM. In addition to thermal processes, axions are abundantly produced by the misalignment (MIS) mechanism [399–401]. The axion DM abundance crucially depends on the cosmological history. Let us first consider the so called *pre-inflationary PQ symmetry breaking scenario*, in which the PQ symmetry is broken before and during inflation and not restored afterwards. This is the cosmological scenario most naturally realized for string theory axions from the open-string sector [47, 402]. After the breakdown of the PQ symmetry, the axion field relaxes somewhere in the bottom of the "wine-bottle-bottom" potential. Near the QCD epoch, topological fluctuations of the gluon fields such as instantons explicitly break the PQ symmetry. This tilting of the "wine-bottle-bottom" drives the axion field toward the CP-conserving minimum, thereby exciting coherent oscillations of the axion field that ultimately represent a condensate of CDM. The fractional cosmic mass density in this homogeneous field mode, created by the MIS mechanism, is [24, 403–405],

$$\begin{aligned} \Omega_a^{\mathrm{MIS}} h^2 &\approx 0.12 \left(\frac{f_a}{9\times 10^{11}\ \mathrm{GeV}}\right)^{1.165} F\,\Theta_{\mathrm{i}}^2 \\ &\approx 0.12 \left(\frac{6\ \mu\mathrm{eV}}{m_a}\right)^{1.165} F\,\Theta_{\mathrm{i}}^2\,, \end{aligned} \tag{89.22}$$

where h is the present-day Hubble expansion parameter in units of 100 km s^{-1} Mpc^{-1}, and $-\pi \leq \Theta_{\mathrm{i}} \leq \pi$ is the initial "misalignment angle" relative to the CP-conserving position attained in the causally connected region which evolved into today's observable Universe. $F = F(\Theta_{\mathrm{i}}, f_a)$ is a factor accounting for anharmonicities in the axion potential. For $F\Theta_{\mathrm{i}}^2 = \mathcal{O}(1)$, m_a should be above $\sim 6\,\mu$eV in order that the cosmic axion density does not exceed the observed CDM density, $\Omega_{\rm CDM}h^2 = 0.12$. However, much smaller axion masses (much higher PQ scales) are still possible if entropy is diluted for example by the late decay of a scalar condensate (see for example [406] and references therein) or the initial value $\Theta_{\rm i}$ just happens to be small enough in today's observable Universe ("anthropic axion window" [407]). In the latter cosmological scenario, however, quantum fluctuations of the axion field during inflation are expected to lead to isocurvature density fluctuations which get imprinted to the temperature fluctuations of the CMB [408,409]. Their non-observation puts severe constraints on the Hubble expansion rate H_I during inflation [410–415], which read, in the simplest cosmological inflationary scenario,

$$H_I \lesssim 5.7\times 10^8\,{\rm GeV}\left(\frac{5\,{\rm neV}}{m_a}\right)^{0.4175}, \tag{89.23}$$

if axions represent all of DM. In Ref. [416] an alternative to the MIS mechanism was proposed: the so-called kinetic MIS mechanism. In this case, the axion field is assumed to have an initial velocity which may be generated, e.g., by a hypothesized explicit breaking of the axion shift symmetry in the early universe. The amount of QCD axion dark matter generated by the kinetic MIS mechanism can fit the observed dark matter abundance for any $f_a \lesssim 1.5\times 10^{11}\,$GeV, down to the minimum value allowed by the SN 1987A constraint.

In the above prediction of the fractional cosmic mass density in axions, the exponent, 1.165, arises from the non-trivial temperature dependence of the topological susceptibility $\chi(T) = m_A^2(T)f_A^2$ at temperatures slighty above the QCD quark-hadron phase transition. Lattice QCD calculations of this exponent [24,367,369,417–419], but also Ref. [368], found it to be close to the prediction of the dilute instanton gas approximation [420] which was previously exploited. Therefore, the state-of-the-art prediction of the axion mass relevant for DM for a fixed initial misalignment angle $\Theta_{\rm i}$ differs from the previous prediction by just a factor of order one.

In the *post-inflationary PQ symmetry breaking scenario*, on the other hand, $\Theta_{\rm i}$ will take on different values in different patches of the present Universe. The average contribution is [24,403–405]

$$\Omega_a^{\rm MIS}h^2 \approx 0.12\left(\frac{30\ \mu{\rm eV}}{m_a}\right)^{1.165}. \tag{89.24}$$

The decay of cosmic strings and domain walls gives rise to a further population of CDM axions, whose abundance suffers from significant uncertainties [365,371–373,404,405,421–432] which arise from the difficulty in understanding the energy loss process of topological defects and the generated axion spectrum in a quantitative way. In fact, in the present state-of-the-art it is still possible that the CDM contribution from the decay of topological defects is subdominant [371] or overwhelmingly large [372] in comparison to the one from the MIS mechanism. Correspondingly, the plausible range of axion masses providing all of CDM in scenarios with postinflationary PQ symmetry breaking is still rather large, namely [371,372]

$$m_a \approx 26\ \mu{\rm eV} - 0.5\ {\rm meV}, \tag{89.25}$$

for models with short-lived (requiring unit color anomaly $N = 1$) domain walls, such as the KSVZ model. On the other hand, the most recent adaptive mesh refinement axion string simulations, which are able to evolve over larger dynamic ranges than static lattice simulations, suggest the mass range $m_a \approx 45 - 300\ \mu$eV, with domain wall decay providing the dominant contribution to DM on the upper end of the uncertainty range [432]. For models with long-lived ($N > 1$) domain walls, such as an accidental DFSZ model [433], where the PQ symmetry is broken by higher dimensional Planck suppressed operators, the mass is predicted to be significantly higher [306,373,433–435],

$$m_a \approx (0.58 - 130)\,{\rm meV}\,. \tag{89.26}$$

However, the upper part of the predicted range is in conflict with stellar energy-loss limits on the axion. The pre- and post-inflationary mass predictions are summarized in Fig. 89.4.

In this post-inflationary PQ symmetry breakdown scenario, the spatial axion density variations are large at the QCD transition, and they are not erased by free streaming. Gravitationally bound "axion miniclusters" form before and around matter-radiation equality [436–438]. A significant fraction of CDM axions can reside in these bound objects [426,439]. The minicluster fraction can be bounded by gravitational lensing [440–442], although more simulations are required to understand whether miniclusters are dense enough and survive in sufficient quantities for lensing bounds to apply.

89.5.1.1 *Ultralight cosmological ALP relics*

The non-thermal production mechanisms attributed to axions are generic to light bosonic weakly interacting particles such as ALPs [376]. The relic abundance is set by the epoch when the axion mass becomes significant, $3H(t) \approx m_a(t)$, and ALP field oscillations begin. For ALPs to contribute to the DM density this epoch must precede that of matter radiation equality. For a temperature independent ALP mass this leads to the bound:

$$m_a \gtrsim 7 \times 10^{-28}\,\mathrm{eV} \left(\frac{\Omega_m h^2}{0.15}\right)^{1/2} \left(\frac{1+z_{\mathrm{eq}}}{3.4 \times 10^3}\right)^{3/2} . \tag{89.27}$$

ALPs lighter than this bound are allowed if their cosmic energy density is small, but they are quite distinct from other forms of DM [443]. Ignoring anharmonicities in the ALP potential, and taking the ALP mass to be temperature independent, the relic density in DM ALPs due to the MIS mechanism is given by

$$\Omega_{\mathrm{ALP}}^{\mathrm{MIS}} h^2 = 0.12 \left(\frac{m_a}{4.7 \times 10^{-19}\,\mathrm{eV}}\right)^{1/2} \left(\frac{f_a}{10^{16}\,\mathrm{GeV}}\right)^2 \left(\frac{\Omega_m h^2}{0.15}\right)^{3/4} \left(\frac{1+z_{\mathrm{eq}}}{3.4 \times 10^3}\right)^{-3/4} \Theta_{\mathrm{i}}^2 . \tag{89.28}$$

An ALP decay constant near the GUT scale gives the correct relic abundance for *ultralight ALPs* (ULAs) with $m_{\mathrm{ULA}} \approx 10^{-19}$ eV [444, 445]. ULAs encompass the entire Earth in a single Compton wavelength, and for large occupation numbers are modelled as a coherent classical field. The coherence time is determined by the mass and virial velocity in the Milky Way, $\tau_{\mathrm{coh}} \sim 1/m_{\mathrm{ULA}} v_{\mathrm{vir}}^2$, with the detailed properties described by a stochastic model with an approximately Rayleigh Jeans distribution [446]. Natural models for ULAs can be found in string and M-theory compactifications [5, 6, 8–13], in field theory with accidental symmetries [447, 448], or new hidden strongly coupled sectors [449, 450].

In addition to the gravitational potential energy, the ULA field also carries gradient energy. On scales where the gradient energy is non-negligible, ULAs acquire an effective pressure and do not behave as CDM. The gradient energy opposes gravitational collapse, leading to a Jeans scale below which perturbations are stable [451]. The Jeans scale suppresses linear cosmological structure formation relative to CDM [452–454]. The Jeans scale at matter-radiation equality in the case that ULAs make up all of CDM is:

$$k_{\mathrm{J,eq}} = 8.7\,\mathrm{Mpc}^{-1} \left(\frac{1+z_{\mathrm{eq}}}{3.4 \times 10^3}\right)^{-1/4} \left(\frac{h^2 \Omega_{\mathrm{ALP}}^{\mathrm{MIS}}}{0.12}\right)^{1/4} \left(\frac{m_{\mathrm{ULA}}}{10^{-22}\,\mathrm{eV}}\right)^{1/2} . \tag{89.29}$$

On non-linear scales the gradient energy leads to the existence of a class of pseudo-solitons known as oscillatons, or axion stars [455]. Axion stars are expected to form in all cosmological scenarios, and for all types of ALPs, including the QCD axion [456–458]. Some axion stars may explode into radio photons, heating the interstellar medium [105].

Cosmological and astrophysical observations are consistent with the CDM model, and departures from it are only allowed on the scales of the smallest observed DM structures with $M \sim 10^{6-8} M_\odot$. The CMB power spectrum and galaxy auto-correlation power spectrum limit the ULA mass to $m_{\rm ULA} \gtrsim \mathcal{O}(\text{few}) \times 10^{-24}$ eV from linear theory of structure formation [443, 459]. Analytic models [460] and N-body simulations [461] for non-linear structures show that halo formation is suppressed in ULA models relative to CDM. This leads to constraints on the ULA mass of $m_{\rm ULA} > 10^{-22}$ eV from observations of high-z galaxies [461–463], and $m_{\rm ULA} > 2 \times 10^{-20}$ eV from the Lyman-alpha forest flux power spectrum [464]. Including the effects of anharmonicities on structure formation with ALPs can weaken these bounds if the misalignment angle $\Theta_{\rm i} \approx \pi$ [465]. A comprehensive study of Milky Way satellites by the DES collaboration resulted in the bound $m_{\rm ULA} > 2.9 \times 10^{-21}$ eV [466].

Cosmological simulations that treat gradient energy in the ULA field beyond the N-body approximation have just recently become available [456, 467–469], and show, among other things, evidence for the formation of axion stars in the centres of ULA halos (various consequences of axion stars are considered in Refs. [470]). These central axion stars have been conjectured to play a role in the apparently cored density profiles of dwarf spheroidal galaxies, and other central galactic regions [128, 456, 471–475]. However, the relationship between the halo mass and the axion star mass [476] leads to problems with this scenario in some galaxies [477–479]. It should be emphasized that many of the conclusions about the role of ULA axion stars in galactic dynamics are based on use of simulation results that do not contain baryons and feedback could be important [480–482].

Inside DM halos the axion gradient energy causes coherence on the de Broglie wavelength and fluctuations on the coherence time [456, 483]. These fluctuations can be thought of as short-lived quasiparticles and lead to relaxation processes that can be described statistically [445, 484] (this relaxation processes also leads to the gravitational condensation of axion stars [457]). The typical relaxation time is:

$$t \sim 10^{10} \text{ years} \left(\frac{m_{\rm ULA}}{10^{-22}\text{ eV}}\right)^3 \left(\frac{v}{100\text{ km s}^{-1}}\right)^2 \left(\frac{r}{5\text{ kpc}}\right)^4 , \tag{89.30}$$

where v and r are the velocity and radius of the orbit in the host DM halo.

Relaxation processes such as these are not observed in galaxies, though there are some circumstances where they may be desirable [445]. An absence of observed relaxation can be used to set limits on the ULA mass. An absence of Milky Way disk thickening excludes $m_{\rm ULA} > 0.6 \times 10^{-22}$ eV [485], while stellar streams give the stronger bound $m_{\rm ULA} > 1.5 \times 10^{-22}$ eV [486]. The survival of the old star cluster in Eridanus II [487] excludes the range of masses 10^{-21} eV $\lesssim m_{\rm ULA} \lesssim 10^{-19}$ eV [488]. Recently Ref. [489] found that ULA DM is excluded for $m_{\rm ULA} \lesssim 3 \times 10^{-19}$ eV at 99% CL because such DM would kinematically heat the stars in the Segue I and Segue II ultra-faint dwarf galaxies to a level inconsistent with their observed dynamics.

Finally, one should note that the beyond-CDM physics of ULAs (Jeans scale, relaxation, axion star formation) of course also applies to the QCD axion on smaller length scales. This is of particular interest inside axion miniclusters [436, 437, 457, 458].

89.5.2 *Electron recoil searches*

In a DM direct detection experiment, a DM ALP featuring a coupling to the electron can be absorbed by the target material, leading to a mono-energetic electronic recoil signal peaked at m_a. This mechanism allowed the EDELWEISS [490], PandaX [234], SuperCDMS [235] and XENON1T [491] collaborations to put the bounds on the axion-electron coupling displayed in Fig. 89.2 in the keV mass range. As discussed in the last sub-section, additionally dynamics beyond the standard misalignment mechanism are needed to explain the axion DM abundance for keV-scale axion masses (*e.g.*, [395,396]). Moreover, such models are subject to stringent constraints from axion decays to photons [395].

89.5.3 *Telescope searches*

The two-photon decay is extremely slow for axions with masses in the CDM regime, but could be detectable for eV masses. The signature would be a quasi-monochromatic emission line from galaxies and galaxy clusters. The expected optical line intensity for DFSZ axions is similar to the continuum night emission. A search for optical line emission in two Abell clusters using spectra from the VIMOS (Visible Multi-Object Spectrograph) integral field unit at the Very Large Telescope (VLT) excludes axions and ALPs with a two photon coupling bigger than $\sim 5\times10^{-12}\,\mathrm{GeV}^{-1}$ in the mass range between 4.5 and 5.5 eV [164], see Fig. 89.1. Spectral data on the dwarf spheroidal galaxy Leo T from the Multi Unit Spectroscopic Explorer (MUSE) at the VLT improve these constraints by more than an order of magnitude for ALP masses between 2.7 and 5.3 eV [163], see Fig. 89.1. More recently, as shown in Fig. 89.1, James Webb Space Telescope (JWST) data has been used to constrain axion decays in the mass range between roughly 0.1 - 4 eV [492,493].

Very low-mass axions in halos produce a weak quasi-monochromatic radio line. Virial velocities in undisrupted dwarf galaxies are very low, and the axion decay line would therefore be extremely narrow. A search with the Haystack radio telescope on three nearby dwarf galaxies gave a limit $|g_{a\gamma\gamma}| < 1.0\times10^{-9}$ GeV^{-1} at 96% CL for $298 < m_a < 363$ μeV [162]. However, this combination of m_a and $g_{a\gamma\gamma}$ does not exclude plausible axion models.

A monochromatic signal is also produced in the conversion of DM axions in the background of slowly varying galactic B-fields [494]. The signal is, however, sensitive to magnetic field power on the scale of the axion mass [495]. Present and future radio telescopes appear to be able to probe ALP DM in the mass range $0.1-100$ μeV for couplings $g_{a\gamma\gamma} \gtrsim 10^{-13}$ GeV^{-1} [495] – several orders of magnitude above the KSVZ benchmark coupling for this mass range.

Resonant conversion of QCD axion or ALP DM in NS magnetospheres may give a detectable signal from individual NSs and populations of NSs for axion masses in the μeV range [496,497]. Recent analyses of radio data in the frequency range 1-40 GHz from several NSs have found no evidence for a narrow peak predicted from this mechanism and therefore put bounds on the axion-photon coupling around $10^{-11}\,\mathrm{GeV}^{-2}$ for $m_a \sim 10$ μeV [134–137], see Fig. 89.1. However, these limits may still suffer from uncertainties because the line intensity is difficult to predict in detail in the complicated environments of NSs [136,498–501]. Still, next-generation radio telescopes such as the Square Kilometer Array (SKA) may reach sensitivity to QCD axion-strength couplings [137]. Stimulated ALP decays in high radiation environments may be detectable, by *e.g.* the future SKA, down to $g_{a\gamma\gamma} \gtrsim 10^{-11}\,\mathrm{GeV}^{-1}$, for masses between μeV and 0.1 meV [502]. Furthermore, in condensed ALP dark matter structures such as ALP stars, a parametric instability may lead to an exponential enhancement of the photon flux from ALP-photon conversion by a factor $\sim \exp\left[|g_{a\gamma\gamma}|\int ds\rho_a^{1/2}(s)\right]$, where the integral is along any photon trajectory through the ALP overdensity [503,504].

Photon propagation on an ULA DM background can induce birefringence that can be compared with upper limits from the CMB [107,142,505] and may also be probed with other sources such as pulsars [506].

89.5.4 *Microwave cavity experiments*

Galactic halo axions in the 1-100 μeV mass range may be detected by their resonant conversion into a quasi-monochromatic microwave signal in a high-Q electromagnetic cavity permeated by a strong, static B field [7,507,508]. The signal should appear as excess electromagnetic power in the cavity at a frequency of the total axion energy, rest mass plus kinetic energy, $\nu = (m_a/2\pi)\left[1+\mathcal{O}(10^{-6})\right]$, with a width representing the axions' virial distribution in the galaxy near Earth (though the frequency spectrum may also contain narrower components from *e.g.* DM streams or other substructure; cavity experiments have enhanced sensitivity to these narrow signals [509]). If the cavity resonant frequency is tuned to the axion frequency then the signal is enhanced by the Q of the cavity. Cavity experiments are thus tuning experiments, typically right circular cylindrical cavities in solenoidal magnets tuned through mechanical motion of conductive or dielectric rods in the hopes of aligning the resonant frequency with the unknown axion mass.

A number of cavity experiments have been executed, and have historically most successful for 1-100 μeV axion masses, where the length scale of the detected photons are at laboratory length scales. At present, the ADMX [69, 70, 510, 511] and CAPP [170, 172] cavity experiments have explored CDM axions with sensitivity to the DFSZ benchmark model over much of the 2.7-4.91 μeV mass range, with the expectation of continuing to scan upwards. A number of other cavity experiments with sensitivity at or near KSVZ coupling have been operated over narrow mass ranges, notably HAYSTAC [175], TASEH [94], QUAX aγ [90], and a higher-order mode of a DMAG (formerly CAPP) cavity [512]. Exclusions are shown in Fig. 89.5.

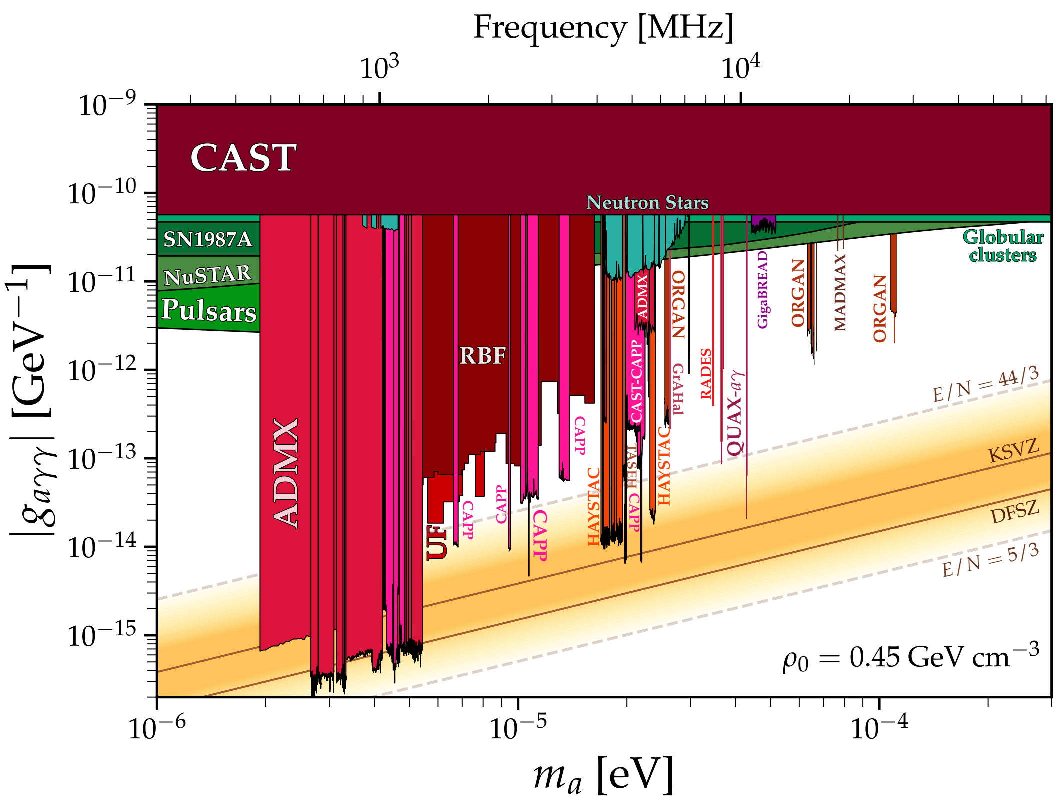

Figure 89.5: Exclusion plot for ALP-photon coupling with closeup on the parameter space of RF cavity experiments as described in the text. This is a zoom-in of Fig. 89.1 that highlights regions probed by axion DM direct detection experiments. In particular, note that all of the terrestrial experiments shaded in red, except that of CAST, assume that the axion is Figure from [63], including data from Refs. [34, 66, 68–70, 79, 82–85, 87, 88, 92, 94, 100, 123, 124, 146, 170, 170–180, 229, 513–518]

89.5.5 *Lower Mass Direct Axion Searches*

The mass range $m_a < 0.1\,\mu eV$ is increasingly of interest because of the strong theoretical motivations to prefer f_a near the GUT scale (see Sec. 89.2.1). As summarized in Fig. 89.1, there are currently no bounds in this region that reach the QCD axion parameter space. Compared to the cavity searches, the wavelength of these axions is much longer than the size of the experiments. When converted to electromagnetic fields in a strong laboratory magnetic field, the resulting signal can be modeled as currents in an equivalent "lumped" element or LC circuit [519–521]. A number of small-scale pathfinder experiments have implemented this approach focusing on the ∼neV region. Some examples are: ABRACADABRA-10cm [64, 65] used a broadband configuration in a toroidal magnet for a broadband search, ADMX-SLIC [72] used a resonant configuration in a solenoid magnet for a narrowband search, BASE [81] used the field of a penning trap for a similar search, and SHAFT [93] used a broadband configuration in a toroidal magnet enhanced with a ferromagnetic core. Future work in this direction is expected to have much more competitive sensitivity [522, 523]. There are several efforts using laser interferometry techniques to access the ∼ neV parameter space: LIDA [86] and DANCE [524]. The ADBC experiment uses an optical bow-tie cavity to search for axion-induced birefringence [67].

89.5.6 *Higher Mass Direct Axion Searches*

For axion masses where the Compton wavelength is much smaller than the characteristic laboratory scale, the conversion volume becomes smaller, the characteristic resonator Q decreases, and the standard quantum noise limit increases. Techniques beyond the standard microwave cavity experiments are necessary to probe QCD axion dark matter. No experiment yet has explored a wide range of masses for high masses, beyond ∼25 μeV, at QCD axions couplings, but a number of promising techniques have been proposed and demonstrated. Only a small representative subset of these demonstrations and proposals can be cataloged here.

One path is to build a resonant structure with a characteristic dimension much longer than the converted photon wavelength. This has been proposed with subwavelength metamaterial-like conductive structure such as ALPHA [525], or VERA [526], and demonstrated with periodic dielectrics as in Orpheus [527], LAMPOST [528], and MADMAX [176]. Another option is to forgo resonance entirely and use a single conversion surface that is many wavelengths in dimension. [529] Experiments working on this concept include BRASS [530] and BREAD [174].

Due to the standard quantum limit readout for these higher-mass detectors to transition to quantum sensors and techniques beyond linear amplifiers [526, 531]. Benefits of quantum squeezing have been demonstrated in HAYSTAC [84, 532], and single photon counters have been shown to be beneficial in QUAX [533] and in dark photon searches using both qubit devices [534] and nanowire detectors [528], among others.

89.5.7 *Non-photon Coupling Axion Searches*

In Fig. 89.2 and Fig. 89.3, the status of the non-photon couplings are discussed in the context of the astrophysical ALP searches. These same non-photon couplings can also be used in direct searches for DM. Much of the recent focus is on the induced oscillating nuclear electric dipole moments (EDMs) [26]. An analysis of the ratio of spin-precession frequencies of stored ultracold neutrons and ^{199}Hg atoms measured by neutron EDM experiments excludes a sizeable region of parameter space in the mass region $10^{-24}\,\text{eV} \leq m_a \leq 10^{-17}\,\text{eV}$ [31], which surpass the limits on anomalous energy loss of SN 1987A by more than seven orders of magnitude and are competitive with the ones from the requirement of successful BBN established in [29].

The oscillating EDMs can also be studied with the precession of nuclear spins in a nucleon spin polarized sample in the presence of an electric field. The resulting transverse magnetization can be searched for by exploiting magnetic-resonance (MR) techniques, which are most sensitive to the low frequencies of the sub-neV axion masses. There are two interactions probed through MR: the electric dipole coupling as searched for by CASPEr-electric (with first result presented in [30]) and an interaction between the axion field gradient and the nuclear spin, which is aimed for by CASPEr-Gradient. The latter has been explored through comagnetometry between a variety of nuclei: H-^{12}C in CASPEr-ZULF [261, 262], Xe-Rb in NASDUCK [535], K-^{3}He [265–267] and K-^{3}He-Rb [273]. The exclusions from these searches are shown in Fig. 89.3.

Sub-μeV ALP masses can also be probed by using the storage ring EDM method proposed in Ref. [536] which exploits a combination of B and E-fields to produce a resonance between the $g-2$ spin precession frequency and the DM ALP field oscillation frequency. This method, however, does not reach the sensitivity to probe the QCD axion prediction for $g_{aN\gamma}$. An non-zero axion electron coupling g_{aee} will lead to an electron spin precession about the axion DM wind [537, 538]. The QUAX a-e experiment exploits MR inside a magnetized material [539]. Because of the higher Larmor frequency of the electron, it is sensitive to masses above 1μeV. Its first run established the bound in the mass range 42.4 to 43.2 μeV [240], slightly better than the solar ν bound, shown in Fig. 89.2. Updated sensitivity projections for MR techniques are detailed in [540], and examples of potential technological improvements in [541, 542].

89.6 Conclusions

There is a strengthening physics case for very weakly coupled light particles beyond the Standard Model. The elegant solution of the strong CP problem proposed by Peccei and Quinn yields a particularly strong motivation for the axion. In many theoretically appealing ultraviolet completions of the Standard Model axions and ALPs occur automatically. Moreover, they are natural CDM candidates. Interestingly, a significant portion of previously unexplored, but phenomenologically interesting and theoretically very well motivated, axion and ALP parameter space can be explored in the foreseeable future by a number of terrestrial experiments and astrophysical observations.

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