Axions – THIS IS PART 1 OF 2

To reduce the size of this section's PostScript file, we have divided it into two PostScript files. We present the following index:

PART 1

PART 2

Axions (A^0) and Other Very Light Bosons, Searches for

AXIONS AND OTHER VERY LIGHT BOSONS

Written October 1997 by H. Murayama (University of California, Berkeley) Part I; April 1998 by G. Raffelt (Max-Planck Institute, München) Part II; and April 1998 by C. Hagmann, K. van Bibber (Lawrence Livermore National Laboratory), and L.J. Rosenberg (Massachusetts Institute of Technology) Part III.

This review is divided into three parts:

Part I (Theory)
Part II (Astrophysical Constraints)
Part III (Experimental Limits)

AXIONS AND OTHER VERY LIGHT BOSONS, PART I (THEORY)

(by H. Murayama)

In this section we list limits for very light neutral (pseudo) scalar bosons that couple weakly to stable matter. They arise if there is a global continuous symmetry in the theory that is spontaneously broken in the vacuum. If the symmetry is exact, it results in a massless Nambu–Goldstone (NG) boson. If there is a small explicit breaking of the symmetry, either already in the Lagrangian or due to quantum mechanical effects such as anomalies, the would-be NG boson acquires a finite mass; then it is called a pseudo-NG boson. Typical examples are axions (A^0) [1], familons [2], and Majorons [3,4], associated, respectively, with spontaneously broken Peccei-Quinn [5], family, and lepton-number symmetries. This Review provides brief descriptions of each of them and their motivations.

One common characteristic for all these particles is that their coupling to the Standard Model particles are suppressed by the energy scale of symmetry breaking, *i.e.* the decay constant f, where the interaction is described by the Lagrangian

$$\mathcal{L} = \frac{1}{f}(\partial_\mu \phi) J^\mu, \tag{1}$$

where J^μ is the Noether current of the spontaneously broken global symmetry.

An axion gives a natural solution to the strong CP problem: why the effective θ-parameter in the QCD Lagrangian $\mathcal{L}_\theta = \theta_{\mathit{eff}} \frac{\alpha_s}{8\pi} F^{\mu\nu a} \widetilde{F}^a_{\mu\nu}$ is so small ($\theta_{\mathit{eff}} \lesssim 10^{-9}$) as required by the current limits on the neutron electric dipole moment, even though $\theta_{\mathit{eff}} \sim O(1)$ is perfectly allowed by the QCD gauge invariance. Here, θ_{eff} is the effective θ parameter after the diagonalization of the quark masses, and $F^{\mu\nu a}$ is the gluon field strength and $\widetilde{F}^a_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}F^{\rho\sigma a}$. An axion is a pseudo-NG boson of a spontaneously broken Peccei–Quinn symmetry, which is an exact symmetry at the classical level, but is broken quantum mechanically due to the triangle anomaly with the gluons. The definition of the Peccei–Quinn symmetry is model dependent. As a result of the triangle anomaly, the axion acquires an effective coupling to gluons

$$\mathcal{L} = \left(\theta_{\mathit{eff}} - \frac{\phi_A}{f_A}\right) \frac{\alpha_s}{8\pi}\, F^{\mu\nu a}\, \widetilde{F}^a_{\mu\nu}\ , \tag{2}$$

where ϕ_A is the axion field. It is often convenient to *define* the axion decay constant f_A with this Lagrangian [6]. The QCD nonperturbative effect induces a potential for ϕ_A whose minimum is at $\phi_A = \theta_{\mathit{eff}} f_A$ cancelling θ_{eff} and solving the strong CP problem. The mass of the axion is inversely proportional to f_A as

$$m_A = 0.62 \times 10^{-3}\mathrm{eV} \ \times (10^{10}\mathrm{GeV}/f_A)\ . \tag{3}$$

The original axion model [1,5] assumes $f_A \sim v$, where $v = (\sqrt{2}G_F)^{-1/2} = 247$ GeV is the scale of the electroweak symmetry breaking, and has two Higgs doublets as minimal ingredients. By requiring tree-level flavor conservation, the axion mass and its couplings are completely fixed in terms of one parameter ($\tan\beta$): the ratio of the vacuum expectation values of two Higgs fields. This model is excluded after extensive experimental searches for such an axion [7]. Observation of a narrow-peak structure in positron spectra from heavy ion collisions [8] suggested a particle of mass 1.8 MeV that decays into e^+e^-. Variants of the original axion model, which keep $f_A \sim v$, but drop the constraints of tree-level flavor conservation, were proposed [9]. Extensive searches for this particle, A^0(1.8 MeV), ended up with another negative result [10].

The popular way to save the Peccei-Quinn idea is to introduce a new scale $f_A \gg v$. Then the A^0 coupling becomes weaker, thus one can easily avoid all the existing experimental limits; such models are called invisible axion models [11,12]. Two classes of models are discussed commonly in the literature. One introduces new heavy quarks which carry Peccei–Quinn charge while the usual quarks and leptons do not (KSVZ axion or "hadronic axion") [11]. The other does not need additional quarks but requires two Higgs doublets, and all quarks and leptons carry Peccei–Quinn charges (DFSZ axion or "GUT-axion") [12]. All models contain at least one electroweak singlet scalar boson which acquires an expectation value and breaks Peccei–Quinn symmetry. The invisible axion with a large decay constant $f_A \sim 10^{12}$ GeV was found to be a good candidate of the cold dark matter component of the Universe [13](see Dark Matter review). The energy density is stored in the low-momentum modes of the axion field which are highly occupied and thus represent essentially classical field oscillations.

The constraints on the invisible axion from astrophysics are derived from interactions of the axion with either photons, electrons or nucleons. The strengths of the interactions are model dependent (*i.e.*, not a function of f_A only), and hence one needs to specify a model in order to place lower bounds on f_A. Such constraints will be discussed in Part II. Serious experimental searches for an invisible axion are underway; they typically rely on axion-photon coupling, and some of them assume that the axion is the dominant component of our galactic halo density. Part III will discuss experimental techniques and limits.

Familons arise when there is a global family symmetry broken spontaneously. A family symmetry interchanges generations or acts on different generations differently. Such a symmetry may explain the structure of quark and lepton masses and their mixings. A familon could be either a scalar or a pseudoscalar. For instance, an SU(3) family symmetry among three generations is non-anomalous and hence the familons are exactly massless. In this case, familons are scalars. If one has larger family symmetries with separate groups of left-handed and right-handed fields, one also has pseudoscalar familons. Some of them have flavor-off-diagonal couplings such as $\partial_\mu \phi_F \bar{d} \gamma^\mu s / F_{ds}$ or $\partial_\mu \phi_F \bar{e} \gamma^\mu \mu / F_{\mu e}$, and the decay constant F can be different for individual operators. The decay constants have lower bounds constrained by flavor-changing processes. For instance, $\mathrm{B}(K^+ \to \pi^+\phi_F) < 3 \times 10^{-10}$ [14] gives $F_{ds} > 3.4 \times 10^{11}$ GeV [15]. The constraints on familons primarily coupled to third generation are quite weak [15].

If there is a global lepton-number symmetry and if it breaks spontaneously, there is a Majoron. The triplet Majoron model [4] has a weak-triplet Higgs boson, and Majoron couples to Z. It is now excluded by the Z invisible-decay width. The model is viable if there is an additional singlet Higgs boson and if the Majoron is mainly a singlet [16]. In the singlet Majoron model [3], lepton-number symmetry is broken by a weak-singlet scalar field, and there are right-handed neutrinos which acquire Majorana masses. The left-handed neutrino masses are generated by a "seesaw" mechanism [17]. The scale of lepton number breaking can be much higher than the electroweak scale in this case. Astrophysical constraints require the decay constant to be $\gtrsim 10^9$ GeV [18].

There is revived interest in a long-lived neutrino, to improve Big-Bang Nucleosynthesis [19] or large scale structure formation theories [20]. Since a decay of neutrinos into electrons or photons is severely constrained, these scenarios require a familon (Majoron) mode $\nu_1 \to \nu_2\phi_F$ (see, *e.g.*, Ref. 15 and references therein).

Other light bosons (scalar, pseudoscalar, or vector) are constrained by "fifth force" experiments. For a compilation of constraints, see Ref. 21.

It has been widely argued that a fundamental theory will not possess global symmetries; gravity, for example, is expected to violate them. Global symmetries such as baryon number arise by accident, typically as a consequence of gauge symmetries. It has been noted [22] that the Peccei-Quinn symmetry, from this perspective, must also arise by accident and must hold to an extraordinary degree of accuracy in order to solve the strong CP problem. Possible resolutions to this problem, however, have been discussed [22,23]. String theory also provides sufficiently good symmetries, especially using a large compactification radius motivated by recent developments in M-theory [24].

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AXIONS AND OTHER VERY LIGHT BOSONS: PART II (ASTROPHYSICAL CONSTRAINTS)

(by G.G. Raffelt)

Low-mass weakly-interacting particles (neutrinos, gravitons, axions, baryonic or leptonic gauge bosons, *etc.*) are produced in hot plasmas and thus represent an energy-loss channel for stars. The strength of the interaction with photons, electrons, and nucleons can be constrained from the requirement that stellar-evolution time scales are not modified beyond observational limits. For detailed reviews see Refs. [1,2].

The energy-loss rates are steeply increasing functions of temperature T and density ρ. Because the new channel has to compete with the standard neutrino losses which tend to increase even faster, the best limits arise from low-mass stars, notably from horizontal-branch (HB) stars which have a helium-burning core of about 0.5 solar masses at $\langle\rho\rangle \approx 0.6\times10^4\,\mathrm{g\,cm^{-3}}$ and $\langle T\rangle \approx 0.7\times10^8\,\mathrm{K}$. The new energy-loss rate must not exceed about $10\,\mathrm{ergs\,g^{-1}\,s^{-1}}$ to avoid a conflict with the observed number ratio of HB stars in globular clusters. Likewise the ignition of helium in the degenerate cores of the preceding red-giant phase is delayed too much unless the same constraint holds at $\langle\rho\rangle \approx 2\times10^5\,\mathrm{g\,cm^{-3}}$ and $\langle T\rangle \approx 1\times10^8\,\mathrm{K}$. The white-dwarf luminosity function also yields useful bounds.

The new bosons X^0 interact with electrons and nucleons with a dimensionless strength g. For scalars it is a Yukawa coupling, for new gauge bosons (*e.g.*, from a baryonic or leptonic gauge symmetry) a gauge coupling. Axion-like pseudoscalars couple derivatively as $f^{-1}\bar\psi\gamma_\mu\gamma_5\psi\,\partial^\mu\phi_X$ with f an energy scale. Usually this is equivalent to $(2m/f)\bar\psi\gamma_5\psi\,\phi_X$ with m the mass of the fermion ψ so that $g=2m/f$. For the coupling to electrons, globular-cluster stars yield the constraint

$$g_{Xe} \lesssim \begin{cases} 0.5\times10^{-12} & \text{for pseudoscalars [3]} \ , \\ 1.3\times10^{-14} & \text{for scalars [4]} \ , \end{cases} \tag{1}$$

if $m_X \lesssim 10\,\mathrm{keV}$. The Compton process $\gamma + {}^4\mathrm{He} \to {}^4\mathrm{He} + X^0$ limits the coupling to nucleons to $g_{XN} \lesssim 0.4\times10^{-10}$ [4].

Scalar and vector bosons mediate long-range forces which are severely constrained by "fifth-force" experiments [5]. In the massless case the best limits come from tests of the equivalence principle in the solar system, leading to

$$g_{B,L} \lesssim 10^{-23} \tag{2}$$

for a baryonic or leptonic gauge coupling [6].

In analogy to neutral pions, axions A^0 couple to photons as $g_{A\gamma}\mathbf{E}\cdot\mathbf{B}\,\phi_A$ which allows for the Primakoff conversion $\gamma \leftrightarrow A^0$ in external electromagnetic fields. The most restrictive limit arises from globular-cluster stars [2]

$$g_{A\gamma} \lesssim 0.6\times10^{-10}\,\mathrm{GeV^{-1}} \ . \tag{3}$$

The often-quoted "red-giant limit" [7] is slightly weaker.

The duration of the SN 1987A neutrino signal of a few seconds proves that the newborn neutron star cooled mostly by neutrinos rather than through an "invisible channel" such as right-handed (sterile) neutrinos or axions [8]. Therefore,

$$3 \times 10^{-10} \lesssim g_{AN} \lesssim 3 \times 10^{-7} \tag{4}$$

is excluded for the pseudoscalar Yukawa coupling to nucleons [2]. The "strong" coupling side is allowed because axions then escape only by diffusion, quenching their efficiency as an energy-loss channel [9]. Even then the range

$$10^{-6} \lesssim g_{AN} \lesssim 10^{-3} \tag{5}$$

is excluded to avoid excess counts in the water Cherenkov detectors which registered the SN 1987A neutrino signal [11].

In terms of the Peccei-Quinn scale f_A, the axion couplings to nucleons and photons are $g_{AN} = C_N m_N / f_A$ ($N = n$ or p) and $g_{A\gamma} = (\alpha/2\pi f_A)\,(E/N - 1.92)$ where C_N and E/N are model-dependent numerical parameters of order unity. With $m_A = 0.62\,\mathrm{eV}\,(10^7\,\mathrm{GeV}/f_A)$, Eq. (3) yields $m_A \lesssim 0.4\,\mathrm{eV}$ for $E/N = 8/3$ as in GUT models or the DFSZ model. The SN 1987A limit is $m_A \lesssim 0.008\,\mathrm{eV}$ for KSVZ axions while it varies between about 0.004 and 0.012 eV for DFSZ axions, depending on the angle β which measures the ratio of two Higgs vacuum expectation values [10]. In view of the large uncertainties it is good enough to remember $m_A \lesssim 0.01\,\mathrm{eV}$ as a generic limit (Fig. 1).

In the early universe, axions come into thermal equilibrium only if $f_A \lesssim 10^8\,\mathrm{GeV}$ [12]. Some fraction of the relic axions end up in galaxies and galaxy clusters. Their decay $a \to 2\gamma$ contributes to the cosmic extragalactic background light and to line emissions from galactic dark-matter haloes and galaxy clusters. An unsuccessful "telescope search" for such features yields $m_a < 3.5\,\mathrm{eV}$ [13]. For $m_a \gtrsim 30\,\mathrm{eV}$, the axion lifetime is shorter than the age of the universe.

For $f_A \gtrsim 10^8\,\mathrm{GeV}$ cosmic axions are produced nonthermally. If inflation occurred after the Peccei-Quinn symmetry breaking or if $T_{\mathrm{reheat}} < f_A$, the

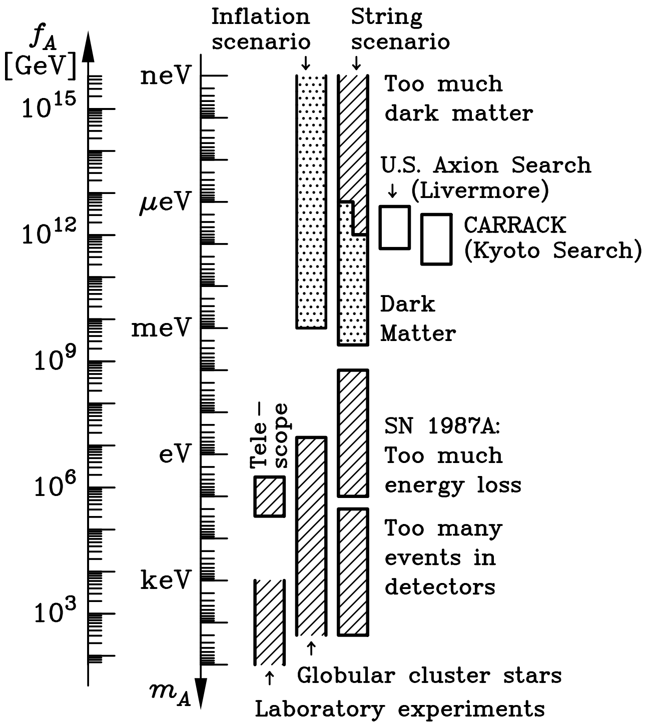

Figure 1: Astrophysical and cosmological exclusion regions (hatched) for the axion mass m_A or equivalently, the Peccei-Quinn scale f_A. An "open end" of an exclusion bar means that it represents a rough estimate; its exact location has not been established or it depends on detailed model assumptions. The globular cluster limit depends on the axion-photon coupling; it was assumed that $E/N = 8/3$ as in GUT models or the DFSZ model. The SN 1987A limits depend on the axion-nucleon couplings; the shown case corresponds to the KSVZ model and approximately to the DFSZ model. The dotted "inclusion regions" indicate where axions could plausibly be the cosmic dark matter. Most of the allowed range in the inflation scenario requires fine-tuned initial conditions. In the string scenario the plausible dark-matter range is controversial as indicated by the step in the low-mass end of the "inclusion bar" (see main text for a discussion). Also shown is the projected sensitivity range of the search experiments for galactic dark-matter axions.

"misalignment mechanism" [14] leads to a contribution to the cosmic critical density of

$$\Omega_A h^2 \approx 1.9 \times 3^{\pm 1} \, (1\,\mu\mathrm{eV}/m_A)^{1.175} \, \Theta_\mathrm{i}^2 F(\Theta_\mathrm{i}) \tag{6}$$

where h is the Hubble constant in units of $100\,\mathrm{km\,s^{-1}\,Mpc^{-1}}$. The stated range reflects recognized uncertainties of the cosmic conditions at the QCD phase transition and of the temperature-dependent axion mass. The function $F(\Theta)$ with $F(0) = 1$ and $F(\pi) = \infty$ accounts for anharmonic corrections to the axion potential. Because the initial misalignment angle Θ_i can be very small or very close to π, there is no real prediction for the mass of dark-matter axions even though one would expect $\Theta_\mathrm{i}^2 F(\Theta_\mathrm{i}) \sim 1$ to avoid fine-tuning the initial conditions.

A possible fine-tuning of Θ_i is limited by inflation-induced quantum fluctuations which in turn lead to temperature fluctuations of the cosmic microwave background [15,16]. In a broad class of inflationary models one thus finds an upper limit to m_A where axions could be the dark matter. According to the most recent discussion [16] it is about 10^{-3} eV (Fig. 1).

If inflation did not occur at all or if it occurred before the Peccei-Quinn symmetry breaking with $T_\mathrm{reheat} > f_A$, cosmic axion strings form by the Kibble mechanism [17]. Their motion is damped primarily by axion emission rather than gravitational waves. After axions acquire a mass at the QCD phase transition they quickly become nonrelativistic and thus form a cold dark matter component. Battye and Shellard [18] found that the dominant source of axion radiation are string loops rather than long strings. At a cosmic time t the average loop creation size is parametrized as $\langle \ell \rangle = \alpha t$ while the radiation power is $P = \kappa\mu$ with μ the renormalized string tension. The loop contribution to the cosmic axion density is [18]

$$\Omega_A h^2 \approx 88 \times 3^{\pm 1} \left[(1 + \alpha/\kappa)^{3/2} - 1 \right] \, (1\,\mu\mathrm{eV}/m_A)^{1.175} \, , \tag{7}$$

where the stated nominal uncertainty has the same source as in Eq. (6). The values of α and κ are not known, but probably $0.1 < \alpha/\kappa < 1.0$ [18], taking the expression in square brackets to 0.15–1.83. If axions are the dark matter, we have

$$0.05 \lesssim \Omega_A h^2 \lesssim 0.50 \, , \tag{8}$$

where it was assumed that the universe is older than 10 Gyr, that the dark-matter density is dominated by axions with $\Omega_A \gtrsim 0.2$, and that $h \gtrsim 0.5$. This implies $m_A = 6$–2500 μeV for the plausible mass range of dark-matter axions (Fig. 1).

Contrary to Ref. 18, Sikivie *et al.* [19] find that the motion of global strings is strongly damped, leading to a flat axion spectrum. In Battye and Shellard's treatment the axion radiation is strongly peaked at wavelengths of order the loop size. In Sikivie *et al.*'s picture more of the string radiation goes into kinetic axion energy which is redshifted so that ultimately there are fewer axions. In this scenario the contributions from string decay and vacuum realignment are of the same order of magnitude; they are both given by Eq. (6) with Θ_i of order one. As a consequence, Sikivie *et al.* allow for a plausible range of dark-matter axions which reaches to smaller masses as indicated in Fig. 1.

The work of both groups implies that the low-mass end of the plausible mass interval in the string scenario overlaps with the projected sensitivity range of the U.S. search experiment for galactic dark-matter axions (Livermore) [20] and of the Kyoto search experiment CARRACK [21] as indicated in Fig. 1. (See also Part III of this Review by Hagmann, van Bibber, and Rosenberg.)

In summary, a variety of robust astrophysical arguments and laboratory experiments (Fig. 1) indicate that $m_A \lesssim 10^{-2}$ eV. The exact value of this limit may change with a more sophisticated treatment of supernova physics and/or the observation of the neutrino signal from a future galactic supernova, but a dramatic modification is not expected unless someone puts forth a completely new argument. The stellar-evolution limits shown in Fig. 1 depend on the axion couplings to various particles and thus can be irrelevant in fine-tuned models where, for example, the axion-photon coupling strictly vanishes. For nearly any m_A in the range generically allowed by stellar evolution, axions could be the cosmic dark matter, depending on the cosmological scenario realized in nature. It appears that our only practical chance to discover these “invisible” particles rests with the ongoing or future search experiments for galactic dark-matter.

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AXIONS AND OTHER VERY LIGHT BOSONS, PART III (EXPERIMENTAL LIMITS)

(by C. Hagmann, K. van Bibber, and L.J. Rosenberg)

In this section we review the experimental methodology and limits on light axions and light pseudoscalars in general. (A comprehensive overview of axion theory is given by H. Murayama in the Part I of this Review, whose notation we follow [1].) Within its scope are searches where the axion is assumed to be dark matter, searches where the Sun is presumed to be a source of axions, and purely laboratory experiments. We restrict the discussion to axions of mass m_A $< O(\text{eV})$, as the allowed range for the axion mass is nominally $10^{-6} < m_A < 10^{-2}$ eV. Experimental work in this range predominantly has been through the axion-photon coupling $g_{A\gamma}$, to which the present review is confined. As discussed in Part II of this Review by G. Raffelt, the lower bound derives from a cosmological overclosure argument, and the upper bound from SN1987A [2]. Limits from stellar evolution overlap seamlessly above that, connecting with accelerator-based limits which ruled out the original axion. There it was assumed that the Peccei-Quinn symmetry-breaking scale was the electroweak scale, *i.e.*, $f_A \sim 250$ GeV, implying axions of mass $m_A \sim O(100\,\text{keV})$. These earlier limits from nuclear transitions, particle decays, *etc.*, while not discussed here, are included in the Listings.

While the axion mass is well determined by the Peccei-Quinn scale, *i.e.*, $m_A = 0.62$ eV $(10^7\,\mathrm{GeV}/f_A)$, the axion-photon coupling $g_{A\gamma}$ is not: $g_{A\gamma} = (\alpha/\pi f_A)\,g_\gamma$, with $g_\gamma = (E/N - 1.92)/2$, where E/N is a model-dependent number. It is noteworthy however, that two quite distinct models lead to axion-photon couplings which are not very different. For the case of axions imbedded in Grand Unified Theories, the DFSZ axion [3], $g_\gamma = 0.37$, whereas in one popular implementation of the "hadronic" class of axions, the KSVZ axion [4], $g_\gamma = -0.96$. The Lagrangian $L = g_{A\gamma}\,\mathbf{E}\cdot\mathbf{B}\,\phi_A$, with ϕ_A the axion field, permits the conversion of an axion into a single real photon in an external electromagnetic field, *i.e.*, a Primakoff interaction. In the case of relativistic axions, $k_\gamma - k_A \sim m_A^2/2\omega \ll \omega$, pertinent to several experiments below, coherent axion-photon mixing in long magnetic fields results in significant conversion probability even for very weakly coupled axions [5].

Below are discussed several experimental techniques constraining $g_{A\gamma}$, and their results. Also included are recent but yet-unpublished results, and projected sensitivities for experiments soon to be upgraded.

III.1. Microwave cavity experiments: Possibly the most promising avenue to the discovery of the axion presumes that axions constitute a significant fraction of the dark matter halo of our galaxy. The maximum likelihood density for the Cold Dark Matter (CDM) component of our galactic halo is $\rho_{\mathrm{CDM}} = 7.5\times10^{-25}\mathrm{g/cm}^3(450\,\mathrm{MeV/cm}^3)$ [6]. That the CDM halo is in fact made of axions (rather than *e.g.* WIMPs) is in principle an independent assumption, however should very light axions exist they would almost necessarily be cosmologically abundant [2]. As shown by Sikivie [7], halo axions may be detected by their resonant conversion into a quasi-monochromatic microwave signal in a high-Q cavity permeated by a strong magnetic field. The cavity is tunable and the signal is maximum when the frequency $\nu = m_A(1 + O(10^{-6}))$, the width of the peak representing the virial distribution of thermalized axions in the galactic gravitational potential. The signal may possess ultra-fine structure due to axions recently fallen into the galaxy and not yet thermalized [8]. The feasibility of the technique was established in early experiments of small sensitive volume,

$V = O(1\,\text{liter})$ [9,10] with High Electron Mobility Transistor (HEMT) amplifiers, which set limits on axions in the mass range $4.5 < m_A < 16.3\,\mu\text{eV}$, but at power sensitivity levels 2–3 orders of magnitude too high to see KSVZ and DFSZ axions (the conversion power $P_{A\to\gamma} \propto g^2_{A\gamma}$). A recent large-scale experiment ($B \sim 7.5\,\text{T}, V \sim 200\,\text{liter}$) has achieved sensitivity to KSVZ axions over a narrow mass range $2.77 < m_A < 3.3\,\mu\text{eV}$, and continues to take data [11]. The exclusion regions shown in Fig. 1 for Refs. [9–12] are all normalized to the best-fit Cold Dark Matter density $\rho_{\text{CDM}} = 7.5 \times 10^{-25}\text{g/cm}^3(450\,\text{MeV/cm}^3)$, and 90% CL. Recent developments in DC SQUID amplifiers [12] and Rydberg atom single-quantum detectors [13] promise dramatic improvements in noise temperature, which will enable rapid scanning of the axion mass range at or below the DFSZ limit. The region of the microwave cavity experiments is shown in detail in Fig. 2.

III.2. Telescope search for eV axions: For axions of mass greater than about 10^{-1} eV, their cosmological abundance is no longer dominated by vacuum misalignment or string radiation mechanisms, but rather by thermal production. Their contribution to the critical density is small, $\Omega \sim 0.01\,(m_A/\text{eV})$. However, the spontaneous-decay lifetime of axions, $\tau(A \to 2\gamma) \sim 10^{25}\text{sec}(m_A/\text{eV})^{-5}$ while irrelevant for μeV axions, is short enough to afford a powerful constraint on such thermally produced axions in the eV range, by looking for a quasi-monochromatic photon line from galactic clusters. This line, corrected for Doppler shift, would be at half the axion mass and its width would be consistent with the observed virial motion, typically $\Delta\lambda/\lambda \sim 10^{-2}$. The expected line intensity would be of the order $I_A \sim 10^{-17}(m_A/3\,\text{eV})^7\text{erg}\,\text{cm}^{-2}\text{arcsec}^{-2}\text{Å}^{-1}\text{sec}^{-1}$ for DFSZ axions, comparable to the continuum night emission. The conservative assumption is made that the relative density of thermal axions fallen into the cluster gravitational potential reflects their overall cosmological abundance. A search for thermal axions in three rich Abell clusters was carried out at Kitt Peak National Laboratory [14]; no such line was observed between 3100–8300 Å (m_A = 3–8 eV) after "on-off field" subtraction of the atmospheric molecular background spectra. A limit everywhere stronger than $g_{A\gamma} < 10^{-10}\text{GeV}^{-1}$ is set, which is seen from Fig. 1 to easily exclude DFSZ axions throughout the mass range.

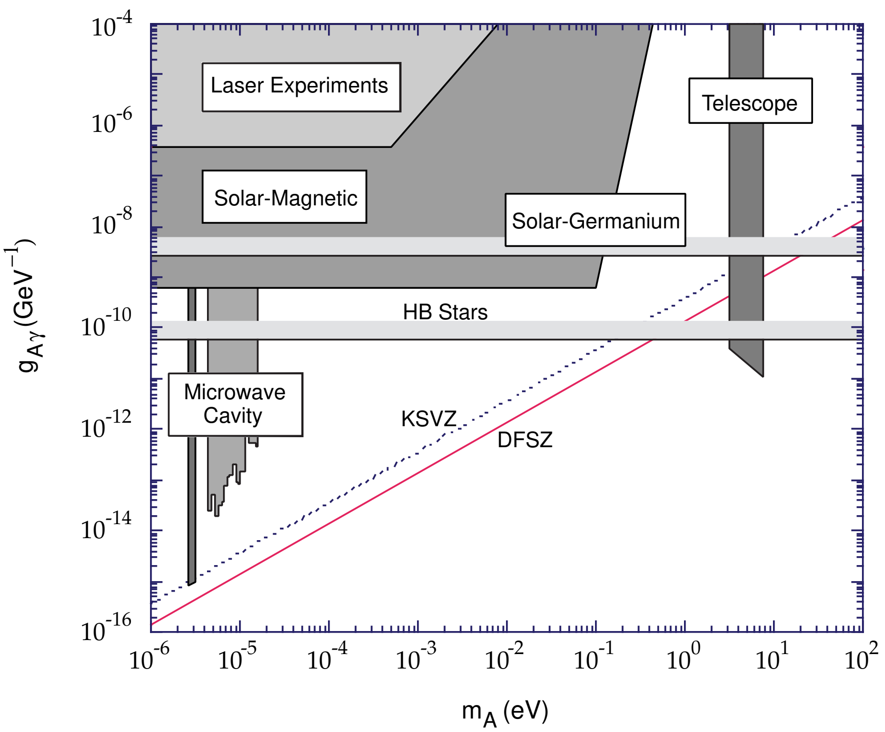

Figure 1: Exclusion region in mass vs. axion-photon coupling $(m_A, g_{A\gamma})$ for various experiments. The limit set by globular cluster Horizontal Branch Stars ("HB Stars") is shown for Ref. 2.

III.3. A search for solar axions: As with the telescope search for thermally produced axions above, the search for solar axions was stimulated by the possibility of there being a "1 eV window" for hadronic axions (*i.e.*, axions with no tree-level coupling to leptons), a "window" subsequently closed by an improved understanding of the evolution of globular cluster stars and SN1987A [2]. Hadronic axions would be copiously produced within our Sun's interior by a Primakoff process. Their flux at the Earth of $\sim 10^{12}\mathrm{cm}^{-2}\mathrm{sec}^{-1}(m_A/\mathrm{eV})^2$, which is independent of the details of the solar model, is sufficient for a definitive

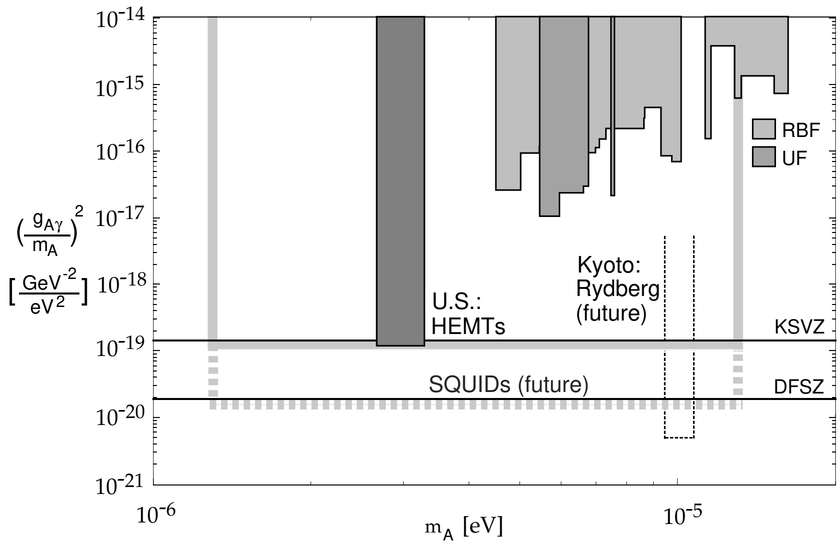

Figure 2: Exclusion region from the microwave cavity experiments, where the plot is flattened by presenting $(g_{A\gamma}/m_A)^2$ vs. m_A. The first-generation experiments (Rochester-BNL-FNAL, "RBF" [9]; University of Florida, "UF" [10]) and the US large-scale experiment in progress ("US" [11]) are all HEMT-based. Shown also is the full mass range to be covered by the latter experiment (shaded line), and the improved sensitivity when upgraded with DC SQUID amplifiers [12] (shaded dashed line). The expected performance of the Kyoto experiment based on a Rydberg atom single-quantum receiver (dotted line) is also shown [13].

test via the axion reconversion to photons in a large magnetic field. However, their average energy is $\sim$ 4 keV, implying an oscillation length in the vacuum of $2\pi(m_A^2/2\omega)^{-1} \sim O(\text{mm})$, precluding the mixing from achieving its theoretically maximum value in any practical magnet. It was recognized that one could endow the photon with an effective mass in a gas, $m_\gamma = \omega_{\text{pl}}$, thus permitting the axion and photon dispersion relationships to be matched [15]. A first simple implementation of this proposal was carried out using a conventional dipole magnet with a conversion volume of variable-pressure helium gas and a xenon proportional chamber as the x-ray detector [16]. The magnet was fixed in orientation to take data for $\sim$ 1000 sec/day. Axions were excluded for $g_{A\gamma} < 3.6 \times 10^{-9}\text{GeV}^{-1}$ for $m_A < 0.03\,\text{eV}$, and $g_{A\gamma} < 7.7 \times 10^{-9}\text{GeV}^{-1}$ for $0.03\,\text{eV} < m_A < 0.11$ eV (95% CL). A more ambitious experiment has recently been commissioned, using a superconducting magnet on a telescope mount to track the Sun continuously. A preliminary exclusion limit of $g_{A\gamma} < 6 \times 10^{-10}\text{GeV}^{-1}$ (95% CL) has been set for $m_A < 0.03$ eV [17].

Another search for solar axions has been carried out, using a single crystal germanium detector. It exploits the coherent conversion of axions into photons when their angle of incidence satisfies a Bragg condition with a crystalline plane. Analysis of 1.94 kg-yr of data from a 1 kg germanium detector yields a bound of $g_{A\gamma} < 2.7 \times 10^{-9}\text{GeV}^{-1}$ (95% CL), independent of mass up to $m_A \sim 1$ keV [18].

III.4. Photon regeneration ("invisible light shining through walls"): Photons propagating through a transverse field (with $\mathbf{E}\|\mathbf{B}$) may convert into axions. For light axions with $m_A^2 l/2\omega \ll 2\pi$, where l is the length of the magnetic field, the axion beam produced is colinear and coherent with the photon beam, and the conversion probability Π is given by $\Pi \sim (1/4)(g_{A\gamma}Bl)^2$. An ideal implementation for this limit is a laser beam propagating down a long, superconducting dipole magnet like those for high-energy physics accelerators. If another such dipole magnet is set up in line with the first, with an optical barrier interposed between them, then photons may be regenerated from the pure axion beam in the second magnet and detected [19]. The overall probability $P(\gamma \to A \to \gamma) = \Pi^2$. Such an experiment has been carried out, utilizing two magnets of length $l = 4.4$ m and $B = 3.7$ T. Axions with mass $m_A < 10^{-3}$ eV, and $g_{A\gamma} > 6.7 \times 10^{-7}\text{GeV}^{-1}$ were excluded at 95% CL [20,21]. With sufficient effort, limits comparable to those from stellar evolution would be achievable. Due to the $g_{A\gamma}^4$ rate suppression however, it does not seem feasible to reach standard axion couplings.

III.5. Polarization experiments: The existence of axions can affect the polarization of light propagating through a transverse magnetic field in two ways [22]. First, as the $E_{\|}$ component, but not the $E_{\perp}$ component will be depleted by the production of real axions, there will be in general a small rotation of the polarization vector of linearly polarized light. This effect will be a constant for all sufficiently light m_A such that the oscillation length is much longer than the magnet ($m_A^2 l/2\omega \ll 2\pi$). For heavier axions, the effect oscillates and diminishes with increasing m_A, and vanishes for $m_A > \omega$. The second effect is birefringence of the vacuum, again because there can be a mixing of virtual axions in the $E_{\|}$ state, but not for the $E_{\perp}$ state. This will lead to light which is initially linearly polarized becoming elliptically polarized. Higher-order QED also induces vacuum birefringence, and is much stronger than the contribution due to axions. A search for both polarization-rotation and induced ellipticity has been carried out with the same magnets described in Sec. (III.4) above [21,23]. As in the case of photon regeneration, the observables are boosted linearly by the number of passes the laser beam makes in an optical cavity within the magnet. The polarization-rotation resulted in a stronger limit than that from ellipticity, $g_{A\gamma} < 3.6 \times 10^{-7}\mathrm{GeV}^{-1}$ (95% CL) for $m_A < 5 \times 10^{-4}$ eV. The limits from ellipticity are better at higher masses, as they fall off smoothly and do not terminate at m_A. There are two experiments in construction with greatly improved sensitivity which while still far from being able to detect standard axions, should measure the QED "light-by-light" contribution for the first time [24,25]. The overall envelope for limits from the laser-based experiments in Sec. (III.4) and Sec. (III.5) is shown schematically in Fig. 1.

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