THE SEARCH FOR DARK MATTER AXIONS

P. SIKIVIE

Theoretical Physics Division, CERN, CH-1211 Genève 23, Switzerland
and
Department of Physics, University of Florida, Gainesville, FL 32611, USA

Axions solve the Strong CP Problem and are a cold dark matter candidate. The combined constraints from accelerator searches, stellar evolution limits and cosmology suggest that the axion mass is in the range $3 \cdot 10^{-3} > m_a > 10^{-6}$ eV. The lower bound can, however, be relaxed in a number of ways. I discuss the constraint on axion models from the absence of isocurvature perturbations. Dark matter axions can be searched for on Earth by stimulating their conversion to microwave photons in an electromagnetic cavity permeated by a magnetic field. Using this technique, limits on the local halo density have been obtained by the Axion Dark Matter eXperiment.

1 Introduction

The standard model of elementary particles has been an extraordinary breakthrough in our description of the physical world. It agrees with experiment and observation disconcertingly well and surely provides the foundation for all further progress in our field. It does however present us with a puzzle.

Indeed the action density includes, in general, a term

$$\mathcal{L}_{\text{stand mod}} = \cdots + \frac{\theta g^2}{32\pi^2} G^a_{\mu\nu} \tilde{G}^{a\mu\nu} \tag{1}$$

where $G^a_{\mu\nu}$ are the QCD field strengths, g is the QCD coupling constant and θ is a parameter. The dots represent all the other terms in the action density, i.e. the terms that lead to the numerous successes of the standard model. Eq. (1) perversely shows the one term that isn't a success. Using the statement of the chiral anomaly [1], one can show three things about that term. First, that QCD physics depends on the value of the parameter θ because in the absence of such dependence QCD would have a $U_A(1)$ symmetry in the chiral limit, and we know QCD

has no such $U_A(1)$ symmetry [2]. Second, that θ is cyclic, that is to say that physics at θ is indistinguishable from physics at $\theta + 2\pi$. Third, that an overall phase in the quark mass matrix m_q can be removed by a redefinition of the quark fields only if, at the same time, θ is shifted to $\theta - \arg \det m_q$. The combination of standard model parameters $\bar{\theta} \equiv \theta - \arg \det m_q$ is independent of quark field redefinitions. Physics therefore depends on θ solely through $\bar{\theta}$.

Since physics depends on $\bar{\theta}$, the value of $\bar{\theta}$ is determined by experiment. The term shown in Eq. (1) violates P and CP. This source of P and CP violation is incompatible with the experimental upper bound on the neutron electic dipole moment unless $|\bar{\theta}| < 10^{-10}$. A new improved upper limit on the neutron electric dipole moment ($|d_n| < 3.0 \cdot 10^{-26}$ e cm) was reported at this conference by P. Geltenbort [3]. The puzzle aforementioned is why the value of $\bar{\theta}$ is so small. It is usually referred to as the "Strong CP Problem". If there were only strong interactions, a zero value of $\bar{\theta}$ could simply be a consequence of P and CP conservation. That would not be much of a puzzle. But there are also weak interactions and they, and therefore the standard model as a whole, violate P and CP. So these symmetries can not be invoked to set $\bar{\theta} = 0$. More pointedly, P and CP violation are introduced in the standard model by letting the elements of the quark mass matrix m_q be arbitrary complex numbers [4]. In that case, one expects $\arg \det m_q$, and hence $\bar{\theta}$, to be a random angle.

The puzzle is removed if the action density is instead

$$\mathcal{L}_{\text{stand mod} + \text{xion}} = \cdots + \frac{1}{2}\partial_\mu a \partial^\mu a + \frac{g^2}{32\pi^2} \frac{a(x)}{f_a} G^a_{\mu\nu} \tilde{G}^{a\mu\nu} \tag{2}$$

where $a(x)$ is a new scalar field, and the dots represent the other terms of the standard model. f_a is a constant with dimension of energy. The $aG \cdot \tilde{G}$ interaction in Eq. (2) is not renormalizable. However, there is a recipe for constructing renormalizable theories whose low energy effective action density is of the form of Eq. (2). The recipe is as follows: construct the theory in such a way that it has a $U(1)$ symmetry which (1) is a global symmetry of the classical action density, (2) is broken by the color anomaly, and (3) is spontaneously broken. Such a symmetry is called Peccei-Quinn symmetry after its inventors [5]. Weinberg and Wilczek [6] pointed out that a theory with a $U_{\text{PQ}}(1)$ symmetry has a light pseudo-scalar particle, called the axion. The axion field is $a(x)$. f_a is of order the expectation value that breaks $U_{\text{PQ}}(1)$, and is called the "axion decay constant".

In the theory defined by Eq. (2), $\bar{\theta} = \frac{a(x)}{f_a} - \det \arg m_q$ depends on the expectation value of $a(x)$. That expectation value minimizes the effective potential. The Strong CP Problem is then solved because the minimum of the QCD effective potential $V(\bar{\theta})$ occurs at $\bar{\theta} = 0$ [7]. The weak interactions induce a small value for $\bar{\theta}$ [8,9], of order 10^{-17}, but this is consistent with experiment.

The notion of Peccei-Quinn (PQ) symmetry may seem contrived. Why should there be a $U(1)$ symmetry which is broken at the quantum level but which is exact at the classical level? However, the reason for PQ symmetry may be deeper than we know at present. String theory contains many examples of symmetries which are exact classically but which are broken by quantum anomalies, including PQ symmetry [10,11,12]. Within field theory, there are examples of theories with *automatic* PQ symmetry, i.e. where PQ symmetry is a consequence of just the particle content of the theory without adjustment of parameters to special values.

The first axion models had f_a of order the weak interaction scale and it was thought that this was an unavoidable property of axion models. However, it was soon pointed out [13,14] that the value of f_a is really arbitrary, that it is possible to construct axion models with any value of f_a. A value of f_a far from any previously known scale need not lead to a hierarchy problem because PQ symmetry can be broken by the condensates of a new technicolor-like interaction [15].

The properties of the axion can be derived using the methods of current algebra [16]. The

axion mass is given in terms of f_a by

$$m_a \simeq 6 \; eV \; \frac{10^6 GeV}{f_a} \; . \tag{3}$$

All the axion couplings are inversely proportional to f_a. For example, the axion coupling to two photons is:

$$\mathcal{L}_{a\gamma\gamma} = -g_\gamma \frac{\alpha}{\pi} \frac{a(x)}{f_a} \vec{E} \cdot \vec{B} \quad . \tag{4}$$

Here $\vec{E}$ and $\vec{B}$ are the electric and magnetic fields, α is the fine structure constant, and g_γ is a model-dependent coefficient of order one. $g_\gamma = 0.36$ in the DFSZ model [14] whereas $g_\gamma = -0.97$ in the KSVZ model [13]. The coupling of the axion to a spin 1/2 fermion f has the form:

$$\mathcal{L}_{a\bar{f}f} = ig_f \frac{m_f}{f_a} a \bar{f} \gamma_5 f \tag{5}$$

where g_f is a model-dependent coefficient of order one. In the KSVZ model the coupling to electrons is zero at tree level. Models with this property are called 'hadronic'.

The axion has been searched for in many places but not found [17]. The resulting constraints may be summarized as follows. Axion masses larger than about 50 keV are ruled out by particle physics experiments (beam dumps and rare decays) and nuclear physics experiments. The next range of axion masses, in decreasing order, is ruled out by stellar evolution arguments. The longevity of red giants rules out 200 keV $> m_a >$ 0.5 eV [18,19] in the case of hadronic axions, and 200 keV $> m_a > 10^{-2}$ eV [20] in the case of axions with a large coupling to electrons [$g_e = 0(1)$ in Eq. 5]. The duration of the neutrino pulse from Supernova 1987a rules out 2 eV $> m_a > 3 \cdot 10^{-3}$ eV [21]. Finally, there is a lower limit, $m_a \gtrsim 10^{-6}$ eV, from cosmology which will be discussed in detail in the next section. This leaves open an "axion window": $3 \cdot 10^{-3} > m_a \gtrsim 10^{-6}$ eV. We will see, however, that the lower edge of this window (10^{-6} eV) is much softer than its upper edge.

2 Axion cosmology

The implications of the existence of an axion for the history of the early universe may be briefly described as follows. At a temperature of order f_a, a phase transition occurs in which the $U_{PQ}(1)$ symmetry becomes spontaneously broken. This is called the PQ phase transition. At these temperatures, the non-perturbative QCD effects which produce the effective potential $V(\bar{\theta})$ are negligible [22], the axion is massless and all values of $\langle a(x) \rangle$ are equally likely. Axion strings appear as topological defects. One must distinguish two scenarios, depending on wether inflation occurs with reheat temperature lower (case 1) or higher (case 2) than the PQ transition temperature. In case 1 the axion field gets homogenized by inflation and the axion strings are 'blown' away.

When the temperature approaches the QCD scale, the potential $V(\bar{\theta})$ turns on and the axion acquires mass. There is a critical time, defined by $m_a(t_1)t_1 = 1$, when the axion field starts to oscillate in response to the turn-on of the axion mass. The corresponding temperature $T_1 \simeq 1$ GeV [23]. The initial amplitude of this oscillation is how far from zero the axion field lies when the axion mass turns on. The axion field oscillations do not dissipate into other forms of energy and hence contribute to the cosmological energy density today [23]. This contribution is called of 'vacuum realignment'. It is further described below. Note that the vacuum realignment contribution may be accidentally suppressed in case 1 if the homogenized axion field happens to lie close to zero.

In case 2 the axion strings radiate axions [24,25] from the time of the PQ transition till t_1 when the axion mass turns on. At t_1 each string becomes the boundary of N domain walls. If $N = 1$,

the network of walls bounded by strings is unstable [26,27] and decays away. If $N > 1$ there is a domain wall problem [28] because axion domain walls end up dominating the energy density, resulting in a universe very different from the one observed today. There is a way to avoid this problem by introducing an interaction which slightly lowers one of the N vacua with respect to the others. In that case, the lowest vacuum takes over after some time and the domain walls disappear. There is little room in parameter space for that to happen and we will not consider this possibility [29] further here. Henceforth, we assume $N = 1$.

In case 2 there are three contributions to the axion cosmological energy density. One contribution [24,25,30,31,32,33,34] is from axions that were radiated by axion strings before t_1. A second contribution is from axions that were produced in the decay of walls bounded by strings after t_1 [31,35,36,29]. A third contribution is from vacuum realignment [23].

Let me briefly indicate how the vacuum alignment contribution is evaluated. Before time t_1, the axion field did not oscillate even once. Soon after t_1, the axion mass is assumed to change sufficiently slowly that the total number of axions in the oscillations of the axion field is an adiabatic invariant. The average number density of axions at time t_1 is

$$n_a(t_1) \simeq \frac{1}{2} m_a(t_1) \langle a^2(t_1) \rangle \simeq \pi f_a^2 \frac{1}{t_1} \tag{6}$$

In Eq. (6), we used the fact that the axion field $a(x)$ is approximately homogeneous on the horizon scale t_1. Wiggles in $a(x)$ which entered the horizon long before t_1 have been red-shifted away [37]. We also used the fact that the initial departure of $a(x)$ from the nearest minimum is of order f_a. The axions of Eq. (6) are decoupled and non-relativistic. Assuming that the ratio of the axion number density to the entropy density is constant from time t_1 till today, one finds [23,29]

$$\Omega_a \simeq \frac{1}{2} \left(\frac{f_a}{10^{12}\text{GeV}} \right)^{\frac{7}{6}} \left(\frac{0.7}{h} \right)^2 \tag{7}$$

for the ratio of the axion energy density to the critical density for closing the universe. h is the present Hubble rate in units of 100 km/s.Mpc. The requirement that axions do not overclose the universe implies the constraint $m_a \gtrsim 6 \cdot 10^{-6}$ eV.

The contribution from axion string decay has been debated over the years. The main issue is the energy spectrum of axions radiated by axion strings. Battye and Shellard [32] have carried out computer simulations of bent strings (i.e. of wiggles on otherwise straight strings) and have concluded that the contribution from string decay is approximately ten times larger than that from vacuum realignment, implying a bound on the axion mass approximately ten times more severe, say $m_a \gtrsim 6 \cdot 10^{-5}$ eV instead of $m_a \gtrsim 6 \cdot 10^{-6}$ eV. My collaborators and I have done simulations of bent strings [31], of circular string loops [31,34] and non-circular string loops [34]. We conclude that the string decay contribution is of the same order of magnitude than that from vacuum realignment. Yamaguchi, Kawasaki and Yokoyama [33] have done computer simulations of a network of strings in an expanding universe, and concluded that the contribution from string decay is approximately three times that of vacuum realignment. The contribution from wall decay has been discussed in detail in ref. [29]. It is probably subdominant compared to the vacuum realignment and string decay constributions.

It should be emphasized that there are many sources of uncertainty in the cosmological axion energy density aside from the uncertainty about the contribution from string decay. The axion energy density may be diluted by the entropy release from heavy particles which decouple before the QCD epoch but decay afterwards [38], or by the entropy release associated with a first order QCD phase transition. On the other hand, if the QCD phase transition is first order [39], an abrupt change of the axion mass at the transition may increase Ω_a. If inflation occurs with reheat temperature less than T_{PQ}, there may be an accidental suppression of Ω_a because the homoger ized axion field happens to lie close to a CP conserving minimum. Because the

RHS of Eq. (7) is multiplied in this case by a factor of order the square of the initial vacuum misalignment angle $\frac{a(t_1)}{f_a}$ which is randomly chosen between $-\pi$ and $+\pi$, the probability that Ω_a is suppressed by a factor x is of order $\sqrt{x}$. Recently, Kaplan and Zurek proposed a model [40] in which the axion decay constant f_a is time-dependent, the value $f_a(t_1)$ during the QCD phase-transition being much smaller than the value f_a today. This yields a suppression of the axion cosmological energy density by a factor $(\frac{f_a(t_1)}{f_a})^2$ compared to the usual case [replace f_a by $f_a(t_1)$ in Eq. (6)]. Finally, the axion density may be diluted by 'coherent deexcitation', i.e. adiabatic level crossing of $m_a(t)$ with the mass of some other pseudo-Nambu-Goldstone boson which mixes with the axion [41].

The axions produced when the axion mass turns on during the QCD phase transition are cold dark matter (CDM) because they are non-relativistic from the moment of their first appearance at 1 GeV temperature. Studies of large scale structure formation support the view that the dominant fraction of dark matter is CDM. Any form of CDM necessarily contributes to galactic halos by falling into the gravitational wells of galaxies. Hence, there is excellent motivation to look for axions as constituent particles of our galactic halo.

There is a particular kind of clumpiness [42,29] which affects axion dark matter if there is no inflation after the Peccei-Quinn phase transition (case 2). This is due to the fact that the dark matter axions are inhomogeneous with $\delta\rho/\rho \sim 1$ over the horizon scale at temperature $T_1 \simeq 1$ GeV, when they are produced at the start of the QCD phase-transition, combined with the fact that their velocities are so small that they do not erase these inhomogeneities by free-streaming before the time t_{eq} of equality between the matter and radiation energy densities when matter perturbations can start to grow. These particular inhomogeneities in the axion dark matter are in the non-linear regime immediately after time t_{eq} and thus form clumps, called 'axion mini-clusters' [42]. They have mass $M_{mc} \simeq 10^{-13}M_\odot$ and size $l_{mc} \simeq 10^{13}$ cm.

3   Axion isocurvature perturbations

If inflation occurs after the Peccei-Quinn phase transition, i.e. if the reheat temperature after inflation T_{RH} is less than the temperature T_{PQ} at which $U_{\mathrm{PQ}}(1)$ is restored (case 1), the quantum mechanical fluctuations of the axion field during the inflationary epoch cause isocurvature density perturbations [43,44] in the early universe. The cosmic microwave background obervations are consistent with purely adiabatic density perturbations and therefore place a constraint, which we now discuss.

Fluctuations generated during inflation in a massless weakly coupled scalar field, such as the inflaton or the axion, are characterized by the power spectrum [45]

$$P_a(k) \equiv \int \frac{d^3x}{(2\pi)^3} < \delta a(\vec{x},t)\delta a(\vec{x}',t) > e^{-i\vec{k}\cdot(\vec{x}-\vec{x}')} = \left(\frac{H_I}{2\pi}\right)^2 \frac{2\pi^2}{k^3} \quad , \tag{8}$$

where $\vec{x}$ are comoving coordinates. Eq. (8) is often written in the shorthand notation $\delta a = \frac{H_I}{2\pi}$. The axion field fluctuations are "frozen" from the time their wavelengths exceed the horizon size H_I^{-1} during the inflationary epoch till the time their wavelenghs reenter the horizon long after inflation has ended.

At the start of the QCD phase transition, the local value of the axion field $a(\vec{x},t)$ determines the local number density of cold axions produced by the vacuum realignment mechanism

$$n_a(\vec{x},t_1) = \frac{f_a^2}{2t_1}\alpha(\vec{x},t_1)^2 \tag{9}$$

where $\alpha(\vec{x},t_1) = a(\vec{x},t_1)/f_a$ is the local misalignment angle. The fluctuations in the axion field

produce perturbations in the axion dark matter density

$$\frac{\delta n_a^{\mathrm{iso}}}{n_a} = \frac{2\delta a}{a_1} = \frac{H_I}{\pi f_a \alpha_1} \qquad (10)$$

where $a_1 = a(t_1) = f_a \alpha_1$ is the initial value of the axion field, at the start of the QCD phase transition, common to our entire visible universe. These perturbations initially obey $\delta \rho_a^{\mathrm{iso}} = -\delta \rho_r^{\mathrm{iso}}$ since the vacuum realignment mechanism converts energy stored in the quark-gluon plasma into axion rest mass energy. Such perturbations are commonly called "isocurvature perturbations" because they do not initially produce a source for the Newtonian potential $\Phi = \frac{1}{2}(g_{00} - 1)$. Note that in case 1 the density perturbations in the cold axion fluid have both adiabatic and isocurvature components. The adiabatic perturbations ($\frac{\delta \rho_a^{\mathrm{ad}}}{3\rho_a} = \frac{\delta \rho_r^{\mathrm{iso}}}{4\rho_r} = \frac{\delta T}{T}$) are produced by the quantum mechanical fluctuations of the inflaton field during inflation, whereas the isocurvature perturbations $[\delta \rho_a^{\mathrm{iso}}(t_1) \simeq -\delta \rho_r(t_1)^{\mathrm{iso}}]$ are produced by the quantum mechanical fluctuations of the axion field during that same epoch.

Isocurvature perturbations make a different imprint on the cosmic microwave background than do adiabatic ones. The CMBR observations are consistent with pure adiabatic perturbations. According to P. Crotty et al. [46], the fraction of cold dark matter perturbations which are isocurvature can not be larger than 31%. This places a constraint on axion models if the Peccei-Quinn phase transition occurs before inflation (case 1). Allowing for the possibility that only part of the cold dark matter is axions, the CMBR constraint of ref. [46] is

$$\frac{\delta \rho_a^{\mathrm{iso}}}{\rho_{\mathrm{CDM}}} = \frac{\delta \rho_a^{\mathrm{iso}}}{\rho_a} \cdot \frac{\rho_a}{\rho_{\mathrm{CDM}}} = \frac{H_I}{\pi f_a \alpha_1} \frac{\Omega_a}{\Omega_{\mathrm{CDM}}} < 0.31 \frac{\delta \rho_m}{\rho_m} \qquad , \qquad (11)$$

where we used Eq. (10). $\frac{\delta \rho_m}{\rho_m}$ is the amplitude of the primordial spectrum of matter perturbations. It is related to the amplitude of large scale (low multipole) CMBR anisotropies through the Sachs-Wolfe effect [47]. The observations imply [48] $\frac{\delta \rho_m}{\rho_m} \simeq 4.6 \; 10^{-5}$.

In terms of α_1, the cold axion energy density is given in case 1 by

$$\Omega_a \simeq 0.15 \left(\frac{f_a}{10^{12}\mathrm{GeV}}\right)^{\frac{7}{6}} \alpha_1^2 \qquad . \qquad (12)$$

where we assumed $h \simeq 0.7$. It has been remarked by many authors, starting with S.-Y. Pi [49], that it is possible for f_a to be much larger than 10^{12}GeV because α_1 may be accidentally small in our visible universe. The requirement that $\Omega_a < \Omega_{\mathrm{CDM}} = 0.22$ [50] implies

$$|\frac{\alpha_1}{\pi}| < 0.5 \left(\frac{10^{12}\mathrm{GeV}}{f_a}\right)^{\frac{7}{12}} \qquad . \qquad (13)$$

Since $-\pi < \alpha_1 < +\pi$ is the a-priori range of α_1 values and no particular value is preferred over any other, $|\frac{\alpha_1}{\pi}|$ may be taken to be the "probability" that the initial misalignment angle has magnitude less than $|\alpha_1|$. If $|\frac{\alpha_1}{\pi}| = 4 \cdot 10^{-3}$, for example, f_a may be as large as 10^{16} GeV.

The presence of isocurvature perturbations constrains the small α_1 scenario in two ways [44]. First, it makes it impossible to have α_1 arbitrarily small since

$$\alpha_1 > \delta \alpha_1 = \frac{\dot{H}_I}{2\pi f_a} \qquad . \qquad (14)$$

Combining Eqs. (13) and (14), we obtain the bound

$$\Lambda_I < 6 \cdot 10^{15}\mathrm{GeV} \left(\frac{f_a}{10^{12}\mathrm{GeV}}\right)^{\frac{5}{24}} \qquad . \qquad (15)$$

Second, one must require axion isocurvature perturbations to be consistent with CMBR observations. Combining Eqs. (11) and (12), and setting $\Omega_{\text{CDM}} = 0.22$, $\frac{\delta\rho_m}{\rho_m} = 4.6 \ 10^{-5}$, we obtain

$$\Lambda_I < 10^{13}\text{GeV} \ \ \Omega_a^{-\frac{1}{4}} \ \left(\frac{f_a}{10^{12}\text{GeV}}\right)^{\frac{5}{24}} \ . \tag{16}$$

Let us keep in mind that the bounds (15) and (16) pertain only if $T_{\text{RH}} < T_{\text{PQ}}$. One may, for example, have $\Omega_a = 0.22$, $f_a \simeq 10^{12}$ GeV, and $\Lambda_I \simeq 10^{16}$ GeV, provided $T_{\text{RH}} \gtrsim 10^{12}$ GeV, which is possible if reheating is sufficiently efficient.

4 Dark matter axion detection

An electromagnetic cavity permeated by a strong static magnetic field can be used to detect galactic halo axions [51]. The relevant coupling is given in Eq. (4). Galactic halo axions have velocities β of order 10^{-3} and hence their energies $E_a = m_a + \frac{1}{2}m_a\beta^2$ have a spread of order 10^{-6} above the axion mass. When the frequency $\omega = 2\pi f$ of a cavity mode equals m_a, galactic halo axions convert resonantly into quanta of excitation (photons) of that cavity mode. The power from axion $\rightarrow$ photon conversion on resonance is found to be [51,52]:

$$
\begin{aligned}
P &= \left(\frac{\alpha \, g_\gamma}{\pi \, f_a}\right)^2 V B_0^2 \rho_a C \frac{1}{m_a} \text{Min}(Q_L, Q_a) \\
&= 0.5 \ 10^{-26} \text{Watt} \left(\frac{V}{500 \text{ liter}}\right) \left(\frac{B_0}{7 \text{ Tesla}}\right)^2 C \left(\frac{g_\gamma}{0.36}\right)^2 \ . \\
&\quad \cdot \left(\frac{\rho_a}{\frac{1}{2} \cdot 10^{-24} \frac{\text{gr}}{\text{cm}^3}}\right) \left(\frac{m_a}{2\pi(\text{GHz})}\right) \text{Min}(Q_L, Q_a)
\end{aligned}
\tag{17}
$$

where V is the volume of the cavity, B_0 is the magnetic field strength, Q_L is its loaded quality factor, $Q_a = 10^6$ is the 'quality factor' of the galactic halo axion signal (i.e. the ratio of their energy to their energy spread), ρ_a is the density of galactic halo axions on Earth, and C is a mode dependent form factor given by

$$C = \frac{\left|\int_V d^3x \vec{E}_\omega \cdot \vec{B}_0\right|^2}{B_0^2 V \int_V d^3x\epsilon|\vec{E}_\omega|^2} \tag{18}$$

where $\vec{B}_0(\vec{x})$ is the static magnetic field, $\vec{E}_\omega(\vec{x})e^{i\omega t}$ is the oscillating electric field and ϵ is the dielectric constant. Because the axion mass is only known in order of magnitude at best, the cavity must be tunable and a large range of frequencies must be explored seeking a signal. The cavity can be tuned by moving a dielectric rod or metal post inside it.

For a cylindrical cavity and a homogeneous longitudinal magnetic field, $C = 0.69$ for the lowest TM mode. The form factors of the other modes are much smaller. The resonant frequency of the lowest TM mode of a cylindrical cavity is $f=115$ MHz $\left(\frac{1m}{R}\right)$ where R is the radius of the cavity. Since 10^{-6} eV $= 2\pi$ (242 MHz), a large cylindrical cavity is convenient for searching the low frequency end of the range of interest. To extend the search to high frequencies without sacrifice in volume, one may power-combine many identical cavities which fill up the available volume inside a magnet's bore [53,54]. This method allows one to maintain $C = 0(1)$ at high frequencies, albeit at the cost of increasing engineering complexity as the number of cavities increases.

Axion dark matter searches were carried out at Brookhaven National Laboratory [55], the University of Florida [56], Kyoto University [57], and by the ADMX collaboration [58,59,60,61,62] at Lawrence Livermore National Laboratory. Thus far, the ADMX experiment has ruled out KSVZ

coupled axions at the nominal halo density of [63] $\rho_a = 7.5 \ 10^{-25}$ g/cm^3 over the mass range $1.90 < m_a < 3.35$ μeV [60,61]. The ADMX experiment is presently being upgraded to replace the HEMT (high electron mobility transistors) receivers we have used so far with SQUID microwave amplifiers. HEMT receivers have noise temperature $T_n \sim 3 \ K$ [65] whereas $T_n \sim 0.05 \ K$ was achieved with SQUIDs[66]. In a second phase of the upgrade, the experiment will be equipped with a dilution refrigerator to take full advantage of the lowered electronic noise temperature. When both phases of the upgrade are completed, the ADMX detector will have sufficient sensitivity to detect axions at even a fraction of the halo density.

The ADMX experiment is equipped with a high resolution spectrometer which allows us to look for narrow peaks in the spectrum of microwave photons caused by discrete flows, or streams, of dark matter axions in our neighborhood. In many discussions of cold dark matter detection it is assumed that the distribution of CDM particles in galactic halos is isothermal. However, there are excellent reasons to believe that a large fraction of the local density of cold dark matter particles is in discrete flows with definite velocities [64]. Indeed, because CDM has very low primordial velocity dispersion and negligible interactions other than gravity, the CDM particles lie on a 3-dim. hypersurface in 6-dim. phase-space. This implies that the velocity spectrum of CDM particles at any physical location is discrete, i.e., it is the sum of distinct flows each with its own density and velocity.

We searched for the peaks in the spectrum of microwave photons from axion to photon conversion that such discrete flows would cause in the ADMX detector. We found none and placed limits [62] on the density of any local flow of axions as a function of the flow velocity dispersion over the axion mass range 1.98 to 2.17 μeV. Our limit on the density of discrete flows is approximately a factor three more severe than our limit on the total local axion dark matter density.

Acknowledgments

References

1. S. Adler, Phys. Rev. **177** (1969) 2426; J.S. Bell and R. Jackiw, Nuov. Cim. **60A** (1969) 47.
2. S. Weinberg, Phys. Rev. **D11** (1975) 3583.
3. C.A. Baker et al., hep-ex/0602020.
4. M. Kobayashi and K. Maskawa, Progr. Theor. Phys. **49** (1973) 652.
5. R. D. Peccei and H. Quinn, Phys. Rev. Lett. **38** (1977) 1440 and Phys. Rev. **D16** (1977) 1791.
6. S. Weinberg, Phys. Rev. Lett. **40** (1978) 223; F. Wilczek, Phys. Rev. Lett. **40** (1978) 279.
7. C. Vafa and E. Witten, Phys. Rev. Lett. **53** (1984) 535.
8. J. Ellis and M.K. Gaillard, Nucl. Phys. **B150** (1979) 141.
9. H. Georgi and L. Randall, Nucl. Phys. **B276** (1986) 241.
10. E. Witten, Phys. Lett. **B149** (1984) 351.
11. K. Choi and J.E. Kim, Phys. Lett. **B154** (1985) 393; *ibidem* **B165** (1985) 689.
12. P. Svrcek and E. Witten, hep-th/0605206.
13. J. Kim, Phys. Rev. Lett. **43** (1979) 103; M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, Nucl. Phys. **B166** (1980) 493.

14. M. Dine, W. Fischler and M. Srednicki, Phys. Lett. **B104** (1981) 199; A. P. Zhitnitskii, Sov. J. Nucl. **31** (1980) 260.

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

16. S. Weinberg in ref.[6]; W.A. Bardeen and S.-H.H. Tye, Phys. Lett. **B74** (1978) 229; J. Ellis and M.K. Gaillard, Nucl. Phys. **B150** (1979) 141; T.W. Donnelly et al., Phys. Rev. **D18** (1978) 1607; D.B. Kaplan, Nucl. Phys. **B260** (1985) 215; M. Srednicki, Nucl. Phys. **B260** (1985) 689; P. Sikivie, in 'Cosmology and Particle Physics', ed. E. Alvarez et al., World Scientific, 1987, pp 143-169.

17. Axion reviews include: J.E. Kim, Phys. Rep. **150** (1987) 1; H.-Y. Cheng, Phys. Rep. **158** (1988) 1; R.D. Peccei, in 'CP Violation', ed. by C. Jarlskog, World Scientific Publ., 1989, pp 503-551; M.S. Turner, Phys. Rep. **197** (1990) 67; G.G. Raffelt, Phys. Rep. **198** (1990) 1.

18. M.I.Vysotsky, Ya.B.Zeldovich, M.Yu.Khlopov and V.M.Chechetkin, JETP Lett. **27** (1978) 502; D. Dicus, E. Kolb, V. Teplitz and R. Wagoner, Phys. Rev. **D18** (1978) 1829 and Phys. Rev. **D22** (1980) 839.

19. G. Raffelt and D. Dearborn, Phys. Rev. **D36** (1987) 2211.

20. D. Dearborn, D. Schramm and G. Steigman, Phys. Rev. Lett. **56** (1986) 26.

21. J. Ellis and K. Olive, Phys. Lett. **B193** (1987) 525; G. Raffelt and D. Seckel, Phys. Rev. Lett. **60** (1988) 1793; M. Turner, Phys. Rev. Lett. **60** (1988) 1797; H.-T. Janka et al., Phys. Rev. Lett. **76** (1996) 2621; W. Keil et al.. Phys. Rev. **D56** (1997) 2419.

22. D. J. Gross, R. D. Pisarski and L. G. Yaffe, Rev. Mod. Phys. **53** (1981) 43.

23. L. Abbott and P. Sikivie, Phys. Lett. **B120** (1983) 133; J. Preskill, M. Wise and F. Wilczek, Phys. Lett. **B120** (1983) 127; M. Dine and W. Fischler, Phys. Lett. **B120** (1983) 137.

24. R. Davis, Phys. Rev. **D32** (1985) 3172 and Phys. Lett. **B180**

25. D. Harari and P. Sikivie, Phys. Lett. **B195** (1987) 361.

26. A. Vilenkin and A.E. Everett, Phys. Rev. Lett. **48** (1982) 1867.

27. P. Sikivie in *Where are the elementary particles*, Proc. of the 14th Summer School on Particle Physics, Gif-sur-Yvette, 1982, edited by P. Fayet et al. (Inst. Nat. Phys. Nucl. Phys. Particules, Paris, 1983).

28. P. Sikivie, Phys. Rev. Lett. **48** (1982) 1156.

29. S. Chang, C. Hagmann and P. Sikivie, Phys. Rev. **D59** (1999) 023505.

30. A. Vilenkin and T. Vachaspati, Phys. Rev. **D35** (1987) 1138; R.L. Davis and E.P.S. Shellard, Nucl. Phys. **B324** (1989) 167; A. Dabholkar and J. Quashnock, Nucl. Phys. **B333** (1990) 815.

31. C. Hagmann and P. Sikivie, Nucl. Phys. **B363** (1991) 247.

32. R.A. Battye and E.P.S. Shellard, Nucl. Phys. **B423** (1994) 260, Phys. Rev. Lett. **73** (1994) 2954 and erratum-ibid. **76** (1996) 2203.

33. M. Yamaguchi, M. Kawasaki and J. Yokoyama, Phys. Rev. Lett. **82** (1999) 4578.

34. C. Hagmann, S. Chang and P. Sikivie, Phys. Rev. **D63** (2001) 125018.

35. D. Lyth, Phys. Lett. **B275** (1992) 279;

36. M. Nagaswa and M. Kawasaki, Phys. Rev. **D50** (1994) 4821.

37. A. Vilenkin, Phys. Rev. Lett. **48** (1982) 59.

38. P. J. Steinhardt and M. S. Turner, Phys. Lett. **B129** (1983) 51; G. Lazarides, R. Schaefer, D. Seckel and Q. Shafi, Nucl. Phys. **B346** (1990) 193.

39. W. G. Unruh and R. M. Wald, Phys. Rev. **D32** (1985) 831; M. S. Turner, Phys. Rev. **D32** (1985) 843; T. DeGrand, T. W. Kephart and T. J. Weiler, Phys. Rev. **D33** (1986) 910; M. Hindmarsh, Phys. Rev. **D45** (1992) 1130.

40. D.B. Kaplan and K.M. Zurek, Phys. Rev. Lett. **96** (2006) 041301.

41. C.T. Hill and G.G. Ross, Nucl. Phys. **B311** (1988) 253.

42. C. J. Hogan and M. J. Rees, Phys. Lett. **B205** (1988) 228; E. Kolb and I. I. Tkachev, Phys. Rev. Lett. **71** (1993) 3051, Phys. Rev. **D49** (1994) 5040, and Ap. J. **460** (1996) L25.

43. M. Axenides, R.H. Brandenberger and M.S. Turner, Phys. Lett. **B126** (1983) 178; P. J. Steinhardt and M. S. Turner, Phys. Lett. **B129** (1983) 51; A.D. Linde, Phys. Lett. **B158** (1985) 375; D. Seckel and M. Turner, Phys. Rev. **D32** (1985) 3178; D.H. Lyth, Phys. Lett. **B236** (1990) 408;

44. M. Turner and F. Wiczek, Phys. Rev. Lett. **66** (1991) 5.

45. N.D. Birrell and P.C.W. Davies, *Quantum Field Theory in Curved Space-Time*, Cambridge University Press, 1982; A. Vilenkin and L.H. Ford, Phys. Rev. **D25** (1982) 175; A.D. Linde, Phys. Lett. **B108** (1982) 389; A. Starobinsky, Phys. Lett. **B117** (1982) 175.

46. P. Crotty, J. Garcia-Bellido, J. Lesgourgues and A. Riazuelo, Phys. Rev. Lett. **91** (2003) 171301

47. R. Sachs and A. Wolfe, Ap. J. **147** (1967) 73; P.J.E. Peebles, Ap. J. **263** (1982) L1; L.F. Abbott and M.B. Wise, Ap. J. **282** (1984) L47.

48. S. Dodelson, *Modern Cosmology*, Academic Press, 2003.

49. S.-Y. Pi, Phys. Rev. Lett. **52** (1984) 1725.

50. D.N. Spergel et al., astro-ph/0603449.

51. P. Sikivie, Phys. Rev. Lett. **51** (1983) 1415 and Phys. Rev. **D32** (1985) 2988.

52. L. Krauss, J. Moody, F. Wilczek and D. Morris, Phys. Rev. Lett. **55** (1985) 1797.

53. C. Hagmann et al., Rev. Sci. Inst. **61** (1990) 1076.

54. C. Hagmann, Ph. D. thesis, unpublished.

55. S. DePanfilis et al., Phys. Rev. Lett. **59** (1987) 839 and Phys. Rev. **40** (1989) 3153.

56. C. Hagmann et al., Phys. Rev. **D42** (1990) 1297.

57. S. Matsuki and K. Yamamoto, Phys. Lett. **B263** (1991) 523; S. Matsuki, I. Ogawa, K. Yamamoto, Phys. Lett. **B336** (1994) 573; I. Ogawa, S. Matsuki and K. Yamamoto, Phys. Rev. **D53** (1996) R1740; A. Kitagawa, K. Yamamoto and S. Matsuki, hep-ph/9908445.

58. C. Hagmann et al., Phys. Rev. Lett. **80** (1998) 2043.

59. S.J. Asztalos et al., Phys. Rev. **D64** (2001) 092003.

60. S.J. Asztalos et al., Ap. J. Lett. **571** (2002) L27.

61. S.J. Asztalos et al., Phys. Rev. **D69** (2004) 011101 (R).

62. L. Duffy et al., Phys. Rev. Lett. **95** (2005) 091304.

63. E.I. Gates, G. Gyuk and M.S. Turner, Astroph. J. **449** (1995) 123.

64. J.R. Ipser and P. Sikivie, Phys. Lett. **B291** (1992) 288; A. Natarajan and P. Sikivie. Phys. Rev. **D72** (2005) 083513.

65. R.F. Bradley, Nucl. Phys. B Proc. Suppl. **72** (1999) 137.

66. M. Mück et al., Appl. Phys. Lett. **72** (1998) 2885; Nucl. Phys. B Proc. Suppl. **72** (1999) 145.