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DARK MATTER AXIONS '93 *

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I. Introduction to axion physics
II. The cosmological axion energy density
III. Galactic halo axion searches

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* Lecture at the Workshop "The Dark Side of the Universe", in Rome, Italy, June 23-25, 1993 and at the 17th John Hopkins Workshop on Current Problems in Particle Theory "Particles and the Universe", in Budapest, Hungary, June 30-August 1, 1993.

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I. INTRODUCTION TO AXION PHYSICS

The axion[1,2,3] was postulated sixteen years ago to explain why the strong interactions conserve P and CP. Consider the Lagrangian of QCD:

$$\begin{aligned}\mathcal{L}_{QCD} = &-\frac{1}{4}G^a_{\mu\nu}G^{a\mu\nu} + \sum_{i=1}^{n}[\bar{q}_i\gamma^\mu D_\mu q_i - m_i q^\dagger_{L_i} q_{R_i} - m_i^* q^\dagger_{R_i} q_{L_i}] \\ &+ \frac{\theta g^2}{32\pi^2}G^a_{\mu\nu}\tilde{G}^{a\mu\nu}.\end{aligned} \tag{1.1}$$

The last term is a 4-divergence and hence does not contribute in perturbation theory. It does however produce non-perturbative effects associated with the existence of QCD instantons[4]. As a consequence, the physics of QCD depends upon the value of the parameter θ. Using the Adler-Bell-Jackiw anomaly of the $U_A(1)$ current, one can readily show that the physics of QCD depends upon θ only through the combination of parameters

$$\bar{\theta} \equiv \theta - arg\ det\ m_q = \theta - arg(m_1 m_2 ... m_n)\ . \tag{1.2}$$

If $\bar{\theta} \neq 0$, QCD violates P and CP. The absence of P and CP violations in the strong interactions therefore places an upper limit upon $\bar{\theta}$. The best constraint comes from the present experimental bound on the neutron electric dipole moment which yields[5,6]:

$$\bar{\theta} \lesssim 10^{-9}\ . \tag{1.3}$$

We must then face the question: why is $\bar{\theta}$ so small? Recall that in the standard model of particle interactions, the quark masses originate in the electroweak

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sector of the theory. This sector must violate P and CP to produce the correct weak interaction phenomenology, in particular $K_L \to 2\pi$ decay. There is no reason in the standard model to expect the overall phase of the quark mass matrix to exactly match the value of θ from the QCD sector in order to set $\bar{\theta} < 10^{-9}$. In particular, if CP violation is introduced in the manner of Kobayashi and Maskawa[7], the Yukawa couplings that give rise to the quark masses are arbitrary complex numbers and hence $arg\ det\ m_q$ and $\bar{\theta}$ have no reason to take on any special value at all.

Peccei and Quinn[1] proposed a solution to this problem by postulating the existence of a global $U_{PQ}(1)$ quasi-symmetry. This $U_{PQ}(1)$ must have the following properties

1. it is a symmetry of the classical theory, i.e. a symmetry of the theory at the Lagrangian level,
2. it is broken explicitly by those non-perturbative QCD effects (instantons and the like) which make the physics of QCD depend upon the parameter θ,
3. it is broken spontaneously by the vacuum expectation value of some scalar field.

To see how the existence of a $U_{PQ}(1)$ quasi-symmetry yields $\bar{\theta} = 0$ (up to tiny corrections due to the CP violating interactions responsible for $K_L \to 2\pi$ decay) consider the theory defined by

$$\begin{aligned}\mathcal{L}_{PQ} = &-\frac{1}{4}G^a_{\mu\nu}G^{a\mu\nu} + \frac{1}{2}\partial_\mu\varphi^\dagger\partial^\mu\varphi + \sum_{i=1}^{n}\left[\bar{q}_i\gamma^\mu D_\mu q_i - (K_i q^\dagger_{L_i} q_{R_i}\varphi + h.c.)\right] \\ &+ \frac{\theta g^2}{32\pi^2}G^a_{\mu\nu}\tilde{G}^{a\mu\nu} - V(\varphi^\dagger\varphi).\end{aligned} \tag{1.4}$$

$\mathcal{L}_{PQ}$ has a classical $U_{PQ}(1)$ symmetry under which $\varphi \to e^{i\alpha}\varphi$ and $q_j \to e^{-i\alpha\gamma_5/2}q_j$ for j=1...n. Assuming the potential V has the shape of a "Mexican hat", the $U_{PQ}(1)$ symmetry is spontaneously broken by the vacuum expectation value of the scalar field φ:

$$\langle\varphi(x)\rangle = v\ e^{i\alpha(x)}. \tag{1.5}$$

The quarks acquire masses

$$m_j = K_j\ v\ e^{i\alpha} \tag{1.6}$$

and hence

$$\begin{aligned}\bar{\theta} &= \theta - arg\ (m_1...m_n) \\ &= \theta - arg\ (K_1...K_n) - n\alpha\ .\end{aligned} \tag{1.7}$$

The important difference between the theory defined by Eq. (1.1) and the theory defined by Eq. (1.4) is that in the former $\bar{\theta}$ is a function solely of the parameters in the theory, whereas in the latter $\bar{\theta}$ is a function also of the dynamical field α. As a result, those non-perturbative effects which make the physics of QCD depend upon the parameter $\bar{\theta}$ will, in the theory defined by Eq. (1.4), produce an effective potential $V_{eff}(\alpha)$. It can be shown that the absolute minimum of $V_{eff}(\alpha)$ occurs at a value of α such that[8]

$$\bar{\theta} = 0\ , \tag{1.8}$$

and hence the theory conserves P and CP. For pedagogical reasons, we did not include the electroweak interactions in our example of Eq. (1.4). Actually, the electroweak interactions only add a few mostly non-essential complications to the implementation of the Peccei-Quinn mechanism. The main difference from the description given above is that the CP violating interactions responsible for $K_L \to 2\pi$ will induce a small value for $\bar{\theta}$. This induced value of $\bar{\theta}$ is however always much smaller than 10^{-9}. For example, $\bar{\theta} \approx 10^{-17}$ if CP violation is introduced into the model in the manner of Kobayashi and Maskawa.[9]

Weinberg and Wilczek[2] pointed out that the Peccei-Quinn solution to the strong CP problem implies the existence of a light pseudo-scalar particle, called the axion. The axion is the Pseudo-Nambu-Goldstone boson associated with the spontaneous breaking of the $U_{PQ}(1)$ quasi-symmetry. Thus the axion field is

$$a(x) = v\ \alpha(x) \tag{1.9}$$

where v and $\alpha(x)$ are defined by Eq.(1.5). One can calculate the properties[2,10] of the axion using current algebra or chiral Lagrangian techniques. The axion mass is given by

$$m_a = \mid N \mid \frac{f_\pi m_\pi}{v} \frac{\sqrt{m_u m_d}}{m_u + m_d} \simeq 0.6eV(\frac{10^7 GeV}{v/\mid N \mid}) \tag{1.10}$$

(in the limit $m_s >> m_u,\ m_d$), where

$$N = 2\sum_f t_f Q_f^{PQ}\ . \tag{1.11}$$

In Eq. (1.11), the sum is over all colored left-handed Weyl fermions f, the Q_f^{PQ} are the Peccei-Quinn charges of these fermions and the t_f are given by $Tr(T_f^\alpha T_f^\beta) = t_f \delta^{\alpha\beta}$ where the $T_f^\alpha(\alpha = 1...8)$ are the $SU^c(3)$ generators for the color representation to which f belongs ($t_3 = t_{\bar{3}} = \frac{1}{2},\ t_6 = t_{\bar{6}} = \frac{5}{2}$, etc.). For example, N = n in the model of Eq.(1.7). I will often use $f_a \equiv \frac{v}{|N|}$. The coupling of the axion to two photons is given by

$$\mathcal{L}_{a\gamma\gamma} = \frac{\alpha}{8\pi}\frac{a}{f_a}[\frac{N_e}{N} - (\frac{5}{3} + \frac{m_d - m_u}{m_d + m_u})]F_{\mu\nu}\tilde{F}^{\mu\nu}\ . \tag{1.12}$$

Here α is the fine structure constant and

$$N_e = 2\sum_f (Q_f^\gamma)^2 Q_f^{PQ}\ , \tag{1.13}$$

where Q_f^γ is the electric charge of fermion f in units of e. In many axion models $\frac{N_e}{N} = \frac{8}{3}$. In particular, this is true of all grand unified axion models which implement the (nearly) successful Georgi-Quinn-Weinberg[11] prediction of $sin^2\theta_w$. Comparing Eq.(1.10) with (1.12), one finds that in these models $\mathcal{L}_{a\gamma\gamma}$ is given uniquely in terms of the axion mass by

$$\mathcal{L}_{a\gamma\gamma} = -g_\gamma \frac{\alpha}{\pi}\frac{m_a}{0.6\ 10^{16}\ eV^2}\ a\ \vec{E}\cdot\vec{B} \tag{1.14}$$

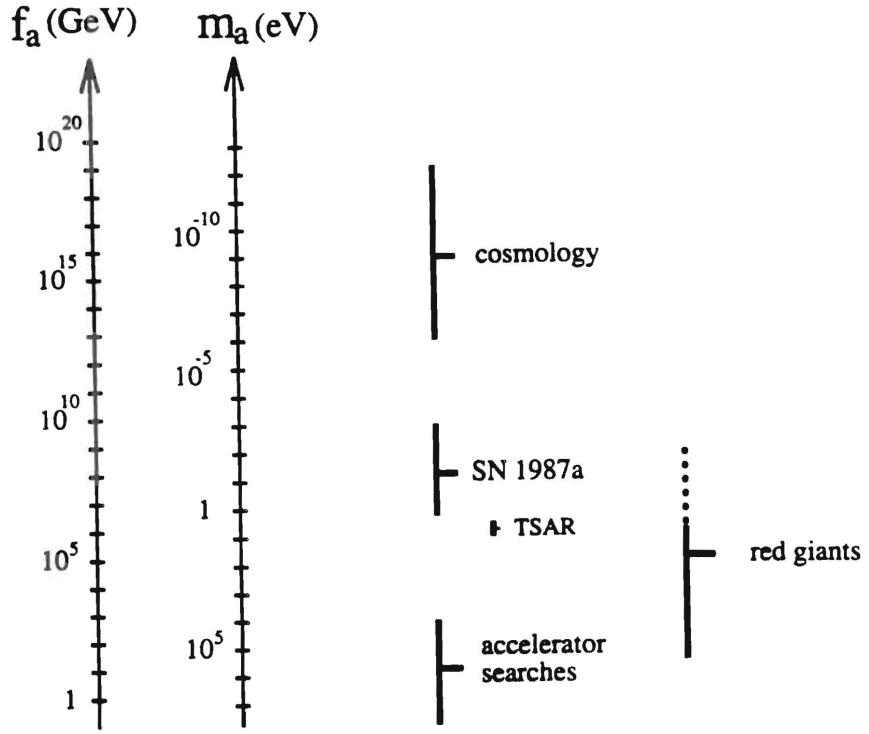

Figure I: Axion mass ranges which have been ruled out so far. TSAR is an acronym for the telescope search by M. A. Bershady et al.[14]

with $g_\gamma = \frac{m_u}{m_d+m_u} \approx 0.36$. In the limit where CP is conserved, the coupling of the axion to a fermion f has the form

$$\mathcal{L}_{a\bar{f}f} = i\, g_f \frac{m_f}{v} a\, \bar{f}\gamma_5 f. \tag{1.15}$$

The coefficients g_f that appear in Eq. (1.15) are rather model-dependent. General formulas for the g_f are given in Refs. (10), including the case where f is a proton or neutron.

When the axion was first proposed, it was thought that the breaking of $U_{PQ}(1)$ occurred at the electroweak scale, i.e. $v \sim 250\ GeV$. The corresponding

axion was searched for in various laboratory experiments but was not found[3]. Soon, however, it was discovered[12] how to construct axion models with arbitrarily large values of v. These were called "invisible" axion models because for $v >> 250\ GeV$, the axion is so weakly coupled that the event rates in the axion search experiments mentioned above are hopelessly small. However such axions are still constrained by astrophysical and cosmological considerations.

Light weakly coupled bosons are severely constrained by stellar evolution because stars emit such particles from their whole volume whereas they emit the more strongly coupled photons (neutrinos in the special case where the 'star' is a supernova core) only from their surface. This idea has been extensively applied[13] to the axion using a variety of stellar objects (the sun, red giants, neutron stars, white dwarfs, and the supernova SN 1987a) and a variety of axion producing thermal processes in those objects (the Primakoff process, Compton-like scattering, axion bremstrahlung in electron-nucleon and nucleon-nucleon scattering, etc.). SN 1987a rules out the axion mass range $10^{-3} eV < m_a <$ few eV. The evolution of red giants rules out $500\ keV < m_a < 10^{-2}\ eV$ for axions with a regular size $[g_e = 0(1)]$ coupling to the electron, and $500\ keV < m_a <$ few eV for axions whose coupling to the electron is much suppressed $[g_e << 1]$. The latter type is usually called "hadronic" axion. It is quite easy to build hadronic axion models (i.e. there is nothing artificial about an axion with $g_e << 1$) and such models can be grand unified. A telescope search for monochromatic photons[14] from the 2γ decay of relic axions ruled out $3eV < m_a < 8\ eV$. Axions with mass $m_a > 10\ keV$ are ruled out by the accelerator based searches mentioned above. By combining the constraints from stellar evolution with those from laboratory searches, one can thus conclude that all axions with $m_a > 10^{-3}\ eV$ are ruled out except possibly for a small window near $m_a = 0(eV)$, for hadronic axions only.

The cosmological bound on the axion mass is of order $m_a \gtrsim 10^{-6}\ eV$. It will be discussed in detail in the next section. There is an additional constraint due to the presence of domain walls in axion models[15,16,17]. The domain walls must be gotten rid of before they dominate the energy density of the universe. This can, in fact, be achieved through inflation or by constructing the axion model in such a way that it has a unique vacuum.

Figure I summarizes the present constraints on the axion.

II. THE COSMOLOGICAL AXION ENERGY DENSITY

There are two main contributions to the present cosmological axion energy density. The first is due to the realignment of the vacuum during the QCD phase transition.[18] The second is from axions radiated by cosmic axion strings.[19,20] Because of space limitations, I will only describe the first contribution here. See refs. [19,20] for a discussion of the second contribution.

At very high temperatures the $U_{PQ}(1)$ symmetry is restored. It becomes spontaneously broken when the temperature drops below $T_{PQ} \approx v$ and the φ-field acquires vacuum expectation value:

$$< \varphi(x) >= v\ e^{i\alpha(x)}\ . \tag{2.1}$$

At these high temperatures, however, the non-perturbative QCD effects (instantons and the like) which will give the axion its mass and thereby lift the degeneracy at the bottom of the "Mexican hat" potential $V(\varphi^\dagger\varphi)$ are highly suppressed. When the axion mass turns on near the QCD phase transition, the axion field $a(x) = v\ \alpha(x)$ starts to oscillate about one of the CP conserving minima of the effective potential $V_{eff}(\alpha)$. The oscillation begins approximately at cosmological time t_1 such that

$$t_1 m_a[T(t_1)] = 0(2\pi) \tag{2.2}$$

where $m_a(T)$ is the temperature dependent axion mass. Soon after time t_1, the axion mass changes sufficiently slowly that the total number of axions in the oscillation of $\alpha(x)$ about the CP conserving minimum is an adiabatic invariant. Using the methods of finite temperature field theory,[21] $T_1 = T(t_1)$ has been estimated to equal about one GeV. The number density of axions at time t_1 due to the initial vacuum misalignment is of order

$$\begin{aligned} n_a^{vac}(t_1) &\approx \rho_a^{vac}(t_1)\frac{1}{m_a(t_1)} \\ &\approx \frac{1}{2}v^2 m_a(t_1) < \alpha^2(t_1) > \approx \pi f_a^2 \frac{1}{t_1} \end{aligned} \tag{2.3}$$

where $f_a = \frac{v}{N}$, as defined earlier. In Eq.(2.3), we have used the fact that the field $a(x) = v\alpha(x)$ is approximately homogeneous on the horizon scale t_1. Wiggles in $\alpha(x)$ which entered the horizon long before t_1 have been red-shifted away[22]. In Eq.(2.3) we have also used Eq.(2.2) and the fact that the initial departure $\alpha(t_1)$ of α from the nearest minimum is of order $\frac{1}{N}$ because N is the number of CP conserving minima[15] at the bottom of the "Mexican hat" potential. The axions of Eq.(2.3) 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 that the present energy density in axions is

$$\rho_a^{vac}(t_0) \approx \rho_{crit}(t_0)\left(\frac{0.6\ 10^{-5} eV}{m_a}\right)^{7/6} \cdot \left(\frac{200\ MeV}{\Lambda_{QCD}}\right)^{3/4}\left(\frac{75\ km\ s^{-1}\ Mpc^{-1}}{H_0}\right)^2 \tag{2.4}$$

where $\rho_{crit}(t_0)$ is the present critical energy density for closing the universe, H_0 is the present value of the Hubble constant and Λ_{QCD} is the QCD scale factor. Eq.(2.4) implies the bound $m_a \gtrsim 0.6\ 10^{-5}\ eV$.

It should be emphasized however that there are many sources of uncertainty in the estimate of the cosmological axion energy density and that Eq.(2.4) only provides us with a rough estimate. For example, the axion energy density can be diluted by the entropy release from heavy particles which decouple before the QCD epoch but decay afterwards[23]. It can also be diluted by the entropy release associated with a first order QCD phase transition. On the other hand, if the QCD phase transition[24] is first order, an abrupt change in the axion mass at the transition may increase $\rho_a(t_0)$ as compared to Eq.(2.4). Finally, if inflation occurs after the Peccei-Quinn phase transition at which $U_{PQ}(1)$ gets spontaneously broken, then

the right-hand side of Eq.(2.4) must be multiplied by $(\alpha(t_1)N)^2$ where $\alpha(t_1)N$ is a randomly chosen number between zero and 0(1). In that case there may be an accidental suppression of the cosmological axion energy density from initial vacuum misalignment because the phase of the Peccei-Quinn field happens to lie close to the CP conserving value at the start of the QCD phase transition. Because of quantum mechanical fluctuations in the axion field during the epoch of inflation[25], the suppression cannot be perfect. The amount of accidental suppression allowed depends upon the inflationary model. See refs. [26] for a discussion.

The axions produced during the QCD phase transition when the axion mass turns on are "cold dark matter" because these axions are non-relativistic from the moment of their first appearance at 1 GeV temperature[18,27]. For the past decade, the 'standard model' for the formation of structure in the Universe[28] has been based, in fact, on the assumptions of Cold Dark Matter (CDM), a flat (Zel'dovich-Harrison) spectrum of primordial adiabatic Gaussian density perturbations and some 'biasing' in the extent to which light traces matter. Hence, there is excellent motivation to search for CDM as the constituent particles of our galactic halo, and in particular to search for dark matter axions as described in the next section. The standard CDM model, based on the assumptions stated above, may require modification in light of some recent measurements, namely the COBE detection of fluctuations in the microwave background on a 10° angular scale and data on galaxy-galaxy correlations at large angles. These are difficult to reconcile with measurements of peculiar motions and galactic rotation curves within the standard CDM model.[29] Several authors[30] have proposed that it be modified by mixing Hot and Cold Dark Matter, $\Omega_{CDM} \simeq 0.7$ and $\Omega_{HDM} \simeq 0.3$ being the preferred ratio, and the leading HDM candidate being a neutrino weighing about 10 eV. From the point of view of the axion search described in the next section and of the other CDM searches as well, it is worth pointing out that in this mixed neutrino + Cold Dark Matter scenario, the neutrino mass density fraction in our galactic halo is likely to be much less than the neutrino fraction in the cosmological energy density.[31] For example, if $H_0 = 55 km/sec.Mpc$, then $\Omega_\nu = 0.3$ for m_{ν_e}, $m_{\nu_\mu} << m_{\nu_\tau} \simeq 9eV$. However the density fraction in our galactic halo $\rho_\nu/\rho_{CDM}|_{halo} < 0.04$ for those parameters.

E. W. Kolb and I. I. Tkachev[32] have recently pointed out that, in the case where there is no inflation after the Peccei-Quinn phase transition, the inhomogeneities in the axion field may, at the start of the matter dominated epoch, evolve into compact objects called "axion miniclusters". From the point of view of the galactic halo axion search described in the next section, it is of course very important to try and determine the fraction of cosmic axions that condense into these miniclusters.

III. CAVITY DETECTOR OF GALACTIC HALO AXIONS

As we saw in the preceding section, it is conceivable that the earth is bathed in a sea of galactic halo axions. If the galactic halo is made up exclusively of axions, their density in the solar neighborhood[33] is approximately $\frac{1}{2}10^{-24}\ gr/cm^3$ and their velocity dispersion is approximately 10^{-3} times the speed of light. A likely mass

range for such axions is ($\hbar = c = 1$)

$$2\pi(242\ MH_z) = 10^{-6}\ eV \lesssim m_a \lesssim 10^{-4}\ eV = 2\pi(24.2\ GH_z)\ . \tag{3.1}$$

The possibility of detecting these dark matter axions constitutes an exciting but nonetheless realistic prospect.

Indeed, axions can be searched for by stimulating their conversion to photons in a strong magnetic field[34]. The relevant coupling is given in Eq.(1.12). In particular, an electromagnetic cavity permeated by a strong magnetic field can be used to detect galactic halo axions. The latter have velocities β of order 10^{-3} and hence their energies

$$\epsilon_a = m_a + \frac{1}{2} m_a \beta^2 \tag{3.2}$$

have a spread of order 10^{-6} above the axion mass. Consider a cylindrical electromagnetic cavity of arbitrary cross-sectional shape, permeated by a large static approximately homogeneous longitudinal magnetic field $\vec{B} = B_0 \hat{z}$. When the frequency $\omega = 2\pi f$ of an appropriate cavity mode equals m_a, galactic halo axions can convert to quanta of excitation (photons) of that cavity mode. Only the $TM_{n\ell 0}$ modes couple in the limit where the cavity is much smaller than the de Broglie wavelength $\lambda_a = 2\pi(\beta m_a)^{-1} \approx 2\pi\ 10^3\ m_a^{-1}$ of the galactic halo axions. The power on resonance from axion $\rightarrow$ photon conversion into the $TM_{n\ell 0}$ mode is[34,35,36]

$$\begin{aligned} P_{n\ell} &= \left(\frac{\alpha}{\pi} g_\gamma \frac{N}{v}\right)^2 V B_0^2 \rho_a C_{n\ell} \frac{1}{m_a} Min(Q_L, Q_a) \\ &= 2.10^{-26} Watt \left(\frac{V}{500\ liter}\right) \left(\frac{B_0}{8\ Tesla}\right)^2 C_{n\ell} \left(\frac{g_\gamma}{0.36}\right)^2 \left(\frac{\rho_a}{\frac{1}{2}10^{-24} gr/cm^3}\right) \cdot \\ &\cdot \left[\frac{m_a}{2\pi(3GH_z)}\right] Min\ (Q_L, Q_a)\ . \end{aligned} \tag{3.3}$$

where V is the volume of the cavity, ρ_a is the density of galactic halo axions on earth, Q_L is the quality factor of the cavity and $Q_a = 10^6$ is the "quality factor" of the galactic halo axion signal, i.e. the ratio of their energy to their energy spread. $C_{n\ell}$ is a model dependent form factor defined by:

$$C_{n\ell} = \frac{|\int_V d^3x \vec{E}_{n\ell} \cdot \vec{B}_0|^2}{B_0^2 V \int_V d^3x\ |\vec{E}_{n\ell}|^2} \tag{3.4}$$

where $\vec{B}_0(\vec{x})$ is the static magnetic field and $\vec{E}_{n\ell} e^{iwt}$ is the oscillating electric field. For a cavity of rectangular cross-section

$$\begin{aligned} C_{n\ell} &= \frac{64}{\pi^4 n^2 \ell^2} \qquad for\ n\ and\ \ell\ odd \\ &= 0 \qquad otherwise\ . \end{aligned} \tag{3.5}$$

For a circular cross-section

$$C_{nm} = \frac{4}{(X_{on})^2} \delta_{om} \tag{3.6}$$

where X_{on} is the nth zero of the Bessel function $J_0(x)$. Eqs. (3.5) and (3.6) show that one should use the lowest TM mode if at all possible. From now on, we will suppress the indices n, ℓ.

To detect the power P, a hole must be made in the cavity wall through which the electromagnetic radiation can be brought to the front end of a microwave receiver. The quality factor Q_L which appears in Eq. (3.3) is the loaded quality factor given by

$$\frac{1}{Q_L} = \frac{1}{Q_w} + \frac{1}{Q_h} \tag{3.7}$$

where $\frac{1}{Q_w}$ is the contribution due to absorption into the cavity walls and $\frac{1}{Q_h}$ is the contribution from the hole. The maximum power that can be brought to the front end of the microwave receiver is $\frac{Q_L}{Q_h} P$.

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 the axion signal. The cavity can be tuned by moving inside the cavity a dielectric rod or metal post. Using Eq. (3.3), one finds that to obtain a given signal to noise ratio s/n, the search rate is

$$\begin{aligned} \frac{df}{dt} &= \frac{2.6GHz}{year}\left(\frac{3n}{s}\right)^2\left(\frac{V}{500\ liter}\right)^2\left(\frac{B_0}{8\ Tesla}\right)^4 C^2\left(\frac{g_\gamma}{0.36}\right)^4 \cdot \\ &\cdot \left(\frac{\rho_a}{\frac{1}{2}10^{-24}gr/cm^3}\right)^2\left(\frac{20K}{T_n}\right)^2\left(\frac{f}{3GHz}\right)^2 \cdot \begin{cases} Q_w/Q_a & \text{if } Q_w < 3Q_a \\ \frac{27}{4}(1-\frac{Q_a}{Q_w})^2 & \text{if } Q_w > 3Q_a \end{cases} \end{aligned} \tag{3.8}$$

where T_n is the sum of the physical temperature of the cavity plus the noise temperature of the microwave receiver. Eq. (3.8) assumes

1. that when $Q_L < Q_a$, i.e. when the cavity bandwidth is larger than the axion bandwidth, one uses the possibility of looking at Q_a/Q_L axion bandwidths simultaneously,
2. that Q_h has been adjusted so as to maximize the search rate. For $Q_w < 3Q_a$, the optimal $Q_h = \frac{1}{2}Q_w$ (and hence $Q_L = \frac{1}{3}Q_w$) whereas for $Q_w > 3Q_a$, the optimal Q_h is such that $Q_L = Q_a$.

Actually, the best possible quality factors attainable at present, using oxygen free copper, are only of order 10^5 in the GHz range. (It does not help to make the cavity of superconducting material since it is permeated by a strong magnetic field in the experiment.) The factor Q_w/Q_a on the RHS of Eq.(3.8) is therefore of order 10^{-1}. On the other hand, noise temperatures as low as 3K (in the 1.2 - 1.6 GHz range) have been reached using commercially available microwave receivers[37].

Eq. (3.8) shows that a galactic halo search is feasible with presently available technology, provided the form factor C can be kept at values of order one at all frequencies. This raises the following issue. A large volume empty cylindrical cavity has a low resonant frequency in its lowest TM mode: $f = 0.115$ GHz $(\frac{1m}{R})$ where R is the radius of the cavity. Thus a large cylindrical cavity is convenient for searching the low frequency end of the range (3.1). How does one search the higher frequencies? This question was addressed in a paper by C. Hagmann et al[38] which also discusses various cavity tuning schemes and presents the results of computer

simulations to optimize cavity design. It was found that the best way to tune a cavity in this experiment is by translating a dielectric rod or metal post sideways inside the cavity. Secondly, the paper concludes that the best way to extend the search to high frequencies is by power-combining many identical cavities which fill up the volume inside the magnet bore. This method avoids the problems of mode localization and resonance crowding which plague the other approaches that were considered. It allows one to maintain $C = 0(1)$ at all frequencies, albeit at the cost of increasing engineering complexity with increasing frequency.

Galactic halo axion searches using cavity detectors have been carried out at Brookhaven National Laboratory by a collaboration between Rochester, Brookhaven and Fermilab[39] (RBF) and at the University of Florida[37] (UF). Development work was also done at KEK[40]. In addition, the possibility of using a beam of Rydberg atoms to detect the microwave photons from a cavity galactic halo axion detector was discussed.[41]

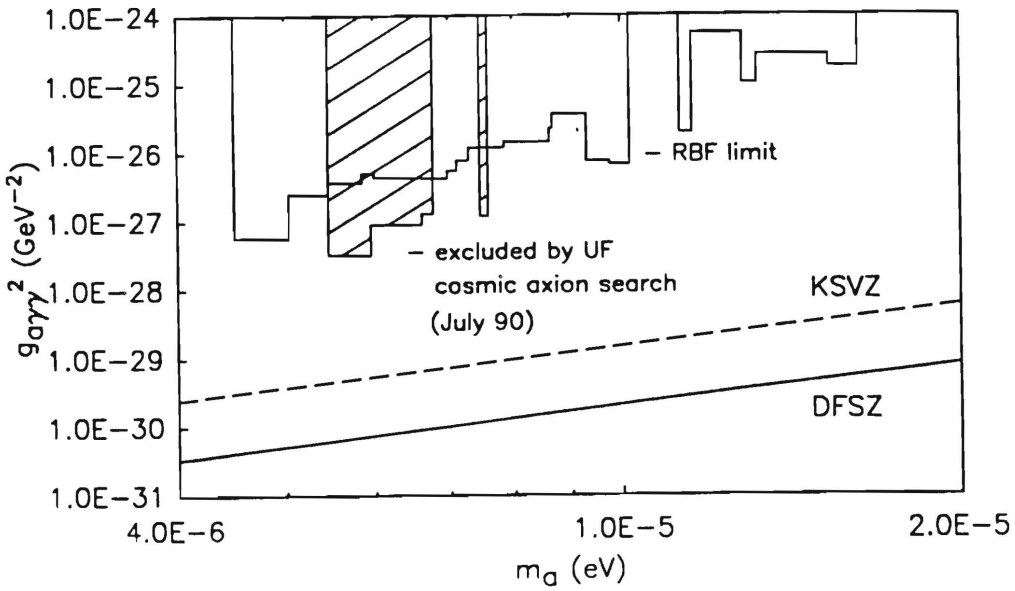

Figure II: Upper limits on the coupling of the axion to two photons. The shaded areas are ruled out if the galactic halo is made of axions.

Fig. II shows the limits that the RBF and UF collaborations placed on the square of the coupling $g_{a\gamma\gamma} \equiv \frac{g_\gamma \alpha}{\pi f_a}$ as a function of the axion mass assuming that the galactic halo is made of axions. The (magnetic field)^2 x volume provided by the magnets used in the experiments was $B_0^2 V = 0.36\ T^2 m^3$ for RBF and $B_0^2 V = 0.45\ T^2 m^3$ for UF. The improvement of a factor 5 or 10 in sensitivity of the UF detector over the Brookhaven one is due, for the most part, to the adoption of a more efficient, computerized data taking method and to the use of more sensitive microwave equipment. Even so, the sensitivity of the UF experiment is still a factor

500 or so short of that required to detect galactic halo axions in the DFSZ model. Note that in other models, such as the KSVZ model, the signal predicted is higher.

Presently a second generation cavity detector is under construction at Lawrence Livermore National Laboratory[42] by a collaboration of physicists from LLNL, Univ. of Florida, The Institute for Nuclear Research in Moscow, MIT, LBL, UC Berkeley and FNAL. The experiment will utilize a commercial magnet with a central field of 8T. The volume for the cavity will have an inner diameter of 50 cm and a length of 110 cm. Hence, $B_0^2 V = 12T^2 M^3$ which is approximately a factor 25 larger than the $B_0^2 V$ values of the RBF and UF pilot experiments. By achieving in addition an improvement in noise temperature, the LLNL experiment is expected to reach a sensitivity which is a factor 40 better than that of the UF experiment. The LLNL experiment will also be the first to use arrays of multiple cavities to expand widely the mass range searched. It will cover $1.5 < m_a < 12.6 \mu eV$. It is planned to start taking data in summer '94, and run for approximately three years.

It appears at present that the cavity detectors cannot reach axion masses much larger than $10^{-5} eV$ because of the complexities involved in segmenting a given magnetic volume into many small cavities. In a recent paper,[43] D. Tanner, Y. Wang and I studied an experimental scheme to search for galactic halo axions in the $10^{-5} eV$ to $10^{-3} eV$ mass range. The detector consists of a large number of parallel superconducting wires embedded in a material transparent to microwave radiation. The wires carry a current configuration which produces a static, inhomogeneous magnetic field $\vec{B}_0(\vec{x})$ within the detector volume. Axions which enter this detector volume have a small probability to convert to photons. Ref. [43] discusses the feasibility of the detector and its sensitivity.

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