DOE/ER/40272--171

DE92 040501

Institute for Fundamental Theory Preprint UFIFT

DARK MATTER AXIONS*

Pierre Sikivie

Institute for Theoretical Physics, University of California, Santa Barbara, CA 93106
and
Department of Physics, University of Florida, Gainesville, Florida 32611

Contents:

* To be published in the proceedings of the IFT Workshop on Dark Matter, held in Gainesville, Florida, on February 14-16, 1992.

DARK MATTER AXIONS*

Pierre Sikivie

Institute for Theoretical Physics, University of California, Santa Barbara, CA 93106
and
Department of Physics, University of Florida, Gainesville, Florida 32611

Contents:

I. INTRODUCTION TO AXION PHYSICS

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

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

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]:

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

* To be published in the proceedings of the IFT Workshop on Dark Matter, held in Gainesville, Florida, on February 14-16, 1992.

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 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 $\bar{\theta}$,
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

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

$\mathcal{L}_{PQ}$ 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, Georgi and Randall[9] found $\bar{\theta} \approx 10^{-17}$ for the Kobayashi-Maskawa model of CP violation.

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

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

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

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

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.

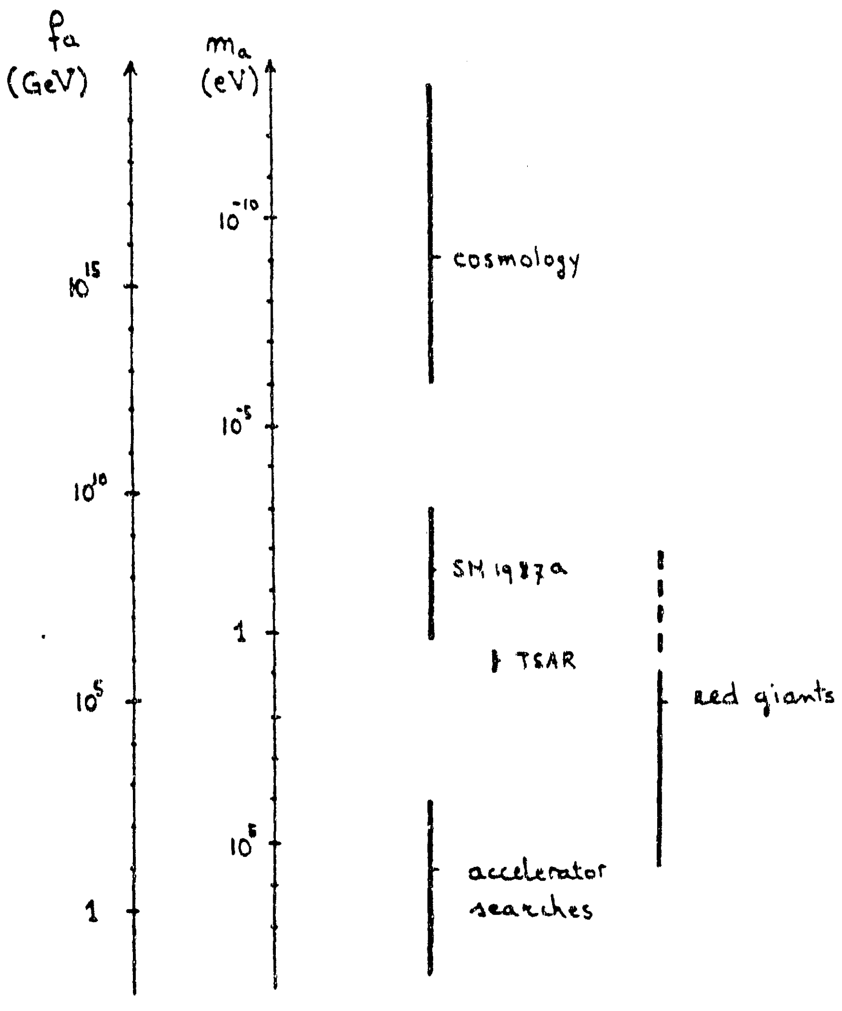

Figure I: Axion mass ranges which have been ruled out so far. TSAR is an acronym for the telescope search by M. A. Bershady et al.[15]

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[13,14] 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[14] 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.

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. However a recent telescope search for monochromatic photons[15] from the 2γ decay of relic axions rules out $3eV < m_a < 8\ eV$.

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[16,17,18]. 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 1 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. The second is from axions radiated by cosmic axion strings. We will discuss the two contributions in succession.

a. The contribution from initial vacuum misalignment[19]

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 finite temperature instanton calculations of Gross, Pisarski and Yaffe[20], $T_1 = T(t_1)$ has been estimated to equal about one GeV[19]. 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[21]. 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[16] 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[22]. 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[23] 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 before the start of the QCD phase transition. It was recently emphasized that because of quantum mechanical fluctuations during the epoch of inflation[24], the suppression cannot be perfect. The amount of accidental suppression allowed depends upon the inflationary model.

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[19,25]. Cold dark matter with a flat (Zel'dovich-Harrison) spectrum of primordial density perturbations and some "biasing" in the extent to which light traces matter yields at present a popular and apparently viable scenario of galaxy formation. Thus the dark matter clustered in halos around galaxies could be axions with mass of order 10^{-5} eV.

b. The contribution from cosmic axion strings

Cosmic axion strings make another sizable contribution to the axion cosmological energy density if inflation does not occur between the Peccei-Quinn phase transition where $U_{PQ}(1)$ becomes spontaneously broken and the QCD epoch[26,27,28]. An axion string is a topological knot in the vacuum expectation value (2.1). To describe it more fully consider the action density for the complex scalar field whose VEV spontaneously breaks the $U_{PQ}(1)$ symmetry:

$$S = \int d^4x \left[\frac{1}{2}\partial_\mu\varphi^\dagger\partial^\mu\varphi - \frac{\lambda}{4}(\varphi^\dagger\varphi - v^2)^2\right] . \tag{2.5}$$

The configuration corresponding to a straight axion string along the $\hat{z}$-axis is

$$\varphi = v \; f(\mu\rho)e^{i\theta} \tag{2.6}$$

in cylindrical coordinates (z, ρ, θ). $\mu \equiv \sqrt{\lambda}\, v$ will be identified with the inverse of the core size of the string. $f(\mu\rho)$ is a function which goes to zero when $\mu\rho \to 0$, approaches one when $\mu\rho >> 1$, and which minimizes the energy per unit length of the string. The latter is then

$$\begin{aligned} \tau &= \int d^2x \left[\frac{1}{2} \mid \vec{\nabla}\varphi \mid^2 + \frac{\lambda}{4}(\varphi^\dagger\varphi - v^2)^2\right] \\ &\approx \int_{\rho>\mu} d^2x \frac{v^2}{2} \mid \vec{\nabla} e^{i\theta} \mid^2 = \pi\, v^2\, \ell n\, \mu L \end{aligned} \tag{2.7}$$

where L is an infra-red cutoff. The RHS of Eq. (2.7) neglects the contribution to τ from the string core. That contribution is of order $\pi\, v^2$, which is smaller than the contribution from the Nambu-Goldstone field outside the string core by a factor $1/\ell n\, \mu L$. For an axion string in the early universe, the infra-red cutoff is provided by the presence of neighboring axion strings with roughly opposite direction. L is then of order the average distance between strings which is of order the horizon scale (see below). For cosmic axion strings at the QCD epoch, we have $\ell n\, \mu L \approx 70$.

Cosmic axion strings dissipate their energy by radiating axions. This is a very efficient process. As a result, the number density of axion strings in the early universe is expected to be of order the lowest possible consistent with causality, i.e. approximately one long string per horizon. The energy density in axion strings at time t is thus of order

$$\rho_s(t) \approx \frac{\tau}{t^2} = \frac{\pi\, v^2}{t^2} \ell n\, \mu t\, . \tag{2.8}$$

Using energy conservation, one obtains from (2.8) the number density of axions radiated by cosmic axion strings[27]:

$$n_a^{str}(t) \approx \frac{1}{t^{3/2}} \int_{t_{PQ}}^{t} \frac{dt'}{t'^{3/2}} \frac{\tau(t')}{\omega(t')} , \tag{2.9}$$

where

$$\frac{1}{\omega(t)} = \int \frac{dE}{d\omega}(t)\frac{d\omega}{\omega} \Big/ \int \frac{dE}{d\omega}(t) d\omega \tag{2.10}$$

and $\frac{dE}{d\omega}(t)$ is the energy spectrum of axions radiated at time t. To obtain the present axion cosmological energy density due to the decay of cosmic axion strings, we need to determine the number density of such axions at time t_1 defined by Eq. (2.2). At time t_1 the axions acquire mass and the axion strings become the boundaries of axion domain walls. What happens to the domain walls is a story told elsewhere[16,17,18]. As Eqs.(2.9) and (2.10) imply, to obtain the number density of radiated axions we need to know their spectrum.

The axions that are radiated at time t are emitted by cosmic axion strings which are bent over a distance scale of order t and which are relaxing to lower energy configurations. There is at present disagreement in the literature about the behavior and the energy spectrum of such a string. One view[26] is that the axion string oscillates 10 to 20 times before it has reached its minimum energy configuration and that the spectrum of radiated axions is concentrated near $\frac{2\pi}{t}$. Let us call this case A. The other view, which my collaborators and I have argued in favor of[27,28], is that, in the limit of large $\ell n\ \mu t$, the axion string will straighten itself out at once (without large scale oscillations) and that the energy spectrum of radiated axions is $\frac{dE}{dk} \sim 1/k$ with a high energy cutoff of order μ and a low energy cutoff of order $\frac{2\pi}{t}$. Recent computer simulations[28] lend support to this second view. Let us call it case B. In case B, $\omega(t) \approx \frac{2\pi}{t}\ell n\ \mu t$ and hence the number density of radiated axions at time t_1 is

$$n_a^{str}(t_1) \approx \frac{v^2}{t_1} \ . \tag{2.11}$$

In case A, on the other hand, one has $\omega(t) \approx \frac{2\pi}{t}$ and

$$n_a^{str}(t_1) \approx \frac{v^2}{t_1}\ell n\ \mu t_1 . \tag{2.12}$$

Most of the axions radiated by cosmic axion strings are non-relativistic today. Hence we may directly compare $n_a^{str}(t_1)$ with $n_a^{vac}(t_1)$ [Eq. (2.3)] to determine the relative importance of axions radiated by strings versus axions produced by initial vacuum misalignment.

Let us assume first that N=1. To some extent, N=1 is favored over N$\neq$ 1 because N=1 axion models have a particularly straightforward mechanism to rid the early universe of axion domain walls. This is an important consideration since we are assuming, for cosmic axion strings to be relevant at all, that there is no inflation after the Peccei-Quinn phase transition and hence the axion domain walls cannot be gotten rid of through inflation. In case A, one would conclude that the contribution to the cosmological axion energy density from axions radiated by cosmic axion strings is roughly a factor $\ell n(\mu t_1) \simeq 70$ times larger than the contribution from initial vacuum misalignment [compare Eqs.(2.3) and (2.12)]. In that case there would be little room left between the cosmological bound and the bound from the supernova. In case B, on the other hand, the cosmological axion energy density from axions radiated by cosmic axion strings is of the same order of magnitude as the contribution from initial vacuum misalignment [compare Eqs.(2.3) and (2.11)]. In that case, the cosmological bound on the axion mass remains $m_a \gtrsim 10^{-6}\ eV$.

Note however that N$\neq$ 1 is also a logical possibility. One can rid the early universe of axion domain walls in this case by introducing a tiny explicit breaking of the $U_{PQ}(1)$ symmetry which slightly lowers one of the N vacua with respect to the others[16]. A small bias of this kind can eliminate the domain walls from the early universe yet be compatible with the Peccei-Quinn solution to the strong CP problem. Comparison of Eqs.(2.3) and (2.11) shows that if N is large, the cosmological axion energy density from cosmic axion strings will dominate over that from initial vacuum misalignment even in case B. However, when $N \neq 1$, there is also a contribution to the cosmological axion energy density from axions produced in the decay of axion domain walls, and this is likely to be the dominant contribution. Its size depends upon the small amount of explicit breaking of the $U_{PQ}(1)$ symmetry introduced to eliminate the domain walls from the universe before they dominate the energy density.

As a conclusion to section II on the cosmological axion energy density, let us state that $\Omega_a < 1$ implies $m_a \gtrsim 10^{-6} eV$ in the sense that $m_a < 10^{-6} eV$ requires that there is inflation after the phase transition during which $U_{PQ}(1)$ gets spontaneously broken and that, by accident, the axion field happens to lie close to the CP conserving value of the effective potential even before the axion mass is turned on.

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[29] 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 ($\hbar = c = 1$)

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

The possibility of detecting these dark matter axions constitutes an exciting but, as we will see, realistic prospect.

Indeed, axions can be searched for by stimulating their conversion to photons in a strong magnetic field[30]. 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[30,31,32]

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

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:

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

where $\vec{B}_0(\vec{x})$ is the static magnetic field and $\vec{E}_{n\ell} e^{i\omega t}$ 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 n^{th} 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[33].

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[34] 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[35] (RBF) and at the University of Florida[34] (UF). Development work was also done at KEK[36]. 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 by S. Matsuki and K. Yamamoto[37].

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.

Recently, a proposal has been put forth to use two large 'Axicell' magnets from the decomissioned mirror fusion test facility at Lawrence Livermore National Laboratory for a cavity galactic halo axion detector[38]. The two magnets together yield $B_0^2 V \simeq 137\ T^2 m^3$ which is an improvement of a factor 300 or so over the

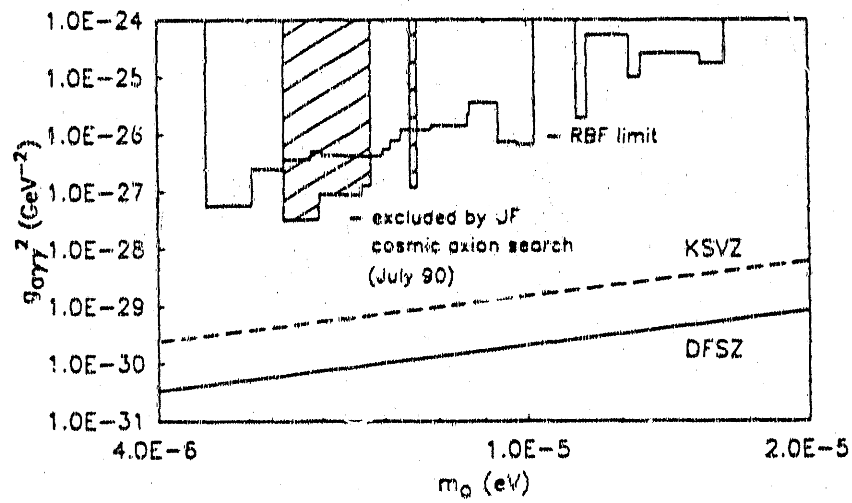

Figure II: Upper limits on the coupling of the axion to two photons. The shaded areas are ruled out if the galactic halo is made of axions.

first generation experiments. The proposed detector would have the sensitivity required to detect galactic halo axions in the DFSZ model over the mass range $0.6 < m_a < 5.6\mu eV$ in Phase I of the experiment and $5.6 < m_a < 16\mu eV$ in Phase II.

IV. A PROPOSAL FOR AXION DETECTION IN THE $m_a \sim 10^{-3}$eV RANGE.

It is possible, as implied by our discussion of the cosmological abundance of axions in section II, that the dark matter is made of axions with mass anywhere in the range $10^{-3}eV < m_a < 10^{-7}eV$. We argued that $m_a \sim 10^{-5}eV$ is the most likely value for which $\Omega_a = 1$, but this value is affected by large uncertainties. In particular, if there is inflation (no inflation) after the PQ phase transition, the most likely value of m_a tends to be shifted downward (upward). It is of course desirable to search as much axion mass range as possible without excessive regard for theoretical input. The cavity detector of galactic halo axions is the easiest to build. However it does not appear at present that cavity detectors can cover the whole $10^{-3}eV - 10^{-7}eV$ mass window. Their range is limited on the one hand by the spatial dimensions of the magnet (bore radius $R < 3m$ implies $m_a \gtrsim 1.6\ 10^{-7}eV$) and, on the other hand, by the complexities involved in segmenting a given magnetic volume into many small cavities to look for high mass axions. The most complex system which has been envisaged up till now is 1024 cavities in the $3m^3$ volume ($R \simeq 0.7m, L = 2m$) afforded by the two large LLNL 'Axicell' magnets mentioned in the previous section. This would reach $m_a \simeq 1.6\ 10^{-5}eV$. It seems difficult to reach much larger axion masses using the cavity detector assuming only presently available technology. An alternative approach, specifically to explore the larger axion masses, is desirable.

In this section we elaborate on a scheme proposed earlier[30,31,39] to search for galactic halo axions in the $m_a \sim 10^{-3}eV$ range. Because of the coupling (1.12) of the axion to two photons, axions will convert to photons (and vice-versa) in an externally applied magnetic field. The cross-section for $a \rightarrow \gamma$ conversion in a region of volume V permeated by a static magnetic field $\vec{B}_0(\vec{x})$ is

$$\sigma = \frac{1}{16\pi^2\beta_a}\left(\frac{\alpha g_\gamma}{\pi f_a}\right)^2 \int d^3k_\gamma \delta(E_a - w)\left|\int_V d^3x e^{i(\vec{k}_\gamma - \vec{k}_a)\cdot\vec{x}}\vec{n} \times \vec{B}_0(\vec{x})\right|^2, \tag{4.1}$$

where $(E_a, \vec{k}_a) = E_a(1, \vec{\beta}_a)$ is the 4-momentum of the axion and $(w, \vec{k}_\gamma) = w(1, \vec{n})$ is the 4-momentum of the photon. $E_a = w$ is required because the magnetic field is static. The extra momentum $\vec{q} = \vec{k}_\gamma - \vec{k}_a$, which is necessary because the photon is massless whereas the axion is massive, must be provided by the inhomogeneity of the magnetic field. For galactic halos axion detection, $k_a \sim 10^{-3}m_a$ and therefore the magnetic field should be made inhomogeneous on the length scale m_a^{-1}.

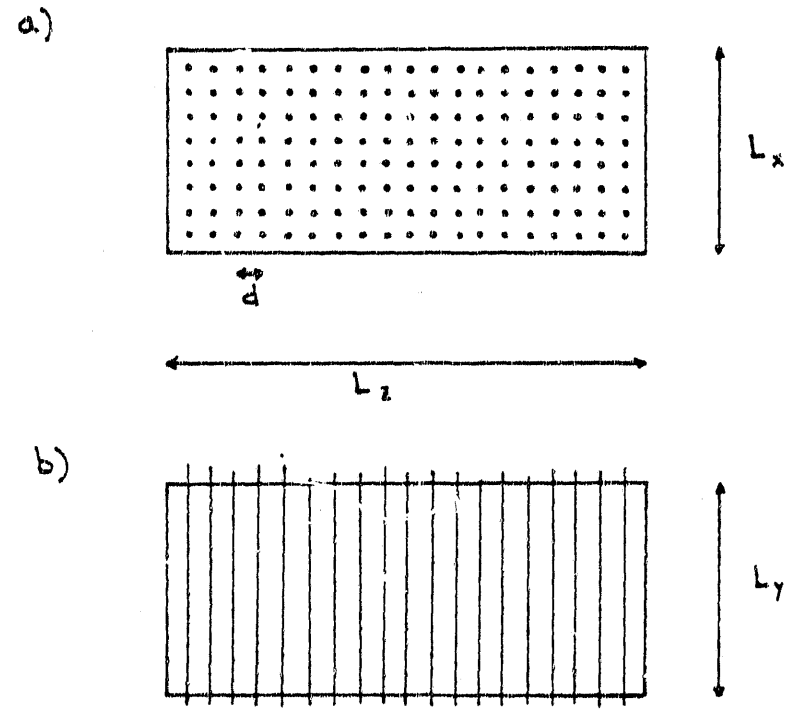

Figure III. Top (a) and side view (b) of the detector of galactic halo axions described in section IV. In an actual design L/d would be much larger than represented here.

The proposed detector is illustrated in Fig. III. A large number of parallel superconducting wires are embedded into a material transparent to microwave radiation. The material keeps the wires from moving. Let $\hat{y}$ be the common direction of the wires. The intersections of the wires with the (x, z) plane form an array of unit cell size d. d must be smaller than the inverse of the axion mass which is being searched for. Through each wire is circulated a current $I(z)$ whose intensity depends only upon the z coordinate of the wire. This produces a coarse grained current density $\vec{J} = \hat{y} I(z) d^{-2}$. To obtain the static inhomogeneous magnetic field

$$\vec{B}_0(\vec{x}) = \hat{x} B_0 \cos(qz), \tag{4.2}$$

one sets the individual currents at $I(z) = -qB_0 \sin(qz)$. B_0 can be, in principle, as large as the critical magnetic field to dri e the superconductor normal. If the magnetic field (4.2) extends over a rectangular volume $V = L_x L_y L_z$, then the conversion cross-section (4.1) is

$$\sigma = \frac{1}{4\beta_a}\left(\frac{\alpha g_\gamma}{\pi f_a}\right)^2 B_0^2 L_x L_y \left[\left(\frac{\sin[m_a(1-\beta_{az})-q]\frac{L_z}{2}}{m_a(1-\beta_{az})-q}\right)^2 + \left(\frac{\sin[m_a(1+\beta_{az})-q]\frac{L_z}{2}}{m_a(1+\beta_{az})-q}\right)^2\right] \tag{4.3}$$

provided $\beta_a \ll 1$ and $L_i \gg m_a^{-1}$ $(i = 1,2,3)$. Eq. (4.3) shows that the bandwidth of the detector in the 3$^{\rm d}$ component of momentum is $\Delta k_z^d \simeq \frac{2}{L_z}$. On the other hand, the width of the axion signal in the same variable is $\Delta k_z^a \simeq 2.10^{-3} m_a$. The power into the detector from $a \to \gamma$ conversion on resonance ($q \simeq m_a$), provided the axion signal falls entirely within the bandwidth of the detector $\left(|q - m_a \pm 10^{-3} m_a| < \frac{1}{L_z}\right)$, is

$$\begin{aligned} P = \sigma\beta_a\rho_a &= \frac{1}{8}\left(\frac{\alpha g_\gamma}{\pi f_a}\right)^2 V L_z B_0^2 \rho_a \\ &= 2.10^{-21}\,\text{Watt}\left(\frac{V L_z}{m^4}\right)\left(\frac{B_0}{8T}\right)^2\left(\frac{m_a}{10^{-3}eV}\right)^2\left(\frac{\rho_a}{\frac{1}{2}10^{-24}gr/cm^3}\right) \end{aligned} \tag{4.4}$$

This power must be collected and brought to the front end of a microwave receiver. Note that the photons all have momentum very nearly parallel to the z-axis and hence can be focussed by a parabolic mirror. Also, their polarization ($\vec{E}$ vector direction) is perpendicular to the superconducting wires. This is desirable because it suppresses the scattering of the photons by the wires. Let ζ be the efficiency with which the power P can be brought to the front end of the microwave receiver and let T be the total (physical plus electronic) noise temperature of the receiver. The signal to noise ratio over the (frequency) bandwidth $\Delta f_a = 10^{-6}\left(\frac{m_a}{2\pi}\right)$ of the axion signal is

$$\frac{s}{n} = \frac{\zeta P}{T}\sqrt{\frac{t}{\Delta f_a}} \tag{4.5}$$

where t is the measurement integration time. The search can be carried out over the whole (frequency) bandwidth $\Delta f_d \simeq \frac{2}{L_z}$ of the detector simultaneously. Thus the search rate for a given signal to noise ratio is

$$\begin{aligned} \frac{\Delta f_d}{t} &= \frac{\Delta f_a}{t}\frac{\Delta f_d}{\Delta f_a} = \left(\frac{n}{s}\right)^2\left(\frac{\zeta P}{T}\right)^2\frac{2Q_a}{m_a L_z} \\ &= \frac{3.5\ 10^{-5}\ eV}{\text{year}}\left(\frac{V^2 L_z}{m^7}\right)\left(\frac{B_0}{8T}\right)^4\left(\frac{\rho_a}{\frac{1}{2}10^{-24}gr/cm^3}\right)^2 \times \\ &\times\left(\frac{\zeta}{0.3}\right)^2\left(\frac{10K}{T}\right)^2\left(\frac{m_a}{10^{-3}eV}\right)^3\left(\frac{n}{4s}\right)^2\left(\frac{g_\gamma}{0.36}\right)^4 . \end{aligned} \tag{4.6}$$

To search at a reasonable rate (one factor of 2 per year, say) near $m_a = 10^{-3}eV$, a detector of linear dimensions at least a couple of meters is required. If $L_x = L_y = L_z = 2m$ and $d = 0.2mm$, there are 10^8 wires arranged in 10^4 planes of 10^4 wires each. All wires in the same plane carry the same current. To extend the search to smaller axion masses keeping a reasonably large search rate, successively larger ($L \sim m_a^{-2/7}$) but less fine grained ($d \sim m_a^{-1}$) would have to be constructed. The detector appropriate for a search near $m_a = 10^{-5}eV$ would be roughly $10m$ in linear dimension and have of order 10^5 wires spaced about $2cm$ apart.

V. THE PHASE SPACE STRUCTURE OF COLD DARK MATTER HALOS

The question naturally presents itself what can be learned from a galactic halo signal if such a signal is ever discovered. The modulation of the signal by the rotation and orbital motion of the Earth has been discussed by M. Turner.[40] However, because the cavity detector of galactic halo axions has unusually high energy resolution additional information may be obtainable.

In discussions of dark matter detection on Earth, the dark matter particles are usually assumed to have an isothermal distribution or an adiabatic deformation of an isothermal distribution.[41] There is good reason to make this assumption as a first approximation. Probably the most direct argument is that it predicts the rotation curves of galaxies to be flat at large radial distances. Indeed, it is natural to make the simplest assumption that explains the main feature of the observational data. From a theoretical viewpoint, thermalization is thought to be the result of a period of "violent relaxation"[42] following the collapse of the protogalaxy. If it is strictly true that the energy and velocity distribution of the dark matter particles is isothermal, then the only information that can be gained from its observation is the corresponding virial velocity and our own velocity relative to its standard of rest. If, on the other hand, the thermalization is incomplete, a signal in a dark matter detector may yield additional information. Thus the question to what extent dark matter is thermalized in a galactic halo deserves serious consideration. J. Ipser and I[43] have looked at this issue and found that there are large deviations from a thermal distribution in that the highest energy particles have discrete values of velocity.

To see why, note that the initial (at a time before the start of galaxy formation) velocity dispersion of cold dark matter particles is very small. The typical initial axion velocity is $v_a(t_G) \sim 10^{-17}\left(\frac{10^{-5}eV}{m_a}\right)$ in the case where there is no inflation after the phase transition during which the $U_{PQ}(1)$ quasi-symmetry gets spontaneously broken. If there is inflation after that phase transition, the initial velocity dispersion of axions is smaller still. Thus the axions are in a thin sheet in $(\vec{r}, \dot{\vec{r}})$ phase space. Initially, before the Galaxy has started to form, this sheet lies near $\dot{\vec{r}} = H(t_G)\vec{r}$. After the Galaxy has formed, the sheet is wrapping itself up. This process is illustrated in Figure IV. In the idealized case where all structure in the density distribution on all scales less than that ($\sim 10^{23}cm$) of the Galaxy as a whole is absent, where the gravitational potential of the Galaxy is spherically symmetric and where all dark matter particles move on purely radial orbits, the number of sheets at our location, which is also the number of peaks in the energy-momentum spectrum of dark matter particles measured on earth, is of order 200. We investigated how this sheet structure is affected by the angular momentum of the dark matter particles, the gravitational field of the galactic disk, the presence

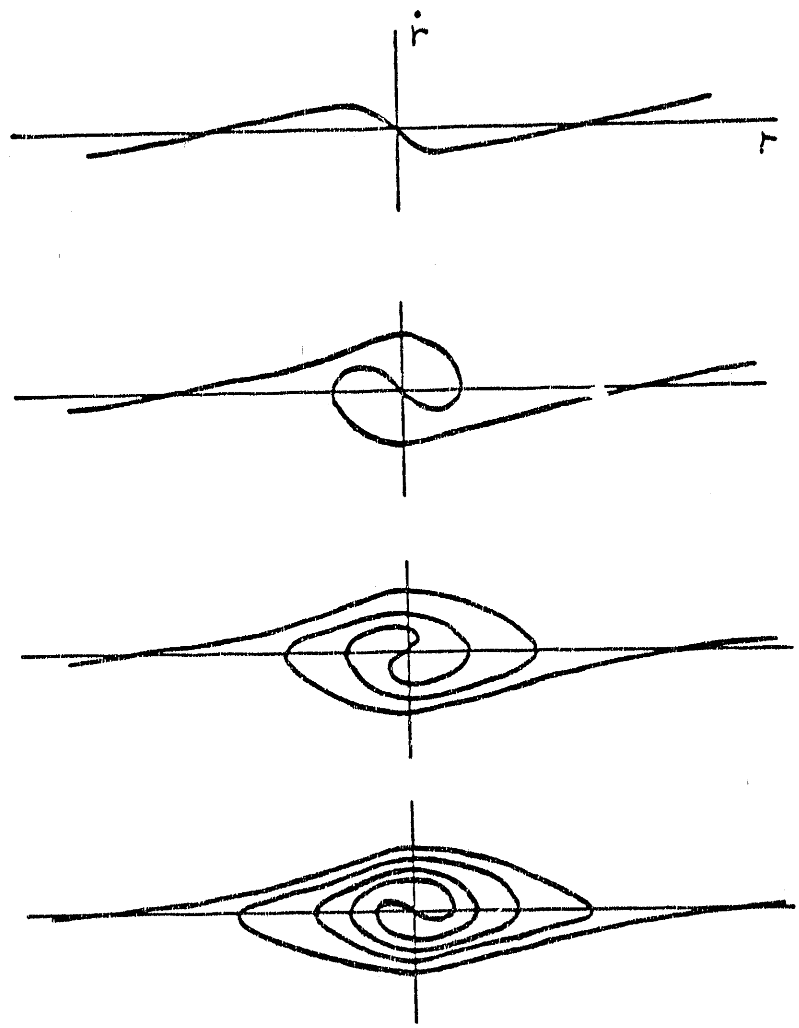

Figure IV. A qualitative picture of the time evolution of the phase-space distribution of cold dark matter particles in a galactic halo.

of stars, globular clusters and large molecular clouds and by primordial density perturbations on scales less than $10^{23}cm$. We found that although these effects change the number of peaks and broaden them, they are far from erasing the peak structure entirely. In particular, the peak due to dark matter particles which are falling onto the Galaxy for the first time as well as the peaks due to particles which have crossed the central parts of the Galaxy only a few times in the past are not erased. They form a record of the Galaxy's history which would become immediately available in a cavity detector of galactic halo axions if a signal is discovered.

Indeed, the energy resolution of a cavity galactic axion detector using the same data processing technique as the U. of Florida pilot detector[33] is set by the stability of the local oscillator with which the output of the cavity is compared.

The quartz oscillator presently in use has a stability of order $\frac{\Delta f}{f} \simeq 10^{-11}$ and thus allows one to reach $\frac{\Delta E}{E} \simeq 10^{-5}$ where $E \simeq 10^{-6} m_a$ is the kinetic energy of the dark matter axions. The resolution is certainly sufficient to measure the energies and the intensities of the peaks in the axion signal. The corresponding sheet velocity can also be readily obtained by modelling the diurnal frequency modulations ($\frac{\Delta f}{f} \simeq 10^{-9}$) of the peaks due to the earth's rotation.

ACKNOWLEDGEMENTS

This work was supported in part by the U.S. Department of Energy under contract No. DE-FG05-86ER40272, and by the National Science Foundation under Grant No. PHY89-04035.

REFERENCES

1. R. D. Peccei and H. Quinn, *Phys. Rev. Lett.* **38** (1977) 1440 and *Phys. Rev.* **D16** (1977) 1791.
2. S. Weinberg, *Phys. Rev. Lett.* 40 (1978) 223; F. Wilczek, *Phys. Rev. Lett.* 40 (1978) 279.
3. Recent 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.
4. G. 't Hooft, *Phys. Rev. Lett.* **37** (1976) 8 and *Phys. Rev.* **D14** (1976) 3432; R. Jackiw and C. Rebbi, *Phys. Rev. Lett.* **37** (1976) 172; C. G. Callan, R. F. Dashen and D. J. Gross, *Phys. Rev. Lett.* **63B** (1976) 334.
5. V. Baluni, *Phys. Rev.* **D19** (1979) 2227; R. J. Crewther, P. Di Vecchia, G. Veneziano and E. Witten, *Phys. Lett.* **88B** (1979) 123.
6. I. S. Altarev et al., *JETPL* **44** (1986) 460; J. M. Pendlebury et al., *Phys. Lett.* **B136** (1984) 327.
7. M. Kobayashi and K. Maskawa, *Progr. Theor. Phys.* **49** (1973) 652.
8. C. Vafa and E. Witten, *Phys. Rev. Lett.* **53** (1984) 535.
9. H. Georgi and L. Randall, *Nucl. Phys.* **B276** (1986) 241.
10. W. A. Bardeen and S.-H. H. Tye, *Phys. Lett.* **74B** (1978) 229; J. Ellis and M. K. Gaillard, *Phys. Lett.* **74B** (1978) 374; T. W. Donnelly et al., *Phys. Rev.* **D18** (1978) 1607. More recent treatments include: M. Srednicki, *Nucl. Phys.* **B260** (1985) 689; P. Sikivie in "Cosmology and Particle Physics", ed. by E. Alvarez et al., *World Scientific* (1987), p. 143-169.
11. H. Georgi, H. R. Quinn and S. Weinberg, *Phys. Rev. Lett.* **33** (1974) 451.
12. J. Kim, *Phys. Rev. Lett.* **43** (1979) 103; M. A. Shifman, A. I. Vainshtein and V. I. Zakharov, *Nucl. Phys.* **B166** (1980) 493; M. Dine, W. Fischler and M. Srednicki, *Phys. Lett.* **104B** (1981) 199; A. P. Zhitnitskii, *Sov. J. Nucl.* **31** (1980) 260.
13. D. Dicus, E. Kolb, V. Teplitz and R. Wagoner, *Phys. Rev.* **D18** (1978) 1829 and *Phys. Rev.* **D22** (1980) 839; M. Fukugita, S. Watamura and M. Yoshimura, *Phys. Rev. Lett.* **48** (1982) 1522 and *Phys. Rev.* **D26** (1982) 1840; N. Iwamoto, *Phys. Rev. Lett.* **53** (1984) 1198; L. M. Krauss, J. E. Moody, and F. Wilczek, *Phys. Lett.* **144B** (1984) 391; D. Dearborn, D. N. Schramm and G. Steigman, *Phys. Rev. Lett.* **56** (1986) 26; G. G. Raffelt, *Phys. Lett.* **166B** (1986) 402; G. G. Raffelt and D. Dearborn, *Phys. Rev.*

D36 (1987) 2201.

14. G. Raffelt and D. Seckel, *Phys. Rev. Lett.* **60** (1988) 1793; M. S. Turner, *Phys. Rev. Lett.* **60** (1988) 1797; R. Mayle et al., *Phys. Lett.* **203B** (1988) 188; T. Hatsuda and M. Yoshimura, *Phys. Lett.* **203B** (1988) 469; A. Burrows, M. S. Turner and R. P. Brinkmann, *Phys. Rev.* **D39** (1989) 1020; N. Iwamoto, Univ. of Toledo preprint (1988); A. Burrows, M. T. Ressell and M. S. Turner, Phys. Rev. **D42** (1990) 3297.
15. M. A. Bershady, M. T. Ressell and M. S. Turner, *Phys. Rev. Lett.* **66** (1991) 1398; M. T. Ressell, Phys. Rev. **D44** (1991) 3001.
16. P. Sikivie, *Phys. Rev. Lett.* **48** (1982) 1156.
17. A. Vilenkin and A. E. Everett, *Phys. Rev. Lett.* **48** (1982) 1867; T. W. Kibble, G. Lazarides and Q. Shafi, *Phys. Rev.* **D26** (1982) 435.
18. P. Sikivie, "Axions in Cosmology", in Where are the Elementary Particles, *Proc. 14th Summer School of Particle Physics* (Gif-sur-Yvette, 1982), eds. P. Fayet et al.
19. L. Abbott and P. Sikivie, *Phys. Lett.* **120B** (1983) 133; J. Preskill, M. Wise and F. Wilczek, *Phys. Lett.* **120B** (1983) 127; M. Dine and W. Fischler, *Phys. Lett.* **120B** (1983) 137.
20. D. J. Gross, R. D. Pisarski and L. G. Yaffe, *Rev. Mod. Phys.* **53** (1981) 43.
21. A. Vilenkin, *Phys. Rev. Lett.* **48** (1982) 59.
22. P. J. Steinhardt and M. S. Turner, *Phys. Lett.* **B129** (1983) 51.
23. For discussions of the uncertainties in the cosmological axion energy density due to the first order nature of the QCD phase transition see: W. G. Unruh and R. M. Wald, *Phys. Rev.* **D32** (1985) 831; M.S. Turner, *Phys. Rev.* **D32** (1985) 843; T. De Grand, T. W. Kephart and T. J. Weiler, *Phys. Rev.* **D33** (1986) 910; M. Hindmarsh, *Phys. Rev.* **D45** (1992) 1130.
24. M. Turner and F. Wilczek, *Phys. Rev. Lett.* **66** (1991) 5; A. Linde, *Phys. Lett.* **B259** (1991) 38 and references therein; D. Lyth, *Phys. Lett.* **B236** (1990) 408.
25. J. Ipser and P. Sikivie, *Phys. Rev. Lett.* **50** (1983) 925.
26. R. Davis, *Phys. Rev.* **D32** (1985) 3172, and *Phys. Lett.* **180B** (1986) 225; 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. M. Quashnock, *Nucl. Phys.* **B333** (1990) 815.
27. D. Harari and P. Sikivie, *Phys. Lett.* **B195** (1987) 361; T. Vachaspati, "Global String Dynamics", Cambridge preprint (Oct. 1988), unpublished.
28. C. Hagmann and P. Sikivie, *Nucl. Phys.* **B363** (1991) 247.
29. M. S. Turner, *Phys. Rev.* **D33** (1986) 889.
30. P. Sikivie, *Phys. Rev. Lett.* **51** (1983) 1415.
31. P. Sikivie, *Phys. Rev.* **D32** (1985) 2988.
32. L. Krauss, J. Moody, F. Wilczek and D. Morris, *Phys. Rev. Lett.* **55** (1985) 1797.
33. C. Hagmann et al., *Phys. Rev.* **D42** (1990) 1297.
34. C. Hagmann et al., *Rev. Sci. Inst.* **61** (1990) 1076.
35. S. DePanfilis et al., *Phys. Rev. Lett.* **59** (1987) 839; Phys. Rev. **D40** (1989) 3153.
36. S. Inagaki et al., in "Cosmic Axions", edited by C. Jones and A. Melissinos, World Scientific Publ., 1990, pp 58-78.
37. S. Matsuki and K. Yamamoto, Phys. Lett. **B263** (1991) 523.
38. K. Van Bibber et al., "A proposed search for dark matter axions in the 0.6

- 16 μeV range", LLNL preprint UCRL-JC-106876 (March 1991).
39. P. Sikivie, N. Sullivan and D. Tanner, unpublished.
40. M. Turner, *Phys. Rev.* **D42** (1990) 3572.
41. See, for example: R. A. Flores, Phys. Lett. **B215** (1988) 73; P. F. Smith and J. D. Lewin, Phys. Rep. **187** (1990) 203; M. Turner in Ref. (40).
42. D. Lynden-Bell, *Mon. Nat. Roy. Astr. Soc.* **136** (1967) 101.
43. P. Sikivie and J. Ipser, Univ. of Florida preprint UFIFT-HET-91-17 (Aug. 1991), and U.C. Santa Barbara preprint NSF-ITP-92-92 (June 1992).

END

DATE FILMED

11 / 16 / 92