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AXIONS IN ASTROPHYSICS AND COSMOLOGY*

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1. Introduction

The axion[1,2] was postulated approximately seven years ago to explain why the strong interactions conserve P and CP.[3] The parameter that sets the amount of P and CP violation in QCD is

$$\bar{\theta} = \theta - arg\ det\ m \quad , \tag{1}$$

where m is the quark mass matrix and θ is the coefficient of $(g^2/32\pi^2)\, G^a_{\mu\nu}\, \tilde{G}^{a\mu\nu}$ in the action density for QCD. Using the Adler–Bell–Jackiw anomaly,[4] one readily shows that QCD depends on θ and $arg\ det\ m$ only through the combination (1). Because $G\tilde{G}$ is a four-divergence, the $\bar{\theta}$ dependence of QCD is due purely to quantum effects. Quantum effects are most important when the coupling constant is large, i.e., in the case of QCD, at energies below a few GeV. They can be more or less reliably calculated using instanton and current algebra techniques.

The present upper limit on the neutron electric dipole moment requires[5] $\bar{\theta} \lesssim 10^{-8}$. If the CP violation necessary to explain $K_L \rightarrow 2\pi$ is introduced into the standard $SU_L(2) \times U_Y(1) \times SU^c(3)$ model of particle interactions in the manner of Kobayaski and Maskawa,[6] then $arg\ det\ m$ is an arbitrary (random) angle and there is absolutely no reason why $\bar{\theta} \lesssim 10^{-8}$. Other methods of introducing CP violation into the standard model also suffer from this difficulty[7] which is believed to be quite general and which has been given the name of "strong CP problem".

Peccei and Quinn[1] proposed the following simple and elegant solution to the problem. Let us postulate a $U_{PQ}(1)$ symmetry for the classical action density under which the quark fields and scalar fields transform generically as follows:

$$q \rightarrow e^{i\beta\gamma_5} q \ , \quad \varphi \rightarrow e^{-2i\beta} \varphi \quad . \tag{2}$$

The Yukawa interactions and scalar self-interactions have the following general form

$$-K\, q_L^\dagger q_R \phi + \text{h.c.} - V(\phi^\dagger \phi) \quad . \tag{3}$$

$V(\phi^\dagger\phi)$ has the shape of a "Mexican hat" and hence

$$\begin{aligned} \langle \phi \rangle &= v e^{i\alpha} \\ m &= K v e^{i\alpha} \\ \bar{\theta} &= \theta - N(\delta + \alpha) \end{aligned} \tag{4}$$

where δ is the overall phase of the matrix K of Yukawa couplings, and N is a model dependent integer.

Because of the $\text{U}_{PQ}(1)$ symmetry, the bottom of the Mexican hat potential is degenerate and the value of α is indifferent at the classical level. But the quantum effects (instantons ...) which make the physics of QCD $\bar{\theta}$–dependent will lift this degeneracy and align α in a particular direction. The most straightforward way to determine the direction of alignment is by minimizing the Yukawa interaction energy and using the fact that QCD produces quark-antiquark condensates $\langle q_L^\dagger q_R \rangle$ which are CP conserving.[8] One readily finds that α aligns in such a way that $\bar{\theta} = 0$. The strong CP problem is thus solved.

Weinberg and Wilczek[2] independently pointed out that the Peccei–Quinn solution to the strong CP problem implies the existence of a light pseudoscalar particle, which they called the axion. The axion is the pseudo–Nambu–Goldstone boson associated with the spontaneous breaking of the $\text{U}_{PQ}(1)$ quasi-symmetry; *i.e.,* it is the degree of freedom corresponding to rolling at the bottom of the

Mexican hat potential. The axion would be massless if $U_{PQ}(1)$ were not broken by QCD instanton effects. One can compute the axion mass using the same considerations as those which determine the alignment of α discussed above. The result is

$$m_a \simeq \frac{f_\pi m_\pi}{v} \left(\frac{N}{6}\right) \simeq 50 \text{ keV} \frac{250 \text{ GeV}}{v} \left(\frac{N}{6}\right) \quad (5)$$

where v and N are the quantities that appear in Eq. (4). In Eq. (5), v and N are normalized to the values they had in the earliest axion models. The coupling of the axion to quarks is

$$\sim i \frac{a}{v} m_q \bar{q} \gamma_5 q \quad . \quad (6)$$

The coupling of the axion to the electromagnetic field is

$$-\frac{4\alpha}{3\pi} N \frac{a}{V} \vec{E} \cdot \vec{B} \quad . \quad (7)$$

The value of the coupling strength[9] given in Eq. (7) holds for grand unified theories in which the unrenormalized value of the electroweak angle is $\sin^2 \theta_w^o = 3/8$. Note that both the axion mass, Eq. (5), and its coupling to the electromagnetic field , Eq. (7), are proportional to N/v. We will call the combination $f_a = v(6/N)$ the axion decay constant. The presence of the factor 6 is due to historical considerations.

It was first thought that the breaking of $U_{PQ}(1)$ occurred at the electroweak scale; *i.e.*, $v \simeq 250$ GeV. The corresponding axion was searched for in K, J/ψ and Υ decays and in reactor and beam dump experiments, but it was not found. Soon, however, it was discovered[10] 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. For a while, it was thought that the strong CP problem was solved without any presently observable consequences whatsoever.

Fortunately, astrophysics and cosmology came to the rescue. As we will see in Sec. III, they provide us with arguments that imply the axion decay constant should lie in the range 10^8 GeV $\lesssim f_a \lesssim 10^{12}$ GeV. A second cosmological constraint arises because axion models have, as a rule, multiple degenerate vacua and hence domain walls. In Sec. II we will describe the properties of these domain walls, the cosmological catastrophe they produce and the ways in which this catastrophe may be avoided. In Sec. IV we give the reasons why axions are an excellent candidate to constitute the dark matter of galactic halos. In Sec. V we describe detectors to look for axions floating about in the halo of our galaxy and for axions emitted by the sun.

2. Axionic Domain Walls

Axion models often have a spontaneously broken exact discrete symmetry.[8] In that case, they have discretely degenerate vacua and hence domain walls. The domain walls are the soliton-like boundaries between regions which happen to be in different vacua.

The exact discrete symmetry in question is the overlap of the group of anomaly free global flavor symmetries of the colored fermions (quarks, ...) with $U_{PQ}(1)$. For example, in the Dine–Fischler–Srednicki model[10] with n quarks

$$[\ \mathrm{SU}_L(n)\times \mathrm{SU}_R(n)\times \mathrm{U}_V(1)\ ]\ \cap\ \mathrm{U}_{PQ}(1)\ =\ Z(2n)\quad . \tag{8}$$

$Z(2n)$ is an exact discrete symmetry of that model. Indeed, as a subgroup of $\mathrm{U}_{PQ}(1)$ it is a symmetry of the classical action, and as a subgroup of the group of *anomaly-free* global symmetries of QCD, it is respected by the quantum effects as well. $Z(2n)$ is spontaneously broken down to $Z(2)$ by the vacuum expectation value $\langle\varphi\rangle = ve^{i\alpha}$ that breaks $\mathrm{U}_{PQ}(1)$ [see Eq. (4)]. Hence the Dine–Fischler–Srednicki model has n degenerate vacua, as many as there are quarks (*e.g.,* $n = 6$). In other axion models, however, the number N of degenerate vacua is different from the number of quarks. In general N is given by the formula[11,12]

$$2N = \frac{2\pi}{T_\theta}\sum_f Q_f\, t_f\quad . \tag{9}$$

Here the sum is over the colored left-handed fermions in the model, Q_f is their Peccei–Quin charge, t_f is their "color-anomaly" defined by $Tr\left(T_f^a T_f^b\right) = 1/2\ t_f\ \delta^{ab}$ where the T_f^a are the generators of $\mathrm{SU}^c(3)$ for the color representation to which the fermions f belong, and T_θ is the period of θ. $T_\theta = 2\pi$ for QCD, the standard $\mathrm{SU}^c(3)\times \mathrm{SU}_L(2)\times \mathrm{U}_Y(1)$ model and the SU(5) grand unified theory (GUT), but $T_\theta = 4\pi$ for the $O(10)$ GUT and $T_\theta = 6\pi$ for the E_6 GUT.[9] For example and for reasons that will soon become clear (see the cosmological domain wall problem below), Georgi and Wise[11] build a three generation SU(5) grand unified axion model with the fermion representation content $3(10)_1 + 3(\bar{5})_1 + 5(5)_{-1} + 5(\bar{5})_{-1}$ where the subscripts indicate the Peccei–Quinn charge of the corresponding multiplet. Using Eq. (9), one readily verifies that this model has a unique vacuum (N=1).

To derive the kinematic properties of the domain walls in axion models, one uses the effective action for the axion a

$$\begin{aligned} S_a &= \int d^4x \left[\frac{1}{2}\partial_\mu a\, \partial^\mu a - \frac{m_a^2 v^2}{N^2} f\left(N\, \frac{a}{v}\right) \right] \\ &= v^2 \int d^4x \left[\frac{1}{2}\partial_\mu \alpha\, \partial^\mu \alpha - \frac{m_a^2}{N^2} f(N\alpha) \right] \end{aligned} \tag{10}$$

where $\alpha = a/v$ denotes collectively all the phases that rotate under a $U_{PQ}(1)$ transformation and f is a periodic function of period 2π, whose Taylor expansion begins with $f(x) = \frac{1}{2}x^2 + \ldots$; for example, $f(x) = 1 - \cos x$. The axion self-interaction potential is then $Z(N)$ symmetric and m_a is the axion mass. A domain wall, in the $x - y$ plane for example, is the static classical solution $\alpha(z)$ obtained by minimizing the energy associated with Eq. (10) with the boundary conditions $\alpha(z) \to 0$ as $z \to -\infty$ and $\alpha(z) \to (2\pi)/N$ as $z \to +\infty$. One readily finds that the axionic domain walls have thickness of order m_a^{-1} and energy per unit surface $\sigma \simeq 8m_a v^2 \simeq 8f_\pi m_\pi v$. The tension in the domain wall equals its surface energy density σ. This follows from energy conservation and the fact that σ is a constant. The energy momentum tensor of a thin domain wall in the $x - y$ plane is thus

$$(T_{\mu\nu}) = \sigma\delta(z)\ \mathrm{diag}\ (1, -1, -1, 0) \quad . \tag{11}$$

Domain walls are a very unusual source of gravity. They are in fact gravitationally repulsive.[13–15] To clarify this statement, let us first remark that the Newtonian limit of Einstein gravity is valid only when T_{00} is much larger than the other components of $T_{\mu\nu}$. Hence, intuition derived from Newtonian gravity is inapplicable to the gravity of domain walls. Einstein's equations for planar domain walls have been solved exactly.[14–15] There is a unique reflection symmetric solution which is free of curvature singularities. It corresponds to a uniform gravitational field in which observers on either side are repelled by the domain wall with constant acceleration $2\pi G_N \sigma$, where G_N is Newton's gravitational constant. More generally it has been shown that,[14] for a wall of arbitrary shape and motion and with arbitrary tension τ and surface energy density σ, the sum of the accelerations towards the wall on both sides, as measured by observers hovering just off the wall, is $4\pi G_N(\sigma - 2\tau)$. For a dust wall ($\tau = 0$) one recovers the Newtonian result. For a domain wall ($\tau = \sigma$), the acceleration has equal magnitude as for a dust wall but opposite direction!

Axionic domain walls also have unusual electromagnetic properties.[9] To investigate these, one writes down the effective action density for photons and axions

$$\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \frac{1}{2}\partial_\mu a \, \partial^\mu a - \frac{m_a^2 v^2}{N^2} f\left(\frac{Na}{v}\right) + \frac{\alpha}{3\pi} \frac{Na}{v} F_{\mu\nu} \tilde{F}^{\mu\nu} \quad . \tag{12}$$

The strength of the $a\gamma\gamma$ coupling given in Eq. (12) is obtained by assuming that there is grand unification with the unrenormalized value of the electroweak angle $\sin^2\theta_W^\circ = 3/8$. The equations of motion derived from Eq. (12) are

$$\vec{\nabla} \cdot \left(\vec{E} - \frac{e^2}{3\pi^2} \frac{Na}{v} \vec{B}\right) = 0$$

$$\vec{\nabla} \times \left(\vec{B} + \frac{e^2}{3\pi^2} \frac{Na}{v} \vec{E}\right) - \frac{\partial}{\partial t}\left(\vec{E} - \frac{e^2}{3\pi^2} \frac{Na}{v} \vec{B}\right) = 0$$

$$\vec{\nabla} \cdot \vec{B} = 0$$

$$\vec{\nabla} \times \vec{E} + \frac{\partial \vec{B}}{\partial t} = 0$$

$$\Box\, a = -\frac{e^2 N}{3\pi^2 v}\, \vec{E} \cdot \vec{B} - \frac{m_a^2 v}{N} \sin \frac{Na}{v} \quad . \tag{13}$$

Consider an axionic domain wall of arbitrary shape and motion. Across the domain wall, $(Na)/v$ changes by 2π. The boundary conditions across the domain wall surface implied by Eqs. (13) in the thin wall approximation are

$$\Delta B_\perp = \Delta \vec{E}_\parallel = 0\ , \quad \Delta E_\perp = \frac{2e^2}{3\pi}\, B_\perp\ , \quad \Delta \vec{B}_\parallel = -\frac{2e^2}{3\pi}\, \vec{E}_\parallel\ . \tag{14}$$

We see that an axionic domain wall becomes electrically charged when traversed by magnetic flux. The electric surface charge density is

$$\sigma = \frac{2e^2}{3\pi}\, \vec{n} \cdot \vec{B} \tag{15}$$

where $\vec{n}$ is the unit normal in the direction of increasing $(Na)/v$. Similarly, an electric field parallel to an axionic domain wall induces a surface current density

$$\vec{K} = -\frac{2e^2}{3\pi}\, \vec{n} \times \vec{E} \quad . \tag{16}$$

These unusual effects are necessary to make sense of the Witten dyon charge[16] in the presence of axionic domain walls. Witten has shown that in a θ–vacuum magnetic monopoles acquire electric charge $q_\theta = C\theta g$ where g is the magnetic charge on the monopole and C is a model dependent constant. When a magnetic monopole traverses an axionic domain wall, the local value of θ changes by 2π and hence the electric charge on the monopole changes by one unit. One may well wonder whence that extra unit of electric charge came or what happens to electric charge conservation. The answer[9] is that the magnetic field of the monopole induces an electric charge density onto the domain wall, Eq. (15). When the monopole approaches the wall, the induced electric charge becomes concentrated near the impact point. It jumps onto the monopole when the monopole traverses the wall. The Witten dyon charge on the magnetic monopole plus the electric charge, Eq. (15), induced onto the axionic domain wall is conserved.

Domain walls exist in any theory in which a discrete symetry is spontaneously broken. In 1974, Zel'dovich, Kobzarev and Okun[17] pointed out that because of these domain walls the spontaneous breaking of an exact discrete symmetry is incompatible with standard cosmology. Their argument is very simple. The universe starts off at some very high temperature at which the discrete symmetry is unbroken. At some critical temperature, the spontaneous breakdown does occur and the order parameter chooses among several equally probable values (or directions), corresponding to the various vacua of the theory. Different regions of the universe will in general settle into different vacua and hence be separated by domain walls. In particular, regions which are outside each other's horizon are causally disconnected and thus totally uncorrelated. Hence, there will be at least on the order of one domain wall per horizon at any given time. The energy density in domain walls today would be

$$\rho_{\text{d.w.}}(t_0) \simeq \frac{\sigma}{t_0} = \rho_{\text{crit}} \left(\frac{\sigma}{10^{-5}\ \text{GeV}^3} \right) \tag{17}$$

where $t_0 \simeq 10^{10}$ years is the age of the universe today and $\rho_{\text{crit}} \simeq 10^{-29}$ gr/cm^3

is its present critical energy density for closure. Since $\sigma \simeq f_\pi m_\pi v >> 10^{-5}$ GeV^3, it is clear that if axions exist and $N > 1$, our present universe would be domain wall dominated many times over. But this can not be. A domain wall dominated universe would be expanding like $R \sim t^2$ (R is the cosmological scale parameter) and at a much higher rate than we observe today.

The cosmological domain wall problem just described can be avoided in a number of ways. Below are the three types of evasion which I am aware of:

1. The inflationary universe scenario[18] provides a solution if the inflationary epoch comes after the Peccei–Quinn phase transition at $T_{PQ} \simeq v$ where $U_{PQ}(1)$ is spontaneously broken. Indeed, inflation will align the phase α of $\langle\varphi\rangle = ve^{i\alpha}$ over enormous distances. Later, when the QCD instanton effects turn on at ~ 1 GeV temperatures, each enormous region will fall entirely into the same vacuum and hence be free of domain walls. For this to work, it is of course necessary that the post-inflation reheating temperature $T_{\rm reheat}$ be less than the temperature $T_{PQ} \simeq v$ at which the $U_{PQ}(1)$ symmetry is restored. We will see in the next section that v should be less than about 10^{12} GeV. On the other hand $T_{\rm reheat}$ must be sufficiently large for the baryon number asymmetry to be produced after inflation, since inflation wipes out any previous baryon number asymmetry. Hence, the set of constraints

$$T_{\substack{\rm baryo-\\ \rm genesis}} < T_{\rm reheat} < T_{PQ} \simeq v \lesssim 10^{12}\ {\rm GeV} \qquad (18)$$

which may be difficult to satisfy in practice.

2. It is possible to construct axion models which have a unique vacuum.[8,19,11,20,21] One way is to build the model in such a way that $N = 1$

where N is the integer given by Eq. (9) [NOTE: if $N = 0$, the Peccei–Quinn mechanism is inoperative; see Eq. (4)]. When $N = 1$, the model only has a discrete Z(2) symmetry which is not spontaneously broken. Hence the vacuum is unique. Many $N = 1$ models have been constructed, *e.g.,* Kim's original "invisible" axion model[10] and the grand unified axion model of Georgi and Wise[11] mentioned above. Another way to construct axion models with a unique vacuum is to embedd the discrete $Z(N)$ symmetry into a gauged[19,20] or an exact global continuous symmetry.[21] In that case, the N vacua are either gauge equivalent and hence not distinct or they are part of a larger continuous degeneracy and hence can be rotated into each other by adding coherent states of massless Nambu–Goldstone bosons.

The argument leading to the cosmological domain wall problem discussed above clearly does not apply to axion models with a unique vacuum. It is not immediately obvious, however, that such models are entirely free of cosmological difficulties because they, in fact, have domain walls, too.[23,12] When one traverses these domain walls one moves away from the unique vacuum and back to it along some topologically nontrivial path. This path is most readily visualized as one turn along the bottom of the Mexican hat potential $V(\varphi^\dagger\varphi)$ of Eq. (3) from the unique vacuum at $\alpha = 0$ through $\alpha = \pi$ and back to the vacuum at $\alpha = 2\pi$. These domain walls are quantum-mechanically unstable[24,12] because holes can be poked in them through some tunneling process. The rate for this process is very much smaller than the $(age)^{-1}$ of the universe, however, so that the domain walls are in fact stable for cosmological purposes. What saves axion models with a unique vacuum from the cosmological disaster of one domain wall per horizon at temperatures $\lesssim 1$ GeV is the earlier appearance of strings.[23,12] When, at temperature $T_{PQ} \simeq v$, the phase transition occurs where $U_{PQ}(1)$ becomes spontaneously broken by $\langle\varphi\rangle = ve^{i\alpha(x)}$, strings appear because π_1 [$U(1)$] $= Z$. When one moves around the string once, the local value of $\alpha(x)$ varies from 0 to 2π. From the usual causality arguments one expects at least on the order of one string per horizon from T_{PQ} onward untill QCD temperatures when the domain walls appear. Each string then becomes the edge of a domain wall. The typical size of a domain wall bounded by a string or of a closed domain wall is the horizon size ($\simeq 10^{-4}$ sec) at QCD temperatures. The probability of finding a domain wall much larger than that is exponentially small.[12] The finite size domain walls oscillate and dissipate away long before they dominate the cosmological energy density.

3. The last evasion of the domain wall problem is based on the observation[8,22] that a *tiny* explicit breaking of the $Z(N)$ symmetry is sufficient to make the domain walls disappear before they dominate the energy density. A soft explicit breaking of $U_{PQ}(1)$ will introduce shifts $\langle\Delta\mathcal{H}\rangle$ in energy density amongst the various vacua, and it will also introduce a finite value of $\bar{\theta}$. When the domain bubbles have average size $r_B = \sigma/\langle\Delta\mathcal{H}\rangle$, the differences in volume energy among bubbles is of order their surface energy, the $Z(N)$ breaking effects become important and the true vacuum takes over. One can show that the $\bar{\theta} \lesssim 10^{-8}$ constraint can be made compatible with the requirement that the domain walls disappear before they dominate the energy density provided $v \lesssim 10^{15}$ GeV. Finally, we note that an explicit breaking of the $Z(N)$ symmetry is of course very artificial if done by hand. On the other hand, this evasion of the domain wall problem is a natural

property of the ultimate theory of the world if the latter has in its low energy effective theory an automatic $U_{PQ}(1)$ which is then explicitly broken by higher order corrections.

3. Astrophysical and Cosmological Constraints on the Axion Decay Constant

The astrophysical constraint[25] arises because stars emit the weakly coupled axions from their whole volume whereas they emit photons only from their surface. Axions are produced in Compton, Primakoff and bremstrahlung type processes when photons collide with nuclei and electrons in stellar interiors. Because the axions are so weakly coupled, they can leave the star without further collisions. It has been shown[25] that if 250 GeV $\lesssim f_a \lesssim 10^8$ GeV, axion emission by stars is too copious to be consistent with our understanding of stellar evolution. If $f_a \lesssim 250$ GeV, the axion is too heavy to be produced in stars; if $f_a \gtrsim 10^8$ GeV, it is too weakly coupled to be produced overabundantly. Since $f_a \lesssim 250$ GeV appears to be ruled out by the unsuccessful laborabory searches, it follows that f_a should be larger than about 10^8 GeV.

The cosmological bound[26] ($fa \lesssim 10^{12}$ GeV) arises because axions are abundantly produced in the early universe when the temperature $T \simeq 1$ GeV. The argument is as follows. When T falls below $T_{PQ} \simeq v$, $U_{PQ}(1)$ becomes spontaneously broken by $\langle\varphi\rangle = ve^{i\alpha(x)}$. The values of $\alpha(\vec{x})$ are at that time randomly chosen since the QCD instanton effects which lift the degeneracy at the bottom of the Mexican hat potential are negligible when T is larger than a few GeV. $\alpha(\vec{x})$ is spatially inhomogeneous. However, all wiggles in $\alpha(\vec{x})$ which fall within the horizon at any given time will start to oscillate thenceforward and thus red-shift away.[27] The result of this is that, at any given time t, $\alpha(x)$ is approximately homogeneous over the horizon scale t. [Of course, if there is inflation with $T_{\rm reheat} < T_{PQ}$, $\alpha(\vec{x})$ is perfectly homogeneous over distances much larger than t. So much the better.] When QCD instanton effects become important at temperatures of order 1 GeV, the axion mass switches on and $\alpha(x)$ begins to oscillate with frequency m_a about the CP conserving minimum at $\alpha = 0$. Thus, at about 1 GeV temperature, a coherent state of nonrelativistic axions suddenly appears. The axions are nonrelativisitic because their momenta are of order the inverse of the horizon scale at QCD temperatures $t_{\rm QCD}^{-1} \simeq (10^{-4}\ \text{sec})^{-1} \simeq 10^{-11}$ eV, which is smaller than the axion mass (for $v \lesssim M_{\rm Planck}$). The axion energy density just after the axion mass has switched on is

$$\rho_a(t_{\rm QCD}) \simeq f_\pi^2\, m_\pi^2\, \alpha^2(t_{\rm QCD}) \simeq \rho_{\rm rad}(t_{\rm QCD})\, \alpha^2(t_{\rm QCD}) \tag{19}$$

where $\rho_{\rm rad}$ is the energy density in radiation. But the nonrelativistic energy density ρ_a decreases with time as R^{-3} (R is the cosmological scale factor), whereas $\rho_{\rm rad} \sim R^{-4}$. Moreover, for $v \gtrsim 10^8$ GeV, the axion fluid is effectively decoupled. It can be shown that the axions do not reheat nor convert into radiation.[26] Thus unless $\alpha^2(t_{\rm QCD})$ is very small, the universe will become axion matter dominated too soon. If we require the axion energy density today to be less than ten times $\rho_{\rm crit}$, we need $\alpha^2(t_{\rm QCD}) \lesssim 10^{-6}$.

How can $\alpha(t_{\rm QCD})$ be so small? If the switch-on of the axion mass were *sudden*, we would have $\alpha(t_{\rm QCD}) \sim O(1)$ and all axion models would be ruled out. "Sudden" means that the switch-on rate $(1/m_a)(dm_a/dt)$ is large compared to the frequency m_a at which $\alpha(t)$ oscillates. The opposite of "sudden" is "adiabatic":

$(1/m_a)(dm_a/dt)$ small compared to m_a. The latter regime is characterized by the adiabatic invariant

$$\oint p\, dq = \pi A^2(t)\; m_a(t) \simeq \text{time independent} \tag{20}$$

where $A(t)$ is the amplitude of the oscillation $\alpha(t) = A(t)\cos(m_a t + \delta)$. Equation (20) tells us that, in the adiabatic regime, the oscillation amplitude decreases while the axion mass is being switched on. Hence, provided the switch-on is sufficiently adiabatic, an excessive axion energy density may be avoided. Let us define a time t_1, such that $m_a^{-1}(t)\; (dm_a/dt) > m_a(t)$ for $t < t_1$, and $m_a^{-1}(t)\; (dm_a/dt) < m_a(t)$ for $t > t_1$. Before t_1, the switch-on is sudden. Hence $A(t_1) = O(1)$. After t_1, the switch-on is adiabatic. Hence

$$A^2(t_{\rm QCD}) = \frac{A^2(t_1)\; m_a(t_1)}{m_a(t_{\rm QCD})} = O(1)\; \frac{m_a(t_1)}{m_a} \quad . \tag{21}$$

The time dependence of the axion mass follows from its temperature dependence which has been calculated:[28] $m_a(t) = m_a\; [\; T(t)\;]$. Using this, the following result was obtained[26] for the axion energy density today

$$\rho_a(t_0) \simeq 5\rho_{\rm crit}\; \left(\frac{f_a}{10^{12}\ {\rm GeV}}\right)^{7/6} \quad . \tag{22}$$

Hence, the constraint $f_a \lesssim 10^{12}$ GeV which applies to all axion models independently of their vacuum structure and of the history of the universe before the temperature reached $O(10)$ GeV.

Steinhardt and Turner[29] have considered entropy production when the temperature of the universe is between 1 GeV and 1 MeV, by out-of-equilibrium decays of a relic particle species or by a first-order QCD phase transition, as a means to dilute the axion energy density and thus to weaken the $f_a \lesssim 10^{12}$ GeV bound.

4. Axions and Galaxy Formation

There is good evidence[30] that individual galaxies possess dark halos with masses exceeding that of the luminous galactic matter by a factor ~ 10. These galactic halos could be made of axions.[31] First, if $f_a \gtrsim 2 \times 10^{10}$ GeV, axionic matter is abundant enough to make up the halos [*cf.,* Eq. (22)]. Second, since axions are effectively decoupled for such large values of f_a, axionic halos are automatically dark. Neutrinos are similar to axions in these two respects and they have indeed been a very popular candidate for the halo matter. However, neutrino halo models have run into rather serious difficulties because the neutrino phase space density tends to be too small to allow them to cluster into galactic halos and because neutrino free streaming greatly inhibits the growth of all matter density perturbations on all mass scales less than about 10^{15} $M_\odot$. Axions, on the other hand, because they are nonrelativistic from the moment of their first appearance at ~ 1 GeV temperatures, have enormous phase space density and vanishingly small free streaming distance.[31] The large phase space density allows them to cluster easily into galactic halos, whereas the absence of free streaming allows the growth of primordial density perturbations to proceed on all scales. Recent computer simulations[32] of the growth of density perturbations in the early universe have shown that indeed cold dark matter (*e.g.,* axions, photinos, gravitinos ...) appears preferable to hot dark matter (*e.g.,* neutrinos).

5. "Invisible" Axion Detectors

One of the exciting aspects of the hypothesis that galactic halos are made of axions is the fact that it can be tested experimentally. How can this be done? One exploits the coupling, Eq. (7), of the axion to the electromagnetic field and the fact that we have available in the laboratory large oscillating electric and magnetic fields with frequencies or wave-vectors of order the axion mass ($\hbar = c = 1$)

$$\begin{aligned} m_a &\simeq 1.24\ 10^{-5}\ \mathrm{eV} \left(\frac{10^{12}\ \mathrm{GeV}}{f_a}\right) \\ &= \frac{2\pi}{10\ \mathrm{cm}} \left(\frac{10^{12}\ \mathrm{GeV}}{f_a}\right) \\ &= (2\pi)\ 3\ \mathrm{GHz} \left(\frac{10^{12}\ \mathrm{GeV}}{f_a}\right) \end{aligned} \quad . \tag{23}$$

The general idea[33] is that an externally applied magnetic or electric field will stimulate the conversion of an axion to a photon through the coupling, Eq. (7). The outgoing photon is relatively easy to detect. We will see below that this process can be used both for the detection of axions floating about in the halo of our galaxy and for the detection of axions emitted by our sun. In addition, one may attempt to observe the static forces with range of order m_a^{-1} due to virtual axion exchange. We refer the reader to the work of Moody and Wilczek[34] who have discussed these effects in detail.

5.1 AXION HALOSCOPE

If the Milky Way halo is composed of axions, their number density in our vicinity is approximately

$$\rho_a \simeq \frac{10^{-24}\ \text{gr}}{\text{cm}^3} \simeq \frac{(.5)\ 10^{14}\ \text{axions}}{\text{cm}^3} \left(\frac{f_a}{10^{12}\ \text{GeV}}\right) \quad . \tag{24}$$

These axions have energies

$$\epsilon_a = m_a \left(1 + \frac{1}{2}\,\beta^2\right) = m_a \left[1 + O(10^{-6})\right] \tag{25}$$

where $\beta \simeq 10^{-3}$ is the galactic virial velocity. Consider an electromagnetic cavity permeated by a strong static magnetic field $\vec{B}_0$. When the frequency ω of one of certain appropriate cavity modes equals the axion mass, there will be resonant conversion of Milky Way halo axions into quanta of excitation (photons) of that mode. For a rectangular cavity, the appropriate modes are $TM_{n\ell o}$ with n and ℓ odd (the longitudinal direction is that of $\vec{B}_0$). The power on resonance into such a mode is[33]

$$\begin{aligned} P_{n\ell} \simeq (.8)\ 10^{-19}\ \text{Watt} \left(\frac{V}{5 \times 10^4\ \text{cm}^3}\right) \left(\frac{B_0}{8\ \text{Tesla}}\right)^2 \\ \times \frac{1}{n^2\ell^2} \left(\frac{m_a}{1.24 \times 10^{-5}\ \text{eV}}\right) \text{Min}\left(1, 10^{-6}\, Q_{n\ell}\right) \end{aligned} \tag{26}$$

where V is the volume of the cavity and $Q_{n\ell}$ is the quality factor for that mode. Because the axion energies have spread of order $10^{-6}\, m_a$, the power [Eq. (26)] is not increased by having quality factors $Q > 10^6$. [However, if the axion were to be found and its mass were known, one could use superconducting cavities with $Q >> 10^6$ to *resolve* the spectrum of galactic halo axion energies.] To scan the allowed range of axion masses in a reasonable amount of time, a very sensitive detector of microwave radiation is required. A typical state of the art detector[35] today has a noise temperature T_N of order 10–20° K in the 1–40 GHz frequency range. The noise equivalent power of such detectors over a bandwidth set by the quality of the cavity [$B \simeq (2f/Q)$] is

$$\mathrm{NEP} = 1.2 \times 10^{-20} \, \frac{\mathrm{Watt}}{\sqrt{\mathrm{Hz}}} \left(\frac{f}{\mathrm{GHz}} \cdot \frac{10^6}{Q} \right)^{1/2} \left(\frac{T_N}{20^\circ \, \mathrm{K}} \right) \quad . \tag{27}$$

Comparison with Eq. (26) suggests that the experiment may be feasible. To keep the thermal noise below the signal, the cavity must be cooled to less than .1° K. Also, to distinguish the signal from fluctuations in thermal and detector noise, it will be necessary to modulate the resonant frequency of the cavity at some audio frequency and carry out phase-sensitive detection of the cavity output.

Rather detailed feasibility studies have been carried out for this experiment. It appears that a set of large cavities (say 150 cm long and 50 cm wide) placed in the bore of an 8 Tesla solenoidal superconducting magnet and equipped with state of the art microwave detectors should be able to cover, in one or two year continuous running time, the range of axion masses between 1 GHz and 30 or 40 GHz [10^{11} GeV $\lesssim f_a \lesssim 3 \times 10^{12}$ GeV] with a signal to noise ratio of three. It is very conceivable that design improvements will extend this range in both directions. Note that this experiment is capable of exploring that special range of values of the axion mass for which axions may provide the critical energy density for closing the universe.

5.2 AXION HELIOSCOPE

The solar axion flux on earth is approximately $.8\times10^{13}\,\mathrm{sec}^{-1}\,\mathrm{cm}^{-2}\,(10^8\,\mathrm{GeV}/f_a)^2$. The solar axions have a broad spectrum of energies centered about 1 keV, the temperature in the sun's interior. Using the coupling given in Eq. (7), one finds[33] the following general cross-section for axion $\rightarrow$ photon conversion in a volume V in which there is a static inhomogeneous magnetic field $\vec{B}_0(\vec{x})$

$$\sigma = \frac{1}{16\pi^2|\vec{\beta}_a|}\left(\frac{e^2N}{3\pi^2v}\right)^2 \sum_\lambda \int d^3k_\gamma\, \delta(k_\gamma - E_a) \times \left|\int_V d^3x\, e^{i\vec{q}\cdot\vec{x}}\, \vec{B}_0(\vec{x})\cdot\vec{\epsilon}(\vec{k}_\gamma,\lambda)\right|^2 \tag{28}$$

where $\vec{q} = \vec{k}_\gamma - \vec{k}_a$ and the sum is over photon polarizations. Consider then a detector of length L in the directions $\vec{n}$ of the sun, inside of which there is a transverse magnetic field $\vec{B}_0 = B_0\hat{t}\cos[(2\pi/d)\,\vec{n}\cdot\vec{x}]$. The cross-section for $a\rightarrow\gamma$ conversion for axions coming from the direction of the sun is

$$\sigma = \frac{1}{16|\vec{\beta}_a|}\left(\frac{e^2NB_0}{3\pi^2v}\right)^2 VL\left[\frac{\sin\left(\frac{2\pi}{d}-q_z\right)\frac{L}{2}}{\left(\frac{2\pi}{d}-q_z\right)\frac{L}{2}} + \frac{\sin\left(\frac{2\pi}{d}+q_z\right)\frac{L}{2}}{\left(\frac{2\pi}{d}+q_z\right)\frac{L}{2}}\right]^2 \tag{29}$$

where V is the detector volume and

$$\begin{aligned} q_z = (\vec{k}_\gamma - \vec{k}_a)\cdot\vec{n} &= E_a - \sqrt{E_a^2 - m_a^2} = \frac{1}{2}\frac{m_a^2}{E_a} \\ &\simeq \frac{2\pi}{16\ \mathrm{cm}}\left(\frac{10^8\ \mathrm{GeV}}{f_a}\right)^2\left(\frac{\mathrm{keV}}{E_a}\right)\ . \end{aligned} \tag{30}$$

On resonance ($q_z \simeq 2\pi/d$) the event rate in the detector is[33]

$$\frac{\#\mathrm{x-ray}}{\mathrm{time}} \simeq \frac{10^{-2}}{\mathrm{sec}}\,\frac{Vd}{(\mathrm{meter})^4}\left(\frac{B_0}{8\ \mathrm{Tesla}}\right)^2\left(\frac{10^8\ \mathrm{GeV}}{f_a}\right)^4\ . \tag{31}$$

Assuming that a signal of one x-ray/ten days can be distinguished from background, it appears that a cubic meter detector can detect solar axions if 10^8 GeV $\lesssim f_a \lesssim 10^9$ GeV. If f_a were to be in this range, axions would provide us with a powerful tool to study the solar interior.

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