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Cosmological Implications of Axinos

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ABSTRACT

Axinos are the supersymmetric partners of axions. They arise in models incorporating both low-energy supersymmetry and the axion solution of the strong CP problem. In the present state of knowledge several of the key properties of axinos, which control their cosmological consequences, are poorly determined. But generically there *are* very significant cosmological consequences, and we attempt to survey the possibilities here. In a wide variety of models the axino is the lightest R odd particle, and destabilizes the more conventional candidates for this title (photino, Higgsino ....) on cosmological time scales. While this consideration perhaps casts some shadow over an important class of dark-matter candidates, it turns out that in a large class of models the axino itself becomes a plausible dark-matter candidate. In other models the axino is heavy, and unstable. Even then axinos are of cosmological interest, because their decay can be the dominant mechanism for production of the lightest R odd particle.

1. Introduction

There are good reasons, both observational and especially theoretical, to suspect that most of the mass of the Universe is in some exotic and hitherto unknown "dark" form.

Quite independently of this, there are also good reasons to suspect the existence of specific, but exotic and as yet unobserved sorts of elementary particles. Among the best motivated and most widely discussed possibilities for widening the spectrum of elementary particles are to include axions on the one hand or sparticles (supersymmetric partner particles) on the other. Axions are desirable since they allow one to implement the Peccei-Quinn mechanism for resolving of the strong CP problem [1,2]; there seem to be no other comparably compelling ideas in this regard. Low energy supersymmetry may be desirable in helping one to understand the approximate masslessness (compared to the Planck mass or the mass where unification of gauge forces occurs) of the Higgs boson which violates $SU(2) \times U(1)$ electroweak symmetry [3,4]. Although existing concrete model implementations tend to be unwieldy and not entirely satisfactory, the basic thought that potentially large radiative corrections to this mass are naturally avoided by cancellations between boson and fermion loops is certainly promising, as are the constraints that supersymmetry imposes on dangerous scalar self-interaction terms even at the classical level.

These two attractive theoretical speculations are not at all mutually exclusive or contradictory. Indeed, axions and supersymmetry both find natural homes, and play important dynamical roles, within superstring theory.

It is quite remarkable that the particles whose existence is suggested by the internal logic of particle physics may very plausibly address the missing matter problem of cosmology. Upon estimating the cosmological production either of axions or of the lightest ordinary supersymmetric particle (abbreviated LOSP; see below) one finds that they are produced during the big bang and persist as a relic gas sufficiently exotic and dark, and with roughly the right average mass density, to provide the desired missing matter.

Two voluminous literatures have developed around these two dark matter candidates. However, strangely, there has been relatively little (though still not inconsiderable!) discussion of the consequences that follow when one simultaneously takes *both* attractive theoretical ideas — the axion solution of the strong

CP problem, and low-energy supersymmetry — seriously. That is what we shall do here.

We find that axinos — the supersymmetric partners of axions — have unique properties, and can have dramatic cosmological implications. The reason the axino is special is that it is very weakly interacting (for the same reasons as the axion), and in a wide class of models — but by no means generically — it is also very light (for similar reasons as the axion). The possible cosmological implications fall into three broad categories, that are realized within different models:

1. The axino may acquire a large mass, larger than that of some other more conventional R-odd particle. Then it is unstable, but generally extremely long-lived compared to familiar elementary particles of comparable mass. In this case the decay of axinos can easily become the dominant source of LOSPs. This effect at the very least modifies estimates of what mass and couplings the LOSP must have in order to provide the dark matter. In some parameter regimes one would simply produce too much mass by this mechanism. One derives constraints on particle physics models by requiring that they avoid this disaster.
2. The axino itself may well be the lightest supersymmetric particle. In this case the LOSP, which previously seemed an attractive dark matter candidate, would be lost. More precisely, the LOSP — generally some linear combination of photino, bino, and Higgsino ... — would be unstable on cosmological time scales against decay into an axino plus ordinary (*i.e.* R even) matter.

(It is noteworthy that although the LOSP is generically unstable on *cosmological* time scales, it is estimated to be quite stable on ordinary macroscopic or *laboratory* time scales. If and when the particles postulated by low energy supersymmetry are discovered, among them an apparently stable one, it will be important to check this apparent stability carefully.)

However light axinos themselves are stable, and *their* contribution to the mass density of the universe must be considered. If it is too large, we must discard the model; if it is too small, the axino has no cosmological significance (and neither does the LOSP). Most interesting of course is the golden mean —

3. The axinos themselves may provide the dark matter. It turns out, that in

a large class of models their contribution to the density is about right for this. For axinos to fulfill this role their mass should be a few keV. They would provide an example of *warm* dark matter.

Various bits and pieces of our discussion have appeared in previous analyses, which will be quoted as appropriate. However as far as we know a systematic discussion attempting to set out the most important alternatives has never appeared. Indeed most or all of the previous authors seem not to have been fully aware of the alternatives, and (because they made different tacit assumptions) reached quite different and often contradictory conclusions. In addition to the work of consolidation, we shall also have occasion to expand and revise some of the previous discussions, especially in regard to the treatment of thermal production of axinos during the big bang.

2. Axino types

Three different methods of implementing Peccei-Quinn symmetry are known. The different implementations lead to only slightly different quantitative predictions for axion phenomenology. As we shall soon see, however, their consequences for axino phenomenology are dramatically and qualitatively different.

The first implementation is essentially due to Kim [5]. In Kim's scheme, one postulates the existence of an $SU(2) \times U(1)$ singlet colored quark Q such that the right-handed component Q_R of Q carries Peccei-Quinn charge +1 while the left-handed component Q_L carries Peccei-Quinn charge -1. The complex scalar field ϕ whose phase will become the axion field is an $SU(3) \times SU(2) \times U(1)$ singlet with Peccei-Quinn charge +2. The important coupling among these fields is

$$L = f\phi\bar{Q}_R Q_L + \text{hermitean conjugate.} \tag{2.1}$$

When ϕ acquires a vacuum expectation value $\langle\phi\rangle = F$ this term generates a mass

$$m_Q = fF \tag{2.2}$$

for the Q quark. If we ignore fluctuations in the magnitude of ϕ (which correspond to very massive quanta) and write $\phi = Fe^{i\frac{a}{F}}$ then a will represent the axion field, with a properly normalized gradient energy term.

The interactions of the axion field with conventional matter degrees of freedom at energies and momenta well below m_Q may be obtained by integrating out the Q quark. The most significant interaction arises from a triangle diagram connecting the axion field to two gluons through an internal Q loop. This gives the effective coupling

$$L_{\text{eff.}} = \frac{a}{F}\frac{g^2}{32\pi^2}\text{tr } G^{\mu\nu}\tilde{G}_{\mu\nu}. \tag{2.3}$$

The second implementation is essentially due to Dine, Fischler and Srednicki (DFS) [6]. In their scheme one postulates the existence of two complex $SU(2)\times U(1)$ doublet scalar Higgs fields h_1 and h_2 carrying electroweak hypercharges $\mp\frac{1}{2}$ and Peccei-Quinn charges Q_d and Q_u respectively. No new quarks are required — instead, the quarks and leptons of the standard model have Peccei-Quinn charges. The right-handed charge $-\frac{1}{3}$ quarks and the right-handed leptons carry Peccei-Quinn charge $-Q_d$ and couple only to h_1 while the right-handed charge $+\frac{2}{3}$ quarks carry Peccei-Quinn charge $-Q_u$ and couple only to h_2. Thus far, we could have been describing the original model of Weinberg and Wilczek [2]. The crucial innovation of DFS is the introduction of an additional $SU(3)\times SU(2)\times U(1)$ singlet complex scalar field ϕ with Peccei-Quinn charge Q_ϕ where $2Q_\phi + Q_u + Q_d = 0$. This field can have an invariant coupling of the form

$$L_{\text{DFS}} = g\phi^2 h_1^\dagger h_2 + \text{hermitean conjugate}. \tag{2.4}$$

Because ϕ is an $SU(3)\times SU(2)\times U(1)$ singlet it is consistent to imagine that it acquires a large vacuum expectation value F. Writing $\phi = Fe^{i\frac{a}{F}}$ as before we again identify the properly normalized axion field a. Strictly speaking one must consider the full symmetry breaking (*i.e.* , the non-zero vacuum expectation values of h_1 and h_2 and chiral symmetry breaking), which introduces small corrections to this expression for a. These corrections can be very important for the phenomenology of physical axions, since they significantly change the couplings of axions to ordinary matter at low energies. This can (and does) happen because the coupling of the zero-order expression for a to matter is exceedingly small, and so even small fractional corrections to the form of a involving fields which couple more strongly can significantly affect the total coupling. However, for the purposes of this paper the zeroth order expression will be sufficient and appropriate.

A troublesome feature of the DFS scheme is that the coupling parameter g must be taken to be quite small, of order v/F, where v is the vacuum expectation value of the fields violating $SU(2)\times U(1)$ (*i.e.* $v \approx \frac{m_W}{e}$). This is because we want h_1 and h_2 to develop vacuum expectation values of order v. For a typical F of interest in axion physics, say $F \approx 10^{12}$GeV we must require $g \lesssim 10^{-9}$. For the moment we simply note this feature. It will play an important role later. An idea which may make this small parameter appear less arbitrary, due to Kim and Nilles [7], will be mentioned later.

Finally a third possibility should be mentioned. It is simply to postulate the existence of a scalar field a which has a coupling of the type (2.3) . This coupling is not renormalizable, and so without further knowledge of the true ultraviolet behavior of the ultimate theory in which it is embedded one cannot derive finite values for the radiative corrections it induces. This is an unfortunate situation, but one we already live with in the theory of gravity. (And one might take the attitude, as we do tacitly for gravity, that radiative corrections due to axion exchange though formally infinite are surely neglible.) Indeed it seems that superstring theory generically contains fields and couplings of precisely this type. Though it necessarily involves us in heavy guesswork, we shall try to analyze this case too.

For ease of reference we shall call the three implementations mentioned above I, II, and III respectively, and refer to the axinos they generate as being of type I, II, and III.

3. Axino mass for the several types

3.1 Tree level results for global supersymmetry [8]

To frame the context, it is appropriate briefly to recall some relevant general features of supersymmetric models.

A globally supersymmetric gauge theory contains gauge supermultiplets $(A^\alpha_\mu, \lambda^\alpha, D_\alpha)$ and chiral matter supermultiplets (ϕ^i, ψ^i_L, F_i). The Fs and Ds are auxiliary fields and generally do not represent propagating degrees of freedom. The scalar potential is constructed from a superpotential W which is a polynomial function in the chiral supermultiplets (complex conjugation is *not* allowed within this polynomial), and so-called D terms. The equation for the scalar potential is

$$V = (F_i)^* F_i + \frac{1}{4} D_\alpha D_\alpha \quad . \tag{3.1}$$

The equations of motion can be solved to eliminate the Fs and Ds, to give

$$F_i^* = \frac{\partial W}{\partial \phi^i} \,, \qquad D_\alpha = e_\alpha \sum_\phi \phi^* T_\alpha \phi \,, \tag{3.2}$$

where T_α is a representation of the α'th generator of the gauge group and e_α is the appropriate gauge coupling.

Now suppose that the Lagrangian is invariant under a global Peccei-Quinn U(1) transformation of the chiral supermultiplets:

$$\phi^i \to \exp(i\lambda Q_i)\phi^i \ . \tag{3.3}$$

Then the essence of Goldstone's theorem, which follows directly from the postulated invariance of the Lagrangian, is the statement that $G \cdot \mathcal{M}_B^2 = 0$, where

$$\mathcal{M}_{Bij}^2 = \frac{\partial^2 V}{\partial \phi_i \partial \phi_j} \tag{3.4}$$

is the boson mass matrix, and

$$G = \frac{1}{N}\left((Q\langle\phi\rangle)_i, -(Q\langle\phi^*\rangle)_i \right) \tag{3.5}$$

and where

$$N = \sqrt{\sum_i (Q\langle\phi\rangle)_i^2} \sim F_{PQ} \ . \tag{3.6}$$

The familar implication of all this is that there exists a massless axion field, which is expressed by G in the (ϕ^i, ϕ^{i*}) basis. (Note of course that $(Q\langle\phi\rangle)_i$ means $Q_i\langle\phi^i\rangle$, with no sum on i.) The normalization factor N insures that if the gradient energy of the ϕ fields was conventionally normalized, then so will be the gradient energy of the scalar field defined by G. Note that N is a more precise version, or generalization, of the parameter we called F before.

The axion field does not remain massless when non-perturbative effects are taken into account. We shall not review that famous story here.

In a supersymmetric theory the axion belongs to a supermultiplet, and it is important to keep tabs on the destiny of each member of this supermultiplet. The supermultiplet contains, in addition to the (pseudo)scalar axion field, a scalar field s — the saxion, and a Majorana fermion field $\tilde{a}$ — the axino. In the absence of supersymmetry breaking these fields would be degenerate with the axion, and the question arises how heavy they become when supersymmetry is broken.

The saxion field is given by

$$\eta = \frac{1}{N}\left((Q\langle\phi\rangle)_i, (Q\langle\phi^*\rangle)_i\right). \tag{3.7}$$

The axino field is given by

$$\Psi_L = \frac{1}{N}\left((Q\langle\phi\rangle)_i, 0\right) \tag{3.8}$$

in the $(\psi^i_L, \lambda^\alpha)$ basis.

For the axino the situation is quite simple, and the result quite interesting. From the general expression for the fermion couplings in the supersymmetric Lagrangian

$$\frac{1}{2}\frac{\partial^2 W}{\partial\phi_i\partial\phi_j}\psi^i_L\psi^j_L + \frac{\partial D_\alpha}{\partial\phi_i}\psi^i_L\lambda^\alpha \tag{3.9}$$

one readily finds (as pointed out by Tamvakis and Wyler [8]) that the fermion mass matrix acting on the axino is given by

$$(\mathcal{M}_F\Psi_L) = \left(-\frac{1}{N}(Q\langle F\rangle)_i, 0\right) \tag{3.10}$$

where the vector $\langle F\rangle_j$ is defined by

$$\langle F\rangle_j \equiv \langle F_j\rangle \equiv \langle \partial W/\partial\phi^j\rangle. \tag{3.11}$$

Now a non-zero vacuum expectation value for the F terms signals supersymmetry breaking. Thus in a theory with low-energy supersymmetry one finds the magnitudes of the components of $\langle F\rangle$ are at most of order $O(m^2_{SUSY})$, *i.e.* perhaps

$\sim$ (1 TeV)^2. The axino mass is therefore expected to receive a contribution of order

$$m_{\tilde{a}} \sim O\left(\frac{m_{SUSY}^2}{F_{PQ}}\right) \tag{3.12}$$

from this form of SUSY breaking. For $m_{SUSY} = 10^3$GeV and $F \sim 10^{12}$GeV, one finds $m_{\tilde{a}} \sim 1$keV.

A more elaborate algebraic analysis is necessary to discuss the saxion mass, and we shall merely report the main result and refer the interested reader to the literature [8,9]. The main result is that the saxion mass in global symmetry is of order m_{SUSY}^2/F if and only if the expectation values of the D terms vanish, $\langle D_\alpha \rangle = 0$ for all α. In other words, if supersymmetry is broken by D terms then Peccei-Quinn symmetry does not enforce extra cancellations and the saxion acquires a mass of order m_{SUSY}, as is typical for all conventional scalar particles in this framework.

It must be emphasized that these results have been derived within the assumption of global supersymmetry, and that they are purely classical (tree-graph) results. We must now discuss the changes wrought by supergravity, and by radiative corrections.

3.2 Supergravity and radiative corrections

Most recent attempts to implement low-energy supersymmetry have exploited the new possibilities offered for supersymmetry breaking terms within supergravity theories incorporating a "hidden" sector. The existence of a hidden sector, that is of matter which couples to the graviton supermultiplet but not to other known forms of matter with any sensible strength, was originally postulated *ad hoc* for model-building purposes, but has come to seem more natural within the context of heterotic string theory. In that theory the fundamental gauge group of the world is E(8)×E(8) and multiplets in the "other" E(8) effectively constitute a hidden sector.

We shall not need to enter into the murky depths of supergravity theory here. In fact so far the deep theory, such as it is, has not allowed any convincing improvement upon a simple pragmatic prescription, as follows.

Supersymmetry allows certain soft breaking terms, whose incorporation does not spoil the ultraviolet behavior of the theory. (That is, their incorporation does not require the inclusion of non-supersymmetric counterterms to cancel infinities in radiative corrections.) These soft terms are of very restricted forms [10]: $m^2\phi^\dagger\phi$, $m^2(\phi_i\phi_j + \text{h.c.})$, $m(\phi_i\phi_j\phi_k + \text{h.c.})$, and $m^2\bar{\lambda}\lambda$ where the ϕ's are the scalar components of the chiral superfields in the theory and λ can be any of the gauginos. These are, of course, superrenormalizable interactions, and occur in the Lagrangian with coefficients having units of mass to a positive power. The pragmatic prescription, which represents the current state of knowledge and conjecture is to represent the effects of supergravity with supersymmetry breaking in a hidden sector at low energy by allowing soft symmetry breaking terms to be added to a globally supersymmetric model. Generally, all the ϕ's and λ's get masses, and the superpotential determines which other soft terms appear. The mass scale for the coefficients of soft breaking terms is $m_{3/2} = \frac{M_S^2}{M_{\text{Pl.}}}$, where M_S is the primary scale of supersymmetry breaking and $M_{\text{Pl.}}$ is of course the Planck mass. The gravitino mass $m_{3/2}$ plays much the same role as our previous m_{SUSY}, and should be taken to be roughly the same order of magnitude. It is interesting that this determines the primary symmetry breaking scale to be $M_S \sim 10^{11}\text{GeV}$. This numerical value is quite close to the most desirable value for the symmetry breaking scale F in axion physics, and invites speculation that these two breakings are closely related. In addition, supergravity may generate effective *nonrenormalizable* corrections to the superpotential. These appear with coefficients proportional to an inverse power of mass. The scale of this mass is naturally taken to be $M_{\text{Pl.}}$. These terms, of course, do not violate supersymmetry. They are negligible in most contexts, but will play an important role at one point below.

Tree level mass terms for fermions in chiral multiplets are *not* soft. As is indicated in (3.8) , the only non-zero components of the axino field involve fermions in chiral (as opposed to gauge) multiplets. Therefore the previous result for the axino mass remains valid in supergravity theories at tree level.

On the other hands, soft terms certainly *can* contribute to the saxion mass at tree level — in fact a saxion mass term is a soft term all by itself directly. Thus in the spirit of supergravity theories the saxion is expected to be one of the vast multitude of particles with a mass of order $O(m_{3/2})$. Since it is R even, it decays predominantly to 2 gluons and is not cosmologically significant.

The most striking point remains that the axino mass is still tiny at the classical or tree-graph level, even taking into account the soft symmetry breaking terms. It remains to consider the effect of radiative corrections. These turn out to be very different for the various types of axion models described above.

Each of these models may be supersymmetrized in a straightforward fashion, although existing constructions of superpotentials which spontaneously implement the desired symmetry breaking pattern are rather awkward. (This seems to be true generally for low-energy supersymmetry models, whether or not they incorporate Peccei-Quinn symmetry.) In the supersymmetrized versions, ϕ is promoted to a chiral supermultiplet, and the axion is the phase of the scalar part of ϕ.

The type I or Kim axion model was also supersymmetrized by Kim [11]. Although for the sake of concreteness we shall discuss his specific model, the main conclusions seem to be independent of its details.

The superpotential is

$$W = f\hat{Q}\hat{\bar{Q}}\hat{\phi} + \text{hidden sector} \tag{3.13}$$

where ˆ means chiral superfield. Kim chooses his hidden sector so that when supersymmetry is broken ϕ gets a vev $\langle\Phi\rangle \sim F$ and hence the quarks Q acquire masses fF. F depends on the parameters of the hidden sector, but is of $O((M_{Pl}m_W)^{1/2})$. To consider the axino mass we need the entire superpotential complete with hidden sector:

$$W = f\hat{Q}\hat{\bar{Q}}\hat{\phi} + (\lambda_1\hat{\phi}\hat{\phi}' + m^2)\hat{Z} + (\lambda_2\hat{\phi}\hat{\phi}' + m'^2)\hat{Z}' + \mu^3 \tag{3.14}$$

The PQ transformation of the chiral superfields is:

$$\hat{\phi} \to \exp(2i\lambda)\hat{\phi} \qquad \hat{\phi}' \to \exp(-2i\lambda)\hat{\phi}' \tag{3.15}$$

$$\hat{Q} \to \exp(-i\lambda)\hat{Q} \qquad \hat{\bar{Q}} \to \exp(-i\lambda)\hat{\bar{Q}} \tag{3.16}$$

$$\hat{Z} \to \hat{Z} \qquad \hat{Z}' \to \hat{Z}'. \tag{3.17}$$

$\lambda_1, \lambda_2, m, m'$, and μ are the hidden sector parameters, chosen so that ϕ has a non-zero expectation value in the ground state.

Let us now apply the preceding analysis to this model. The only F terms which acquire vacuum expectation values are [5]

$$F_Z = \frac{\partial W}{\partial Z} \quad \text{and} \quad F_{Z'} = \frac{\partial W}{\partial Z'}. \tag{3.18}$$

Since Z and Z' have zero PQ charge, the axino has mass zero at tree level. This conclusion is valid at tree level in supergravity also.

However, the axino does acquire a significant mass at one loop. Supergravity breaking is expected to induce the soft term $m_{3/2}Af\phi\overline{Q}Q$ in the scalar potential, where $m_{3/2}$ is the gravitino mass and A is a number of order 1 (whose precise value is determined by the details of the supergravity theory.) $\phi, \overline{Q}$, and Q are the scalar parts of their respective superfields. The Feynman diagram exhibited in Figure 1 gives the axino a mass $m_{\tilde{a}} = (3/16\pi^2)f^2 A m_{3/2}$, as was shown by Moxhay and Yamamoto [12].

Moxhay and Yamamoto go on to suggest, on the basis of very specific dynamical assumptions, that $f^2 \sim O(\alpha_s)$. In general, though, (and certainly if the PQ symmetry is broken in a hidden sector, as suggested by Kim) f is not uniquely determined. So it seems appropriate to summarize the situation for type I axino models in the estimate

$$m_{\tilde{a}} \sim 10\text{GeV}\left(\frac{m_{3/2}f^2}{100\text{GeV}}\right) . \qquad \text{type I} \tag{3.19}$$

Now let us consider the type II or DFS axion scheme. In the DFS model, ϕ couples weakly to the Higgs scalars of electroweak symmetry breaking instead of to heavy quarks. This coupling is through the term (2.4) . The DFS model cannot be supersymmetrized exactly as it stands because a $\phi^2 h_1^i \epsilon_{ij} h_2^j$ term in the scalar potential cannot be obtained from a superpotential. A very similar model can be constructed, as follows. (The next few paragraphs are rather technical and can be skimmed without loss of continuity. In fact, we recommend this course to all but the most fanatic model-builders.)

Consider the N=1 supergravity extension of the standard model, as described for example in section 7 of Lahanas and Nanopoulos [4]. It includes the following chiral supermultiplets (family indices suppressed):

Chiral Supermultiplet	Transformation Under SU(3) × SU(2) × U(1)$_Y$
$\hat{Q} = \begin{pmatrix} \hat{U} \\ \hat{D} \end{pmatrix}$	$\left(3, 2, \frac{1}{6}\right)$
$\hat{U}^c$	$\left(\bar{3}, 1, -\frac{2}{3}\right)$
$\hat{D}^c$	$\left(\bar{3}, 1, \frac{1}{3}\right)$
$\hat{L} = \begin{pmatrix} \hat{N} \\ \hat{E} \end{pmatrix}$	$\left(1, 2, -\frac{1}{2}\right)$
$\hat{E}^c$	$(1, 1, 1)$
$\hat{H}_1 = \begin{pmatrix} \hat{H}_1^0 \\ \hat{H}_1^- \end{pmatrix}$	$\left(1, 2, -\frac{1}{2}\right)$
$\hat{H}_2 = \begin{pmatrix} \hat{H}_2^+ \\ \hat{H}_2^0 \end{pmatrix}$	$\left(1, 2, \frac{1}{2}\right)$

The low energy superpotential is

$$W = h_t \hat{H}_2^T \epsilon \hat{Q} \hat{U}^c + h_b \hat{H}_1^T \epsilon \hat{Q} \hat{D}^c + h_\tau \hat{H}_1^T \epsilon \hat{L} \hat{E}^c + m_4 \hat{H}_2^T \epsilon \hat{H}_1$$

where h_t, h_b, and h_τ are Yukawa couplings for the top and bottom quarks and the τ lepton. m_4 is a mass parameter which cannot be set to zero naturally and which affects the electroweak symmetry breaking in the model.

An axion of essentially the DFS type can be introduced by replacing the m_4 term in the superpotential with

$$k \hat{H} \hat{H}_2^T \epsilon \hat{H}_1$$

where $\hat{H}$ is a new superfield whose scalar part is to get a vev $\langle H \rangle \sim F$. The axion is the phase of H. k is a dimensionless parameter which must be small (see below) if the axion is to be invisible. The PQ transformation is given by:

$$\hat{H} \rightarrow exp(i\alpha Q_H)\hat{H}$$

$$\hat{H}_1 \rightarrow \exp(i\alpha Q_d)\hat{H}_1 \qquad \hat{H}_2 \rightarrow \exp(i\alpha Q_u)\hat{H}_2$$

$$\hat{D}^c \rightarrow \exp(-i\alpha Q_d)\hat{D}^c \qquad \hat{U}^c \rightarrow \exp(-i\alpha Q_u)\hat{U}^c$$

where $Q_H + Q_u + Q_d = 0$. This relationship among the PQ charges differs from that of the original DFS model only by a '2'.

The scalar potential is V = F terms + D terms + V_{soft} where V_{soft} are the soft terms induced by supergravity breaking. Everywhere that Lahanas and Nanopoulos have m_4, we must replace it by $k\langle H\rangle$. Since m_4 was $O(m_W)$ — actually O(10GeV) in the specific models of Lahanas and Nanopoulos — and since $\langle H\rangle \sim F \sim 10^{12}$ GeV, k must be of order 10^{-10}. There is no underlying motivation for k to be this small – it just has to be if the axion is to be invisible. This fine tuning is analogous to the fine tuning of g in (2.4) that was necessary in the original DFS model.

The scalar potential V in a direction where squark and slepton fields do not acquire vacuum expectation values is given by

$$V = V_F + V_D + V_{soft}$$

with

$$V_F = \left|\frac{\partial W}{\partial H_1}\right|^2 + \left|\frac{\partial W}{\partial H_2}\right|^2 + \left|\frac{\partial W}{\partial H}\right|^2 = k^2|H|^2|H_2|^2 + k^2|H|^2|H_1|^2 + k^2|H_2^T\epsilon H_1|^2$$

where the 3rd term can be neglected and

$$V_{soft} = m_1^2|H_1|^2 + m_2^2|H_2|^2 + m_H^2|H|^2 + kBm_{3/2}(HH_2^T\epsilon H_1 + h.c.).$$

V_D may be taken over unchanged from [4] since $\hat{H}$ is a gauge group singlet. H is the scalar part of $\hat{H}$, and similarly for the other superfields. The parameters m_1, m_2,and m_H are masses of $O(m_{3/2})$, and B is a constant of order 1. In principle, they are all determined by the breaking of the underlying supergravity model. Actually their values are directly determined at some high energy scale by supergravity breaking, and are renormalized as the energy scale is lowered. Electroweak symmetry breaking occurs at the energy scale where the renormalized parameters satisfy [4]

$$(m_1^2 + k^2F^2)(m_2^2 + k^2F^2) - (kBm_{3/2}F)^2 < 0.$$

When this condition is satisfied, the neutral components of H_1 and H_2 acquire vevs.

There seems to be no simple way of ensuring that the field H gets a large expectation value of order F. One possibility is to add a hidden sector, as Kim did in [11] and [7].

Let us now determine the mass of the axino in this model. The axino is given by

$$\tilde{a} = \frac{Q_H F\tilde{H} + Q_u\langle H_2\rangle \tilde{H}_2 + Q_d\langle H_1\rangle \tilde{H}_1}{\sqrt{Q_H^2F^2 + Q_u^2\langle H_2\rangle^2 + Q_d^2\langle H_1\rangle^2}}$$

After electroweak symmetry breaking, the F terms F_{H_1} and F_{H_2} get vacuum expectation values of order 10^3GeV (F_H is proportional to $1/F^2$ and hence is negligible.) Thus the argument leading to the estimate following (3.12) applies, *viz.*

$$m_{\tilde{a}} \sim 1\text{keV}\left(\frac{\text{F}}{10^{12}\text{GeV}}\right)^{-1}. \qquad \text{type II} \qquad (3.20)$$

In this type of model, as opposed to the type I or Kim model, there is no large one loop contribution to the axino mass. The reason for this difference is simply that k is so small. The type I axino may have a significantly strong coupling, namely that to the heavy quark, parametrized by f. Thus, it may be subject to significant radiative corrections. On the other hand the type II axino is truly weakly coupled at all scales, and is not significantly perturbed by radiative corrections.

It is appropriate to mention another possible approach to the construction of an axino extension to the DFS scheme [7], which seems to have attractive features (and in particular, may offer a possible rationale for the small coupling.) That is, we may stick to the original invariant form (2.4) but *interpret it directly as a term in the superpotential*, written in terms of superfields. This term would usually not be considered, because it is non-renormalizable. However, as mentioned before, the pragmatic approach to supergravity suggests that one ought to allow this term, with a coefficient inversely proportional to $M_{\mathrm{Pl.}}$. The term in the superpotential becomes

$$\frac{g}{M_{Pl.}}\hat{\phi}^2\hat{h}_1^T\epsilon\hat{h}_2 \tag{3.21}$$

In the scalar potential, this term will yield terms of the types

$$\left|\frac{\partial W}{\partial h_1}\right|^2 \to g^2\frac{F^4}{M_{\mathrm{Pl.}}^2}h_2^\dagger h_2 \tag{3.22}$$

and

$$\left|\frac{\partial W}{\partial \phi}\right|^2 \to g^2\frac{F^2}{M_{\mathrm{Pl.}}^2}\left|h_1^T\epsilon h_2\right|^2 \tag{3.23}$$

The second of these is completely harmless; the first is certainly acceptable for $g \leq 10^{-2}$. We will not investigate model building along these lines in detail (see [7] for an example) but even the simple considerations above indicate that there are possibilities for identifying the small parameter k with some independent parameters already required to be small on other grounds, *e.g.* $\frac{M_S}{M_{\mathrm{Pl.}}}$ or $\frac{F}{M_{\mathrm{Pl.}}}$.

Finally let us consider the type III axino scheme. The effective theory here looks much like the Kim scheme, but without the quark Q. One might suppose that in the absence of the heavy quark field those supergravity corrections to the squark mass which ultimately generated a sizable axino mass could not arise. One might therefore be tempted to expect a light axino in the type III scheme. However, the lesson we learn from the type I scenario is that it is not possible to draw firm conclusions without knowledge of the high energy theory which gives the effective coupling (2.3).

4. Elementary processes

The couplings of axinos are of course closely related to the familiar coupling of axions [13]. The most important of these may be derived by supersymmetry transformations of the axion couplings. Thus from the axino-gluon-gluon vertex

$$L_{\mathrm{aGG}} = \frac{\alpha_{\mathrm{strong}}}{8\pi} \frac{\mathrm{N}}{\mathrm{F}} a \frac{1}{2} \epsilon_{\mu\nu\rho\delta} \mathrm{G}^{c\mu\nu} \mathrm{G}^{c\rho\delta} \tag{4.1}$$

follows the axino-gluino-gluon vertex

$$L_{\tilde{\mathrm{a}}\tilde{\mathrm{G}}\mathrm{G}} = \frac{\alpha_{\mathrm{strong}}}{8\pi} \frac{\mathrm{N}}{\mathrm{F}} \bar{\tilde{a}} \gamma_5 \sigma^{\mu\nu} \tilde{G}^c G^c_{\mu\nu} . \tag{4.2}$$

N is the number of flavours of quarks with Peccei-Quinn charge. For type I models, N=1, while for type II models, N=6. These effective couplings are valid at energy scales between the Peccei-Quinn symmetry breaking scale F and the electroweak symmetry breaking scale. The constraints on the mass of the axion can be expressed as 10^{10}GeV $<$ F/N $<$ 10^{13}GeV [15]. Therefore, in the rest of this paper "F" means "F/N" and will usually appear as (F/10^{12}GeV). There are also effective couplings between axions and fermions and hence between axinos, fermions and sfermions, but we will not need these in what follows. The characteristic decay rates which are perhaps the most important qualitative feature of axino physics follow immediately from the form of the vertex (4.2) .

In principle there are two possibilities: either the axino is not the lightest supersymmetric particle, or it is. Given the estimates (3.19) and (3.20) , and existing experimental lower bounds on the mass of ordinary superparticles, the first case appears rather unlikely. Nevertheless, we consider it now. For the sake of concreteness, let us suppose that the photino is the lightest supersymmetric particle or LSP and that the only available decay channel is $\tilde{a} \to \tilde{\gamma}\gamma$. (The qualitative conclusion does not depend on this.) If the axino is much heavier than the photino, the axino will decay with a lifetime of order [16]

$$\tau_{\tilde{a}} \sim 10^{-1}\mathrm{sec} \left(\frac{F}{10^{12}\ \mathrm{GeV}}\right)^2 \left(\frac{m_{\tilde{a}}}{10^2\ \mathrm{GeV}}\right)^{-3} . \tag{4.3}$$

This is an unusual lifetime for such a heavy particle, and could have important cosmological consequences as we shall see. However it is difficult to imagine practical *laboratory* experiments which could sense the existence of an unstable axino, since its feeble couplings to ordinary matter at low energies preclude any significant production rate.

If the axino is the LSP, the situation is completely different. The lightest ordinary supersymmetric particle or LOSP (here taken to be the photino $\tilde{\gamma}$) will decay into it with a lifetime of order [16]

$$\tau_{\tilde{\gamma}} \sim 10^{-1}\text{sec.} \left(\frac{F}{10^{12}\ \text{GeV}}\right)^2 \left(\frac{m_{\tilde{\gamma}}}{10^2\ \text{GeV}}\right)^{-3} . \tag{4.4}$$

This estimate unfortunately depends cubically on the mass $m_{\tilde{\gamma}}$ and inverse quadratically on the coupling F, both of which are uncertain. However, over a wide range of the most interesting values of these parameters, the lifetime falls within a very interesting region. That is, it is short on cosmological but quite long on laboratory time scales. While the feebly coupled axino will not be produced in significant numbers in any accelerator, the LOSP will be relatively easy to produce at accelerators with sufficiently high reach in energy. Thus we might be presented with particles that would appear at first sight to be absolutely stable, but which could reveal their instability to searchers looking for their decay products downstream in a beam dump type experiment. A particularly interesting possibility opened up by this line of thought is that the LOSP could be *charged* or *colored*. For truly stable particles, this possibility is all but ruled out by the sensitive searches for trace contaminations of stable anomalous relics bound to ordinary matter, using mass spectroscopy. Given any reasonable rate of production in the big bang (certainly expected for the particles contemplated here), such charged or colored stable particles would almost surely not have escaped notice. However the existence of charged particles whose lifetime is short on cosmological time scales are not so constrained, and this speculative possibility perhaps gains a certain small degree of respectability from our considerations.

If the LOSP is some neutralino state other than the photino, several decay processes and mixing angles must be considered in determining its lifetime, but the result does not differ qualitatively from (4.4). If the LOSP carries an additive quantum number, somewhat different decay processes must be considered. Suppose, for example, that the LOSP is a sneutrino. Then the relevant decay processes are $\tilde{\nu} \to \tilde{a} Z \nu$ or $\tilde{\nu} \to \tilde{a} W e$. They proceed from elementary $\tilde{a}\tilde{Z}Z$ or $\tilde{a}\tilde{W}W$

vertices by attaching $\tilde{\nu}\nu$ or $\tilde{\nu}e$ lines respectively. (If the $\tilde{\nu}$ is lighter than the W, one must either allow the Z or W become virtual, and attach light fermion lines; or resort to $\tilde{\nu} \rightarrow \tilde{a}\nu\gamma$, which can arise from an underlying $\tilde{a}\tilde{Z}\gamma$ vertex.) These processes are all somewhat slower than the photino decay considered above, but still correspond to lifetimes much less than the age of the universe.

5. Axino cosmology

In forming expectations for axino cosmology, three factors are crucial. These are whether it is the axinos or the LOSPs that are stable, the lifetime of the quasi-stable species, and the potency of equilibration (*i.e.* production and destruction) processes. Together with the mass of the stable species, these factors determine the relic density, whose ratio to the critical density is the most important cosmological output. Since each of these factors is affected by various parameters and model choices which are poorly determined in the present state of knowledge, one must explore a variety of possibilities that may at first appear bewildering. However, we believe that by keeping in mind that it is these three factors that control the cosmology you will be able to follow the thread.

Let us begin with some general observations. Since the axinos are coupled only weakly to standard model particles, it seems likely that they will decouple early, when still relativistic. If they are stable (*i.e.* if they are the LSP), they will contribute to the present mass density of the universe, as dark matter. Because they decouple while relativistic, their abundance at decoupling is roughly equal to the photon density at that time. During adiabatic expansion these two numbers remain comparable. However their ratio does change in a calculable way because the entropy of annihilating species increases the photon number density but does not affect the axinos. This is the same effect that lowers the temperature of the relic neutrino gas relative to the microwave background. Taking it into account, we can use the observed photon density to estimate the mass density in axinos. We obtain a bound on the axino mass by requiring that the ratio $\Omega_{\tilde{a}}$ of the axino density to the critical density, and the Hubble constant h measured in units of (100 km s^{-1} Mpc^{-1}) satisfy $\Omega_{\tilde{a}}h^2 < 1$. Using standard estimates [14], this corresponds to

$$m_{\tilde{a}} < 12.8eV\left(\frac{g_*(T_D)}{g_{eff}}\right) \tag{5.1}$$

where $g_{eff} = 0.75g = 1.5$ for a two component axino, and $g_*(T_D)$ is the effective number of degrees of freedom (1 for bosonic degrees of freedom, 7/8 for fermionic) at the decoupling temperature T_D. In the minimal standard model if a species decouples before electroweak symmetry breaking, $g_*(T_D) = 106.75$. In a supersymmetric model with 2 complex Higgs doublets, if all the supersymmetric particles are in thermal equilibrium when the axino decouples then $g_*(T_D) = 228.75$. We conclude that if the axino is stable and is present in thermal equilibrium at high temperatures and decouples when relativistic, then

$$m_{\tilde{a}} < 2\text{keV}. \tag{5.2}$$

This 2 keV bound assumes that all the supersymmetric partners of the standard model particles are in equilibrium when the axino decouples. If only some of them are in equilibrium, the bound is between 2 keV and 1 keV.

Now let us consider the situation in the models we described above. In all three models, the axinos are handily maintained in thermal equilibrium at high temperatures through the reaction depicted in Figure 2. For $T < F$, the rate is given by

$$\Gamma \sim \frac{\alpha_s^3}{16\pi F^2} T^3. \tag{5.3}$$

The condition for thermal equilibrium is $\Gamma > H$, where the expansion rate or Hubble parameter is given according to standard assumptions [14] by $H = 1.66(g_*)^{1/2}T^2/M_{PL} \simeq 25T^2/M_{PL}$ for $g_* = 228.75$ at high temperatures. Putting these together, we find that the axinos decouple at

$$T_D \sim 10^{11} GeV \left(\frac{F}{10^{12} GeV}\right)^2 \left(\frac{\alpha_s}{0.1}\right)^{-3} \tag{5.4}$$

and are in equilibrium at higher temperatures. Thus the axino decouples at a high temperature, while still relativistic.

(A subtle point arises here, which merits explicit notice. The crucial anomalous couplings of axions and of axinos to gauge particles persist for virtual momenta substantially larger than the masses of the fermions in the triangle graphs which generate them. Therefore, the reaction in Figure 2 proceeds even when the temperature is much higher than the masses of the fermions in the triangle graphs but still lower than F. This reflects the fact that the anomalous Ward identities are *local* equations, that is that they hold for large momentum transfer and not only for the integrated charge. Perhaps the simplest way to understand this is to recognize that the anomalous contributions to the triangle graph arise from the necessity to include regulator terms. The effect of these regulator terms is most transparent if Pauli-Villars regulators are used: the anomaly arises because the violation of chiral symmetry associated with the large mass of the regulator particles persists even as the mass is taken to infinity. Because the essential contribution arises from these infinitely massive particles, it is not affected by finite momenta inserted at the external legs.)

In type I axino models $m_{\tilde{a}}$ from (3.19) is generically much greater than 2 keV. The type I models seem to be viable only if one or another of the following *a priori* unlikely conditions is true. First, f may be very small, so that $m_{\tilde{a}} \leq$ 2keV. Or second, the axino may be so heavy that it is not the LSP (as we mentioned above, this is not easy to arrange.) Or third, some of the cosmological hypotheses which went into the calculation are not true for the actual evolution of our universe. To generate a sufficient suppression of axino number density, requires a *big* departure from the standard radiation-dominated big bang to occur at temperatures $T \leq T_D$. Perhaps the most attractive idea along these lines is to suppose that a period of inflation intervened between T_D and the present, and the universe never subsequently reheated above this temperature.

Sophisticated readers will notice a close resemblance between the problems encountered here and the usual gravitino problem of supergravity theories. In that context, the most popular hypothesis seems to be the third one mentioned above. A significant difference between the two cases is that the decoupling temperature for axinos is smaller than what one finds for gravitinos, at least for $F \ll M_{\rm Pl}$. This exacerbates the problem, already serious for gravitinos, of reconciling a low reheating temperature with a plausible scenario for baryogenesis.

Let us further pursue the cosmological implications of heavy, unstable axinos in type I models, *i.e.* the second alternative above.

As above, let us suppose for the sake of concreteness that the photino is the LSP, and the only available decay channel is $\tilde{a} \rightarrow \gamma\tilde{\gamma}$. (Again the qualitative conclusion does not depend on this.) Unless the axinos decay fairly rapidly, the universe will become matter dominated by axinos for a brief period. The temperature at which matter and radiation contribute equally to the density is $T_{EQ} = 5.5\Omega_o h^2$eV [14] and the heavy axinos contribute $\Omega_{\tilde{a}} \sim m_{\tilde{a}}/2$keV. Therefore, the universe becomes matter dominated at

$$T_{EQ} \sim 300\mathrm{MeV}\left(\frac{m_{\tilde{a}}}{100\mathrm{GeV}}\right) \quad . \tag{5.5}$$

The lifetime (4.3) for the axino decay $\tilde{a} \to \tilde{\gamma}\gamma$ corresponds to a temperature [14]

$$T_{\text{axino decay}} \sim 3MeV\left(\frac{F}{10^{12}GeV}\right)^{-1}\left(\frac{m_{\tilde{a}}}{100GeV}\right)^{3/2} \quad . \tag{5.6}$$

Hence, the condition for an epoch of axino matter domination to occur is

$$m_{\tilde{a}} < 10^6\mathrm{GeV}\left(\frac{F}{10^{12}\mathrm{GeV}}\right) \quad . \tag{5.7}$$

When the axinos decay, they will increase the entropy per comoving volume by a factor of [14]

$$100\left(\frac{m_{\tilde{a}}}{100\mathrm{GeV}}\right)^{-1/2}\left(\frac{\mathrm{F}}{10^{12}\mathrm{GeV}}\right) \quad . \tag{5.8}$$

The temperature after the axinos decay and most of the photinos thermalize is

$$T_{RH} \sim 10\mathrm{MeV}\left(\frac{F}{10^{12}GeV}\right)^{-1}\left(\frac{m_{\tilde{a}}}{100GeV}\right)^{3/2} \quad . \tag{5.9}$$

If the axinos decay during or after the time of nucleosynthesis, the entropy production and the plethora of photons energetic enough to dissociate light nuclei will surely destroy the brilliant agreement between the standard calculations of element abundances and observations. Since this is unacceptable, we must require that the reheating temperature (5.9) is higher than 1 MeV. This gives a lower bound for the mass of the axino of

$$m_{\tilde{a}} > 20\mathrm{GeV}\left(\frac{F}{10^{12}\mathrm{GeV}}\right)^{2/3} \tag{5.10}$$

This is not a significant additional constraint.

More significant is the following consideration. As we have discussed, the axinos are relativistic when they decouple. Their density is therefore comparable to the photon density. If each decay of an axino ultimately produces at least one LSP with a mass of at least several GeV, we are evidently in grave danger of producing far more than critical density in LSPs. This can only be avoided if the LSPs produced in this way annihilate very efficiently.

As usual, for concreteness let us suppose that the axinos decay into photinos. If the axinos decay before photino annihilation reactions freeze out, then the relic photino abundance will be unaffected by axino decays. The freeze out of photino annihilations occurs generically at $\mathrm{T} \sim m_{\tilde{\gamma}}/20$. (This is generally true for species which decouple when nonrelativistic and only depends logarithmically on the annihilation cross section.) If the axino decays occur after the photinos freeze out, then the relic photino abundance is increased relative to its value in the absence of axino decays by a factor of $(m_{\tilde{\gamma}}/20\mathrm{T}_{\mathbf{axino\ decay}})^2$. Since the relic photino abundance is usually not too different from the closure density, a significant increase in the relic abundance can not be tolerated. The condition that axino decays do not increase the relic photino abundance is

$$\mathrm{T}_{\mathbf{axino\ decay}} > \left(\frac{m_{\tilde{\gamma}}}{20}\right) . \tag{5.11}$$

From (5.11) and (5.6) we conclude that in order to avoid producing more than the critical density in photinos, the axino must be heavier than

$$m_{\tilde{a}} > 3\mathrm{TeV}\left(\frac{F}{10^{12}\mathrm{GeV}}\right)^{2/3}\left(\frac{m_{\tilde{\gamma}}}{10\mathrm{GeV}}\right)^{2/3} \quad . \tag{5.12}$$

Given the estimate (3.19), this constraint seems very difficult to satisfy. The heavy unstable type I axino is a marginal possibility at best.

This concludes our discussion of type I axino cosmology. Clearly the overall conclusion must be that it is difficult to incorporate these models into a reasonable cosmological scenario.

Now let us consider type II axino models.

In this scenario, the axino is the LSP, and we must ensure that the decay of the LOSP does not upset nucleosynthesis. As usual, we consider the case where the

LOSP is a photino and decays via $\tilde{\gamma} \to \tilde{a}\gamma$. The entropy produced in this decay will not be significant since the universe will never have been LOSP dominated. Therefore the only constraint on the LOSP decay comes from the requirement that the photons from the decay must not be energetic enough to dissociate light nuclei. The single high energy decay photon will rapidly be degraded into many photons with energies $E \sim (1\mathrm{MeV})^2/(50T)$ [17] by pair production off photons in the high energy tail of the thermal photon distribution. Thus, the decay will yield photons with energies high enough to dissociate deuterium only if the temperature is below about 9 keV. Hence in the type II axino scenario, photinos must decay before a time of about 10^4 seconds so that the deuterium produced during nucleosynthesis is not dissociated. From the photino lifetime (4.4), we derive a lower bound on the photino mass of

$$m_{\tilde{\gamma}} > 2\mathrm{GeV}\left(\frac{\mathrm{F}}{10^{12}\mathrm{GeV}}\right)^{2/3}. \tag{5.13}$$

If the LOSP is not a photino, but rather a different mixture of neutralinos, the lower bound (5.13) on its mass will not be affected qualitatively. If the LOSP is a sneutrino the lifetime is longer, and the bound somewhat more restrictive. (If its decay products included only an axino and neutrinos, then the argument leading to the lower bound (5.13) on its mass does not apply. However this seems very unlikely, given the decay schemes we mentioned before.)

Now let us consider the cosmological implications of the type II axinos themselves. The reaction in Figure 2 keeps axinos in thermal equilibrium at high temperatures. The axinos will decouple at the temperature T_D given by (3.4) just as in the type I scenario. At the (supersymmetry violating) weak phase transition, the type II axinos acquire their small mass. Since they are so light, they are almost certainly the stable LSPs, and hence they contribute to the overall density of the universe as dark matter. From (5.2) and the discussion surrounding it, we must have $m_{\tilde{a}} < 2keV$, and of course if $m_{\tilde{a}} \sim 2keV$ then $\Omega_{\tilde{a}}h^2 \sim 1$. But this is exactly what is suggested by the estimate (3.20) !

Thermally produced light type II axinos are therefore reasonable candidates to form the wanted dark matter. They can contribute $\Omega_{\tilde{a}}h^2 \sim 1$ for reasonable values of $m_{\tilde{a}}$. It is very interesting to note that these light axinos are 'warm' dark matter. They are colder than hot dark matter like neutrinos because they decouple so early. They are hotter than cold dark matter, because they do decouple when relativistic. To be specific, let us consider axinos with a mass of 2 keV so that they contribute significantly to Ω. They will therefore become nonrelativistic when their temperature falls below 2 keV. At this point, the rest of the universe will be at a temperature of

$$\mathrm{T_{rad}} \sim 2\mathrm{keV} \left(\frac{g_*(\mathrm{T_D})}{g_*(1\mathrm{keV})} \right)^{1/3} \sim 10\mathrm{keV} \ . \tag{5.14}$$

From this point until matter radiation equality $\mathrm{T}_{\tilde{a}} \sim R^{-2}$ while $\mathrm{T_{rad}} \sim R^{-1}$ where R is the scale factor. Hence, when the universe becomes matter dominated at $T_{rad} \sim 5.5\Omega_o h^2 eV$ [14], the axinos are moving at a speed of about $v/c \sim 10^{-3}$.

In the type II axino scenario, there are two dark matter candidates — axinos which are warm dark matter and axions which are cold dark matter. For F between about 10^{11}GeV and 10^{12}GeV, both give significant contributions to the cosmological density. For larger values of F, the axions are dominant. For smaller values of F, the axinos are dominant.

Warm dark matter may offer significant advantages in the context of structure formation in the universe [18]. Because keV mass axinos are moving more slowly than neutrinos when the universe becomes matter dominated, the damping scale in an axino dominated universe is shorter than in a neutrino dominated universe. In fact, the damping scale is very close to 1 Mpc which is the scale of galactic perturbations. By way of comparison, the damping scale in a neutrino dominated universe is 40 Mpc/(m_ν/30eV). In an axino dominated universe, therefore, the first structures to form are galaxy-sized, rather than the supercluster-sized structures that form first in a hot dark matter model. While hot dark matter has some attractive features including the abundance of large scale structure it produces, its downfall is the failure to form galaxies sufficiently early, which can be traced directly to the large damping scale. Warm dark matter certainly corrects this difficulty.

Warm dark matter may also have advantages over cold dark matter. In the cold dark matter scenario the damping scale is very small ($\ll$ Mpc), and subgalactic-sized objects form first. While cold dark matter does an excellent job of reproducing the observed features of the universe on small scales, say less than about $20h^{-1}$ Mpc, there is growing concern that it does not reproduce the wealth of structure that observers are discovering on large scales. Warm dark matter may offer some help here. When the spectrum of density perturbations for warm dark matter is normalized in the same way as that for cold dark matter, *e.g.* $\frac{\delta\rho}{\rho} = 1$ on the scale $8h^{-1}$ Mpc, warm dark matter has more power on large scales than does cold dark matter. This additional power on large scales may be what is needed to to account for the structure on large scales.

In sum, a universe dominated by type II axinos with mass of order 1 keV may be very attractive. It seems to incorporate the best features of both hot and cold dark matter. Galaxies form first as in cold dark matter, but there is more power on large scales than in cold dark matter. These advantages, together with the observation that type II axinos provide a theoretically compelling warm dark matter candidate, suggest that warm dark matter may be worthy of further detailed numerical study.

6. Additional remarks, and conclusion

The type III axino is plausibly even much lighter than 1keV. If that is true, the axino will destabilize the LOSP without itself becoming a dark matter candidate.

Other interesting, theoretically motivated particles with possible cosmological implications that can be analyzed in the same spirit as above are the dilaton and its supersymmetric partner, the dilatino. They arise in superstring theory, but there is no consensus (even in the logarithm of the order of magnitude) as to the value of their mass. Their coupling is very weak; roughly speaking they behave like axions and axinos but with $\frac{1}{F} \to \frac{1}{M_{\rm Pl}}$. It would be worthwhile to analyze the bewildering variety of alternative possibilities the existence of these particles opens up, but we shall not attempt that here.

Let us summarize. We have found that the existence of axinos, which inevitably accompanies any attempt to implement the two appealing ideas of low-energy supersymmetry and Peccei-Quinn symmetry simultaneously, can drastically affect the status of the dark matter problem. Axinos can and generally do destabilize the lightest ordinary superparticles, which are popular dark matter candidates. While removing these candidates, they seem themselves to provide a *warm* dark matter candidate in one class of models.

Finally, a methodological remark. The models analyzed here seem (certainly when viewed in isolation, without usable guidance from a larger framework) contrived. This actually seems to be true of models incorporating low energy supersymmetry generally, with or without Peccei-Quinn symmetry. It may well be that the difficulties in model-building are a sign that supersymmetry is broken by some different mechanism, which is poorly represented by the Higgs field paradigm. If so, it will become important to analyze the consequences of this hypothetical new mechanism for axino cosmology. It is also possible that the complexity of the models is merely a sign that our present universe has undergone many layers of symmetry breaking, and that there really are lots and lots of massive and supermassive fields and condensates, so that to ask both for economy and for realism in the effective theory is to ask for the impossible.

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16. The lifetime for the decay of a heavy axino into a light photino is the same as that of a heavy photino decaying into a light axino. This calculation was done by Kim *et al.* in the limit of equal top and stop masses. (J. E. Kim, A.

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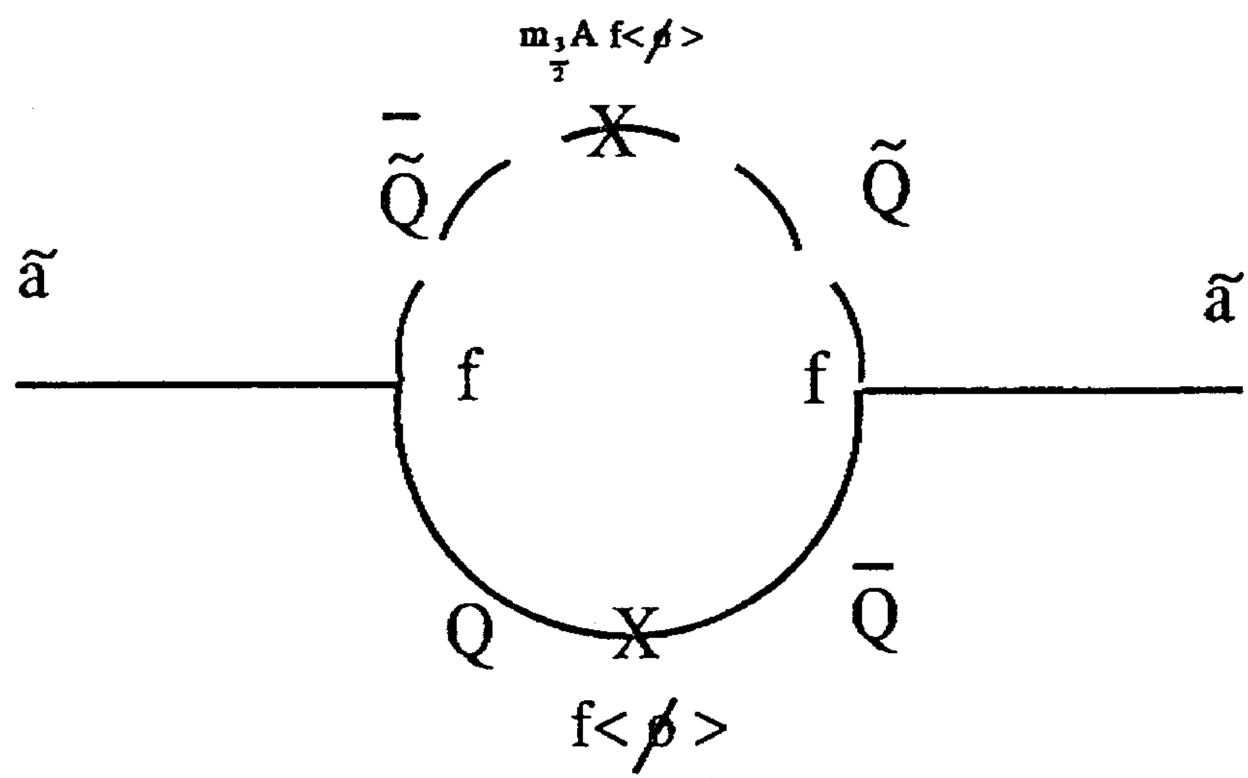
FIGURE CAPTIONS

1. Feynman diagram for the radiative correction which generates a large axino mass in type I axino models.
2. Feynman diagram for a process maintaining the thermal equilibrium abundance of axinos at high temperature.



$m_{3/2} A f\langle\phi\rangle$
$\bar{\tilde{Q}}$
$\tilde{Q}$
$\tilde{a}$
$\tilde{a}$
f
f
Q
$\bar{Q}$
$f\langle\phi\rangle$

FIGURE 1

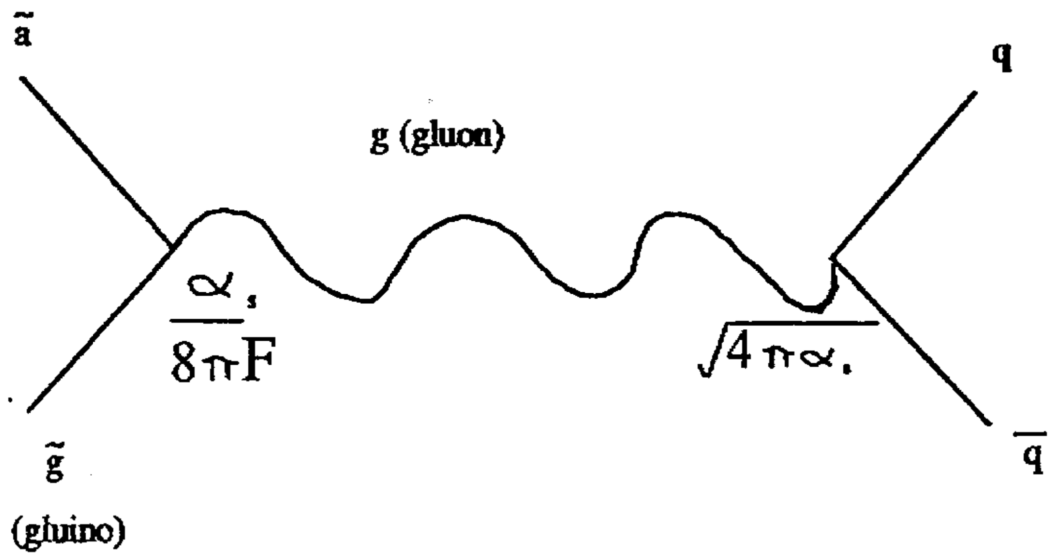


FIGURE 2