INVISIBLE A^0 (AXION)

As discussed in the note on "Axions and Other Light Bosons," the so-called invisible axion models decouple the scale of the Peccei–Quinn symmetry breaking from the electroweak scale, and avoid the constraints from negative accelerator searches for the axion. Two classes of models are discussed commonly in the literature. One introduces new heavy quarks that carry Peccei–Quinn charge while the usual quarks and leptons do not (KSVZ axions or "hadronic" axions) [1]. The other does not need additional quarks but requires two Higgs doublets, and all quarks and leptons carry Peccei–Quinn charges (DFSZ axions or "GUT" axions) [2]. All models contain at least one electroweak singlet scalar boson that acquires an expectation value and breaks the Peccei–Quinn symmetry.

The common property of all axion models is the effective coupling

$$\mathcal{L} = \left(\theta_{\mathit{eff}} - \frac{\phi_A}{f_A}\right) \frac{g_s^2}{32\pi^2}\, F^{\mu\nu a}\, \widetilde{F}^a_{\mu\nu} \;, \tag{1}$$

where ϕ_A is the axion field, θ_{eff} is the effective QCD vacuum angle after the diagonalization of the quark masses, g_s the QCD coupling constant, $F^{\mu\nu a}$ the gluon field strength and $\widetilde{F}^a_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}\, F^{\rho\sigma a}$. It is often convenient to *define* the axion decay constant f_A with this Lagrangian [3]. The QCD instanton effect induces a potential for ϕ_A whose minimum is at $\phi_A = \theta_{\mathit{eff}} f_A$, cancelling θ_{eff} and solving the strong CP problem. The mass of the axion is related to f_A by

$$m_A = 0.62 \times 10^{-3}\text{eV} \;\times (10^{10}\text{GeV}/f_A)\;. \tag{2}$$

The constraints on the axion mass from various experiments, astrophysics, and cosmology are derived from the interactions of the axion with either photons, electrons, or nucleons. We use the following notation for their coupling constants $G_{A\gamma\gamma}$ and G_{Aff},

$$\mathcal{L}_{A\gamma\gamma} = -\frac{1}{4} G_{A\gamma\gamma}\phi_A\, F^{\mu\nu}\, \widetilde{F}_{\mu\nu} = G_{A\gamma\gamma}\,\phi_A \boldsymbol{E}\cdot\boldsymbol{B}\;, \tag{3}$$

$$\mathcal{L}_{Aff} = G_{Aff}\, \partial_\mu\phi_A\, \overline{f}\gamma^\mu\gamma_5\, f\;, \tag{4}$$

Table 1: The coupling constants of the axion to the matter particles in DFSZ and KSVZ models, taken from Ref. [5] where the results of Ref. [6] were used. These dimensionless coupling constants are related to those in the Lagrangian by $G_{Aii} = c_i/2f_A$ for $i = \gamma, e, p, n$. The parameter β is an arbitrary angle whose tangent is defined by the ratio between the expectation values of the two Higgs doublets in the DFSZ model. A rational number E/N in the KSVZ model depends on the number of new quarks and their charges. The coupling to nucleons are subject to certain ambiguities in hadronic matrix elements [4] that are not shown here. All entries have small uncertainties from the current quark masses.

	DFSZ	KSVZ
c_γ	0.0017	$0.0023(E/N - 1.92)$
c_e	$(1/3)\cos^2\beta$	0
c_p	$-0.10 - 0.45\cos^2\beta$	-0.39
c_n	$-0.18 + 0.39\cos^2\beta$	$+0.04$

for $f = e, p, n$. The relations of these coupling constants to f_A (and m_A) are model dependent, and are listed in Table 1.

For illustrative purposes, we depict various constraints on f_A (and m_A) for the case of the KSVZ model in Fig. 1, using only representative constraints. What follows is a brief discussion of each of the constraints shown in the figure. The bounds on the DFSZ axion are similar.

Astrophysics puts a lower bound on f_A, because a small f_A leads to a large coupling of the axion to nucleons, electrons, and photons and thus to a large "exotic" energy-loss rate. In horizontal-branch (HB) stars, the Primakoff process $\gamma + (^4He, e^-) \rightarrow (^4He, e^-) + A^0$ would be the dominant axionic energy-loss mechanism. It would accelerate the consumption of nuclear fuel and thus shorten the helium-burning lifetime of these stars. The observable number fraction of HB stars in globular clusters would be significantly reduced relative to theoretical expectations unless $G_{A\gamma\gamma} < 0.6 \times 10^{-10}$ GeV^{-1} [5]. The duration of the neutrino burst from supernova (SN) 1987A observed at the Kamiokande and IMB detectors was

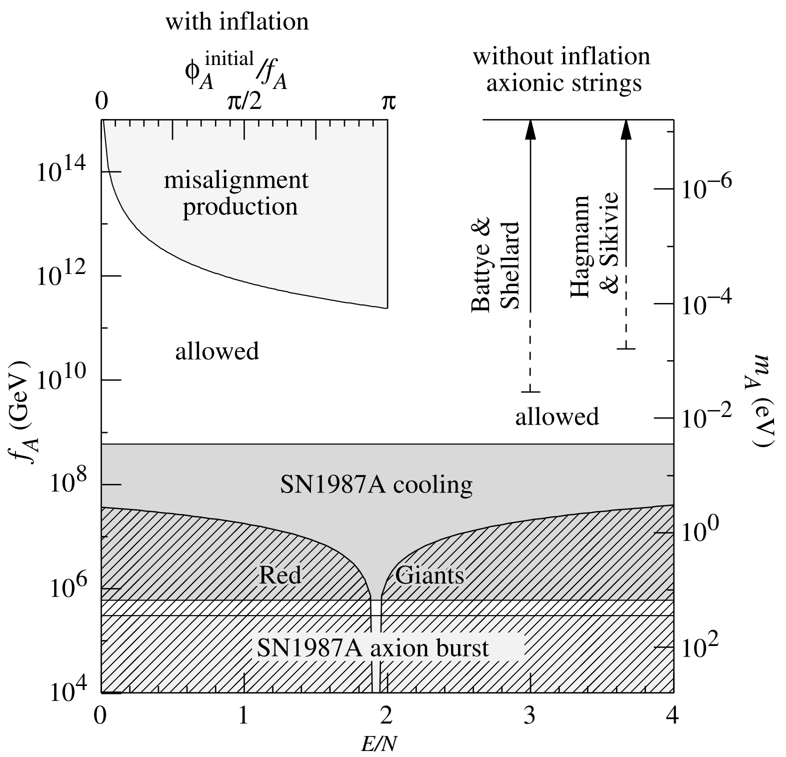

Figure 1: An illustration of the astrophysical and cosmological constraints on the axion decay constant f_A (and equivalently on m_A) in the KSVZ model. The constraint on the DSFZ model is similar except for the small window at $f_A \sim 10^6$ GeV, and one needs inflation. Shaded regions are excluded based on the arguments given in the text, though both sides have large uncertainties. The lower bound on f_A from red giants depends on the parameter E/N. All the other constraints do not. If there is inflation, there is an upper bound on f_A from misalignment production, which depends on the initial value of the axion field ϕ_A^{initial} after inflation is over. If there is not, cosmic strings generated by Peccei–Quinn symmetry breaking produce axions, which contribute to the present mass density. Estimates of the resulting mass density vary. Here, upper bounds on f_A from two groups are shown, each of them with rather large ambiguity shown by dashed lines.

consistent with expectations, while axion emission would have cooled the core and shortened the burst duration [7]. The dominant emission process is axion bremsstrahlung in nucleon-nucleon collisions, a process that needs to be calculated in a hot and dense nuclear medium where many-body effects are important. Early calculations overestimated the emission rate to some extent. A more realistic treatment leads to a somewhat diminished limit of about $f_A > 0.6 \times 10^9$ GeV [8], [9], although the treatment of many-body effects in this result is still under study. The cooling argument does not exclude $f_A \lesssim 0.6 \times 10^6$ GeV; in this range of f_A, the axions produced in the SN core are trapped [10]. Still, $f_A \lesssim 0.3 \times 10^6$ GeV is excluded because the trapped axions result in a burst similar to that of the neutrinos and can produce signals in water Cherenkov detectors [11]. For KSVZ axions with $E/N \approx 2$, there exists an apparent small window between these two SN arguments.

Cosmology usually puts an upper bound on f_A, because the predicted cosmic mass density in axions is proportional to $f_A^{1.175}$. The DFSZ model and the KSVZ model with more than one new quark leads to domain walls that have to be diluted away by inflation. On the other hand, the axion field does not know during inflation where the true minimum of its potential is, and is "misaligned" [12]. It begins a coherent oscillation from its misaligned initial value after the QCD phase transition and contributes to the present energy density [13] as $\Omega_a h^2 \simeq 0.2(f_A/10^{12}\ \text{GeV})^{1.175}(\phi_A^{\text{initial}}/2\pi f_A)^2 \leq 1$ for a small misalignment $\phi_A^{\text{initial}}/f_A \lesssim 1$. The KSVZ model with a single new quark does not produce domain walls and does not need inflation. Without inflation, there are cosmic strings created at the time of the Peccei–Quinn symmetry breaking, that emit axions and eventually decay (or collapse). There is an ongoing controversy on the estimate of the relic energy density of the emitted axions [14]. Furthermore, cosmological bounds change if there is additional entropy production [15] or a dissipation of the coherent oscillation into lighter particles [16].

It has been widely argued that a fundmental theory will not possess global symmetries; gravity, for example, is expected to violate them. Global symmetries such as baryon number typically arise as an indirect consequence of gauge symmetries and renormalizability (accidental symmetry). It has been noted [17] that the Peccei-Quinn symmetry, from this perspective, must also arise as an accidental one and must hold to an extraordinary degree of accuracy in order to solve the strong CP problem. See, for example, Ref. [17] for a possible resolution to this problem; string theory also provides sufficiently good symmetries (see for a review, [18]).

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