Astrophysical Constraints on the Couplings of Axions, Majorons, and Familons

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The spontaneous breaking of a global symmetry leads to the existence of Nambu-Goldstone bosons, which through their coupling to electrons and/or photons, can transport energy from the cores of stars and affect significantly the course of stellar evolution. We find by following in detail the evolution of stars that if the couplings to electrons and/or photons is too strong, helium never ignites—in contradiction with the observational evidence. Our limits restrict the axion mass to less than 0.01 eV, the familon breaking scale to $> 7 \times 10^9$ GeV, and the triplet Majoron vacuum expectation value to < 9 keV.

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To solve the strong *CP* problem, Peccei and Quinn[1] proposed a global (pseudo) symmetry whose spontaneous breakdown led to the existence of an almost massless (pseudo) Goldstone boson, the axion.[2] The original axion, however, interacted too strongly and had a mass and lifetime which were excluded by laboratory experiments[3] as well as astrophysical data.[4] Dine, Fischler, and Srednicki (DFS)[5] modified the Peccei-Quinn proposal to yield a much more weakly interacting—"invisible"—axion which was not in violation of laboratory[3] and/or astrophysical[4] constraints.

The couplings of the DFS invisible axion to electrons and photons were constrained by comparing the axion emission rate to the nuclear-energy generation rate in red giant stars and the Sun.[6] The axion-quark couplings have been constrained by considerations of the cooling rates of neutron stars.[7] Kim[8] independently introduced a new, more weakly coupled axion whose couplings have been constrained by astrophysical considerations.[9]

Other weakly coupled Nambu-Goldstone bosons are the familon, associated with the spontaneous breakdown of a global family symmetry[10] and several varieties of majorons associated with different schemes for breaking a global lepton-number symmetry.[11] For a review of Nambu-Goldstone bosons, see Gelmini, Nussinov, and Yanagida.[12]

In this Letter we report the results of a detailed steller-evolution calculation which improves—by an order of magnitude or more—the constraints on the couplings to electrons and to photons of a general class of Nambu-Goldstone bosons. The previous astrophysical constraints were derived in a linear approximation to the overall energetics of steller systems. In reality, stellar evolution is a coupled nonlinear problem in which feedback plays an important role. Our fully consistent, multizoned, stellar-evolution calculations are capable of yielding far stronger constraints on the allowed couplings of a variety of Goldstone bosons. The reason for this additional leverage is that a small emissivity in Goldstone bosons at a critical zone in a star, (or, at a critical epoch in its evolution) can alter the structure (and/or subsequent evolution) while leaving the gross energetics virtually unaffected.

To set the stage for our discussion of the effects on stellar structure and evolution of Nambu-Goldstone bosons, we first describe the standard results obtained in their absence. Next, we define the couplings of a general class of Goldstone bosons to electrons and to photons and outline the details of our calculations. We then derive new limits on these couplings from considerations of the Sun and, for our best limits, red giant stars. We concluded by noting that our constraints imply that Nambu-Goldstone bosons can only have a negligible effort on the structure and evolution of main sequence stars.

The standard theory of stellar evolution, while neglecting several potentially important phenomena such as rotation and magnetic fields, reproduces many of the observed features relating to the structure and evolution of stars. For a star similar to the Sun, this theory predicts formation in a relatively short collapse phase ($\sim 10^7$ yr) which corresponds to the observed T Tauri stars. The core temperature increases until fusion of hydrogen to helium occurs. During this main-sequence stage, which lasts for more than 90% of the star's life (~ 10 Gyr for a star like the Sun), hydrogen is converted to helium. When the hydrogen is exhausted, the star expands to become a red giant star reaching a luminosity some 2000 times that of the Sun. Such stars will typically have a helium core of $\sim 0.4 M_\odot$. This maximum luminosity for the giant branch is reached just prior to helium ignition for a wide range of masses below $2M_\odot$.[13] Furthermore, comparison of calculated temperature and luminosity distributions are in good agreement with observations of M67 and NGC188 (two clusters with evolved stars of mass near $M_\odot$).[14] The red-giant envelope is convective and material which has been processed through nuclear reactions in the core is dredged up to the surface; the predictions of the standard theory agree well with the observed surface abundances of red giant stars. Eventually, the helium core gets sufficiently hot to ignite helium. During this helium-burning phase, which for solar-mass stars lasts a few times 10^8 yr, the luminosity is some 40 times that of the Sun. The observed ratio of the numbers of main sequence (e.g., H-burning) stars to the He-burning giants is in good agreement with the ratio of lifetimes predicted for these stages.

Our strongest constraints will follow from the simple requirement that red giant stars ignite helium; such a requirement is unaffected by the solar-neutrino problem. We will find that even a small emissivity in light Nambu-Goldstone bosons will cool the cores of red giant stars sufficiently rapidly to prevent He ignition, thus leading to a contradiction with the observed luminosities and numbers of red giants.

Low-mass or massless scalars (s) may be produced in stars throughout their coupling to electrons via the Compton ($\gamma + e \to e + s$) or bremsstrahlung[15] ($e + Z \to e + s$) processes or, through their coupling to photons, via the Primakoff[16] ($\gamma + e, Z \to e, Z + s$) process. For sufficiently weakly interacting scalars, the mean free path will be large compared to a stellar radius; once produced, the scalars will leave the star, carrying away energy and, thus, cooling the interior. The star will respond to this energy drain by contracting and increasing its central temperature, thus raising the rate at which nuclear fuel is burned. A sufficiently high energy-loss rate will have dramatic effects on the central temperature and density of a star, as well as on the stellar lifetime.

As Gelmini, Nussinov, and Yanagida[12] show, Goldstone bosons will, in general, have only pseudoscalar couplings to identical fermions. In the Lagrangean density we may write the scalar-electron coupling as

$$\mathscr{L}_{se} = g_{se} [s(\bar{e} i \gamma_s e)]. \tag{1}$$

The coupling strength g_{se} may be related to the axion mass $[10^{11} g_{ae} \sim (1.4\ \text{eV}^{-1}) m_a]$, the familon breaking scale $(g_{fe} = 2m_e/V)$, or the triplet majoron vacuum expectation value $[10^{11} g_{me} \sim (1.6\ \text{MeV}^{-1}) V_T]$; the singlet majoron of Chikashige, Mohapatra, and Peccei[11] couples to fermions only through higher loops[17]; the Kim[8] axion has no tree-level couplings to fermions.[18] Following Fukugita, Watamura, and Yoshimura[6] and Krauss, Moody, and Wilczek[15] the energy-loss rates as a function of density (ρ) and temperature (T) for the Compton and bremsstrahlung processes are taken to be

$$\epsilon_C = 1.35 \times 10^{-22} g_{se}^2 (1+X) T^6 I(\rho,T) G(\omega_0,T)\ \text{ergs g}^{-1}\,\text{s}^{-1}, \tag{2}$$

$$\epsilon_b = 33.4 g_{se}^2 (1+X) \rho T^{2.5} G_0(\omega_0,T)\ \text{ergs g}^{-1}\,\text{s}^{-1}. \tag{3}$$

The function $I(\rho,T)$ accounts for corrections due to relativistic effects and degeneracy (in their absence, $I=1$); it is well fitted for conditions of interest by

$$I(\rho,T) = (1.01 - 2.9 \times 10^{-9} T)[1 + 2 \times 10^{-6}(1+X)\rho]. \tag{4}$$

At high densities, plasma effects suppress the emissivities by the factor $G = \exp(-\hbar\omega_0/kT)$, where $\omega_0^2 = 2\pi \times e^2 m_e^{-1}(1+X)\rho$ is the plasma frequency, and X is the hydrogen mass fraction.

For the general coupling of scalars to photons, it is convenient to write

$$\mathscr{L}_{s\gamma} = (\alpha/m_e) C_{s\gamma} [s(F\tilde{F})_{\text{em}}] = (4\alpha/m_e) C_{s\gamma} [s(\mathbf{E}\cdot\mathbf{B})]. \tag{5}$$

For the DFS axion,[5] $10^{12} C_{a\gamma} \approx (2.4\ \text{eV}^{-1}) m_a$; it appears to us that the value of $C_{a\gamma}$ used by Fukugita, Watamura, and Yoshimura[6] is too small by a factor of 2 ($C_{a\gamma}^{\text{FWY}} = \frac{1}{2} C_{a\gamma}$), and accordingly we have multiplied their rate by a factor of 4. The emissivity for the Primakoff process was taken to be

$$\epsilon_P = 1.84 \times 10^{-6} C_{s\gamma}^2 (3+X) T^4 J(\rho,T) G(\omega_0,T)\ \text{ergs g}^{-1}\,\text{s}^{-1}. \tag{6}$$

For $T > m_s$ and ω_0, J is well approximated by[6]

$$J(\rho,T) = 1 - 0.72 \ln[|(\hbar\omega_0)^2 - (m_s C^2)^2|^{1/2}/kT]. \tag{7}$$

In Table I we list the emissivities for a range of temperatures and densities; the entries are normalized to $g_{se} = 1.4 \times 10^{11}$ and $C_{s\gamma} = 2.4 \times 10^{-12}$ (these are the values appropriate to the DFS axion with $m_a = 1$ eV). In gen-

TABLE I. Axion emission for $m_a = 1$ eV.

T	Brems	Compt	Prim
	$x = 1.000$ $\rho = 10^2$		
1.0 (+07)	3.802(-01)	4.756(-02)	7.068(-01)
2.0 (+07)	2.247(+00)	3.086(+00)	1.508(+01)
3.0 (+07)	6.284(+00)	3.459(+01)	8.725(+01)
4.0 (+07)	1.300(+01)	1.896(+02)	2.999(+02)
5.0 (+07)	2.280(+01)	7.029(+02)	7.775(+02)
1.0 (+08)	1.301(+02)	3.777(+04)	1.466(+04)
1.5 (+08)	3.596(+02)	3.446(+05)	8.021(+04)
2.0 (+08)	7.393(+02)	1.459(+06)	2.695(+05)
	$x = 1.000$ $\rho = 10^4$		
1.0 (+07)	1.725(+01)	2.243(-02)	1.773(-01)
2.0 (+07)	1.514(+02)	2.161(+00)	4.401(+00)
3.0 (+07)	4.820(+02)	2.763(+01)	2.579(+01)
4.0 (+07)	1.067(+03)	1.618(+02)	1.008(+02)
5.0 (+07)	1.947(+03)	6.239(+02)	2.930(+02)
1.0 (+08)	1.202(+04)	3.629(+04)	7.068(+04)
1.5 (+08)	3.412(+04)	3.399(+05)	4.279(+04)
2.0 (+08)	7.107(+04)	1.449(+06)	1.508(+05)
	$x = 1.000$ $\rho = 10^6$		
1.0 (+07)	6.369(-01)	3.982(-05)	6.546(-05)
2.0 (+07)	2.909(+02)	1.996(-01)	8.456(-02)
3.0 (+07)	3.464(+03)	9.520(+00)	1.850(+00)
4.0 (+07)	1.478(+04)	1.078(+02)	1.216(+01)
5.0 (+08)	4.007(+04)	6.173(+02)	4.604(+01)
1.0 (+08)	5.455(+05)	7.914(+04)	1.773(+03)
1.5 (+08)	2.014(+06)	9.648(+05)	1.020(+04)
2.0 (+09)	4.787(+06)	4.693(+06)	4.401(+04)

eral, bremsstrahlung is only important in low-mass stars with high densities and low temperatures. For the relatively low densities and intermediate temperatures characteristic of solar-type stars, Primakoff emission will be important (provided that there is a direct scalar-photon coupling). At the high temperatures relevant for massive stars, hydrogen-shell-burning stars, and all helium-burning stars, Compton emission will dominate.

The energy-loss rates described above were incorporated in a stellar-structure code derived from that of Eggleton.[19] This code which has successfully modeled stars from $0.2M_\odot$ to $> 100M_\odot$, was used to follow the evolution of stars of various masses near $M_\odot$.

In following the evolution of stars which can cool by emitting weakly coupled Nambu-Goldstone bosons, we find that helium ignition provides the strongest constraint on the couplings of scalars to electrons and/or photons. If the scalar-electron couplings are too strong, Compton emission cools the helium in the region behind the hydrogen-burning shell and insulates the core from conductive heating. With helium ignition suppressed, the star continues to evolve up the giant branch, increasing its luminosity—in excess of that observed. One may set crude limits by using the observed main-sequence parameters of the Sun ($L_\odot = 3.9\times 10^{33}$ ergs/sec, $\tau_\odot = 4.55\times 10^9$ yr). However, these limits were much weaker than those described below.

For the DFS axion[5] we find a barely acceptable stellar model if $m_a = 0.01$ eV. This model has a maximum luminosity nearly a factor of 2 larger than that observed; it does, however, ignite helium. The axion luminosity at the time of ignition was only 0.003 of the total luminosity, but nearly 100 times the neutrino luminosity. It should be noted that our limit, $m_a \leq 0.01$ eV, is likely conservative. Racine[20] claims an uncertainty in the absolute magnitude of the cluster M67 at ± 0.13 magnitude or less than a 15% uncertainty. Not only is the maximum luminosity for $m_a = 0.01$ too large but, we *assumed* that the helium-shell flashes in this model will continue and succeed in propagating into the core; for somewhat stronger coupling, the flashes died out and the helium failed to ignite.

Our limit on m_a translates into a more general constraint on g_{se} (the Primakoff process for axions in the temperature and density regime of interest is negligible). For helium ignition,

$$g_{se} \leq 1.4\times 10^{-13}. \tag{8}$$

To constrain the scalar-photon coupling we set $g_{se} = 0$ and turned $C_{s\gamma}$ down until helium ignition occurred. In this manner we find

$$C_{s\gamma} \leq 2.4\times 10^{-13}. \tag{9}$$

For the case of the DFS axion[5] this would correspond to $m_a \leq 0.1$ eV.

With the upper limits on g_{se} and $C_{s\gamma}$ obtained from the requirement that helium ignition occur (*and* that the luminosity at the tip of the giant branch not exceed that observed), we have followed the main sequence evolution of stars in the mass range $0.2M_\odot$ to $100M_\odot$. The scalar emissivity was negligible compared to the stellar luminosity over the entire range, varying from $L_s/L_{\rm tot} = 1\times 10^{-4}$ for a $0.2M_\odot$ star down to $L_s/L_{\rm tot} = 8\times 10^{-8}$ for a $100M_\odot$ star. Since this ratio provides a measure of the enhancement in the rate at which hydrogen is consumed, the stellar lifetimes are unaltered; the estimates of the ages of the oldest stars[21] are unchanged. As a corollary, the weak coupling required by helium ignition ensures that the standard model of the Sun[22] is unaffected by the existence of Nambu-Goldstone bosons.

The above constraints on g_{se} and $C_{s\gamma}$ are of significance for various theories of spontaneously broken global symmetries. For the general pseudoscalar coupling of Goldstone bosons to electrons[12]

$$g_{se} = 2m_e/V. \tag{10}$$

Our limit, $g_{se} \leq 1.4\times 10^{-13}$, corresponds to $V > 7\times 10^9$ GeV. For example, for the familon model, the family-symmetry-breaking scale must exceed 7×10^9 GeV. Our result also contains the triplet vacuum expectation value in the majoron model of Gelmini and Roncadelli[11] (the majoron of Chikashige, Mohapatra, and Peccei[11] has no tree-level couplings to electrons):

$$V^{-1} = \sqrt{2}G_F V_T; \quad V_T < 9 \text{ keV}. \tag{11}$$

Care must be taken in translating our results into limits on the invisible axion. In particular, the value of the axion decay constant, f_a, depends on the normalization of the axion current; most of the apparent discrepancies between various papers in the literature are traceable to this fact. We have followed the conventions of Bardeen and Tye[23] (BT) which, for three generations, yield $[m_a/(1\text{ eV})][f_a/(10^7\text{ GeV})] = 3.6$. Fukugita, Watamura, and Yoshimura[6] and Sikivie[24] have the same normalization. However, for Srednicki,[6] $f_a = 2f_a(\text{BT})$ and, for DFS,[5] Kaplan,[24] and Krauss *et al.*,[24] $f_a = \frac{1}{2}f_a(\text{BT})$. With the BT normalization, the DFS[5] invisible-axion couplings to electrons and photons (for three generations) are

$$g_{ae} \approx m_e/f_a \approx (1.4\times 10^{-11}\text{ eV}^{-1})m_a, \tag{12a}$$

$$C_{a\gamma} \approx (8.8\times 10^{-5}\text{ GeV}^{-1})f_a^{-1} \approx (2.4\times 10^{-12}\text{ eV}^{-1})m_a. \tag{12b}$$

Our limits require $m_a \leq 0.01$ eV and $f_a \geq 4\times 10^9$ eV.

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