Red-Giant Evolution, Metallicity, and New Bounds on Hadronic Axions

W. C. Haxton and K. Y. Lee
Department of Physics, FM-15, University of Washington, Seattle, Washington 98195
(Received 31 January 1991; revised manuscript received 2 May 1991)

We explore stellar cooling by nuclear axion emission, identifying those special isotopes that dominate this process for temperatures $\sim 10^7$–10^9 K. We argue that such nuclear energy-loss mechanisms are distinctive because the effects track metallicity. Three observables associated with evolution of stars along the red-giant and horizontal branches are shown to impose new and restrictive constraints on axions in the hadronic window.

PACS numbers: 97.20.Li, 14.80.Gt, 95.30.Cq, 97.10.Cv

The elegant solution to the strong *CP* problem proposed by Peccei and Quinn[1] has the consequence that a new light pseudoscalar, the axion, should arise from spontaneous symmetry breaking.[2] A variety of laboratory searches and astrophysical arguments restrict the axion mass to two possible ranges, the hadronic axion window of $1 \lesssim m_a \lesssim 5$ eV and a window of cosmological interest, $10^{-6} \lesssim m_a \lesssim 10^{-3}$ eV.[3,4] The first range exists only for axions of the type discussed by Kim, where the coupling to electrons vanishes at tree level.[5]

The lower bound on the hadronic axion window results from studies of SN 1987A.[6,7] Axions produced by nucleon-nucleon bremsstrahlung can contribute to core cooling, thereby reducing the energy and duration of the competing neutrino burst. As the burst from SN 1987A was consistent with conventional supernova theory, dramatic axion cooling did not occur. Thus either the axion is so weakly coupled that few are produced ($m_a \lesssim 10^{-3}$ eV) or so strongly coupled that the trapping is comparable to or greater than that for neutrinos. The latter lower "ledge" on the hadronic axion window corresponds roughly to $m_a \sim 3 \pm 3$ eV, with the uncertainty due in part to approximations in supernova modeling and in part to the absence of a precise relationship between m_a and the strength of the aNN coupling.

The upper ledge of the hadronic axion window is the bound imposed by the helium-burning lifetimes of red-giant stars.[8] While the absence of a direct axion-electron coupling lessens the energy loss in the case of hadronic axions, Primakoff production still leads to the interesting constraint

$$m_a \lesssim (1\ \text{eV}) \left| \frac{0.72}{E/N - 1.95} \right| .$$

This limit is somewhat uncertain due to the dependence of the underlying $a\gamma\gamma$ coupling on the ratio of the coefficients of the electromagnetic and color anomalies, E/N. While in the simplest theories $E/N = 8/3$ (so that $m_a \lesssim 1$ eV), other choices yield a weaker mass limit. For instance, Kaplan[9] constructed a model where $E/N = 2$, which would then imply $m_a \lesssim 14$ eV.

In this Letter we explore nuclear axion emission in stars, a mechanism that has several interesting virtues. First, it directly tests[10] the coupling of hadronic axions to nucleons with a sensitivity that competes favorably with the red-giant Primakoff process. This is important because, while the coupling to nucleons depends on at least one poorly constrained parameter (the flavor-singlet axial-vector matrix element S), those couplings that lead to weak nuclear emission tend to produce strong supernova limits (the opacity is decreased). Thus it becomes more difficult to pry open the hadronic axion window. We find that there is a substantial region in S where our limits overlap those from SN 1987A. Second, the anomalous energy production is directly tied to stellar metallicity, an observable. Thus one can study variations in bolometric magnitudes or stellar lifetimes as a function of metallicity, constructing sound statistical arguments for ruling out axions of a given mass or coupling. We extract three such observables from Raffelt's recent analysis of red-giant-branch (RGB) and horizontal-branch (HB) stellar evolution.[11] Third, the nuclear physics governing axion emission is strongly constrained by known γ-decay rates.

Kaplan[9] and Srednicki[12] have discussed the couplings of hadronic axions to nucleons. The couplings depend on the symmetric and antisymmetric reduced matrix elements for the octet axial current, which conventionally are determined from neutron and hyperon β decay and flavor SU(3) symmetry. We use the values $F = 0.48$ and $D = 0.77$.[7] Ignoring small corrections[12] depending on the quark mass ratio $w = m_u/m_s$, one finds that the aNN coupling can be written as

$$\mathcal{L} \sim a \bar{N} \gamma_5 (g_0 + g_3 \tau_3) N , \qquad (1)$$

where $g_0 = -1.77 \times 10^{-8} [m_a/(1\ \text{eV})](1 + 2.94S)$ and $g_3 = -2.75 \times 10^{-8} m_a/(1\ \text{eV})$. The flavor-singlet axial-vector matrix element S ($= \Delta u + \Delta d + \Delta s$ in Ref. 7) is poorly constrained. The naive quark model (NQM) gives $S = 0.68$, while an estimate based on the European Muon Collaboration (EMC) measurement of the spin-dependent muoproduction structure function is $S = -0.09$.[7] We explore $-1 < S < 2$ to encompass these results and the broad range mentioned by Kaplan.[9]

Note that the coupling to neutrons, $(g_0 - g_3)/2$, vanishes for $S \sim 0.2$, a value between the NQM and EMC estimates.

In the long-wavelength limit this coupling generates $M1$ axion emission in nuclei. The rate relative to γ decay is[13]

$$\frac{w_a}{w_\gamma} = \frac{1}{2\pi\alpha}\frac{1}{1+\delta^2}\left[\frac{g_0\beta + g_3}{(\mu_0 - \frac{1}{2})\beta + \mu_1 - \eta}\right]^2, \quad (2)$$

where δ is the $E2/M1$ mixing ratio, μ_0 and μ_1 are the isoscalar and isovector magnetic moments, and the nuclear-structure-dependent terms are

$$\eta = -\frac{\langle J_f || \sum_{i=1}^{A} \mathbf{l}(i)\tau_3(i) || J_i \rangle}{\langle J_f || \sum_{i=1}^{A} \boldsymbol{\sigma}(i)\tau_3(i) || J_i \rangle}$$

and

$$\beta = \frac{\langle J_f || \sum_{i=1}^{A} \boldsymbol{\sigma}(i) || J_i \rangle}{\langle J_f || \sum_{i=1}^{A} \boldsymbol{\sigma}(i)\tau_3 || J_i \rangle}.$$

As $\mu_0 - \frac{1}{2} \sim 0.38 \ll \mu_1 \sim 4.71$, the denominator (particularly in the case of strong $M1$ transitions) should be dominated by μ_1. The transitions in odd-A nuclei of present interest are characterized by $\beta \sim \pm 1$ depending on whether the unpaired nucleon is a proton or neutron.

The axion luminosity due to a particular nucleus will depend on the abundance of that metal, the Boltzman factor for populating an excited nuclear state, and the $M1$ strength of the transition to the ground state. The stellar temperatures of interest ($T_8 \sim 1$, or $kT \sim 8.6$ keV) are low on the scale of typical nuclear transition energies: The cooling rate is optimized, for fixed $B(M1)$, at $4kT \sim 35$ keV. From solar abundance tables and tabulated transition energies and γ-decay rates one can estimate the axion emission rates for various metals. Stellar axion emission over the range $0.1 \lesssim T_8 \lesssim 10$ is largely governed by three nuclei, ^{57}Fe, ^{55}Mn, and ^{23}Na. The first is an odd-neutron nucleus, while the latter two are odd-proton nuclei. The characteristics of the relevant $M1$ transition in each nucleus are given in Table I, along with the results of large-basis shell-model calculations of the matrix element ratios β and η.[14] The respective cooling rates (erg/g s) are

TABLE I. Dominant metals, solar abundance by mass fraction, and $M1$ transition properties for hadronic axion cooling, $0.1 \lesssim T_8 \lesssim 10$. (W.u. denotes Weisskopf unit.)

Isotope	Abundance	ΔE (keV)	$B(M1)$ (W.u.)	β	η
^{57}Fe	3.26×10^{-5}	14.4	8.2×10^{-3}	-1.19	0.80
^{55}Mn	1.48×10^{-5}	126	4.3×10^{-2}	0.79	-3.74
^{23}Na	4.00×10^{-5}	440	0.23	0.88	-1.20

$$\delta E = \frac{Z}{Z_\odot}\left[\frac{m_a}{1\text{ eV}}\right]^2 \times \begin{cases} 0.91(1-9.7S)^2 e^{-\beta_1}/(1+2e^{-\beta_1}), & ^{57}\text{Fe},\\ 5.6\times10^4(1+0.99S)^2 e^{-\beta_2}/(1+1.33e^{-\beta_2}), & ^{55}\text{Mn},\\ 7.2\times10^8(1+1.07S)^2 e^{-\beta_3}/(1+1.5e^{-\beta_3}), & ^{23}\text{Na}, \end{cases} \quad (3)$$

where $Z/Z_\odot$ is the metallicity relative to the solar value, $\beta_1 = (14.4\text{ keV})/kT$, $\beta_2 = (126\text{ keV})/kT$, and $\beta_3 = (440\text{ keV})/kT$ (see Fig. 1). Despite the relative weakness of the ^{57}Fe transition, its favorable Boltzman factor produces the largest rate near $T_8 \sim 1$.

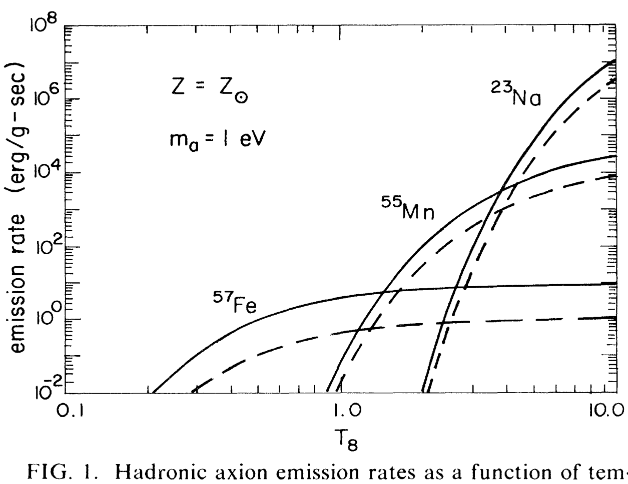

FIG. 1. Hadronic axion emission rates as a function of temperature for $m_a = 1$ eV and $Z = Z_\odot$. The solid (dashed) lines correspond to $S = 0.68$ (-0.09), the NQM (EMC) value.

Raffelt,[11] in deriving an improved constraint on the neutrino magnetic moment, discussed the theory and observational data for red-giant evolution before and after the helium flash. A RG contains a dense helium core ($\rho \sim 10^6$ g/cm^3) that supports itself by electron degeneracy pressure, surrounded by a thin shell, at the interface with the hydrogen envelope, in which hydrogen burning occurs. As additional helium is synthesized and added to the core, the gravitational potential at the burning shell increases, thereby requiring a faster rate of hydrogen burning to maintain the necessary gas pressure. Thus the star brightens. The core density must also increase in response to its greater mass, eventually reaching a point where $3\alpha \rightarrow {}^{12}$C burning can ignite. The resulting "helium flash" at the tip of the red-giant branch results in a quick expansion of the core. The star relaxes into a new mode where helium burns at the center of the helium core while hydrogen shell burning continues at a much reduced rate due to its lower gravitational potential. The overall luminosity of the star is reduced after the helium flash, and the star begins its slow evolution along the horizontal branch.

Raffelt noted that absolute measurements of the brightness of HB stars, M_{RR}, have been made at the RR Lyrae strip (surface temperature $10^{3.85}$ K) by the Baade-Wesselink method.[15] One can then compare this with the stellar evolution prediction for the brightness of stars entering the HB, provided one makes a small correction for the brightness difference between RR Lyrae and zero-age HB stars, Δ_{RR}. In a given globular cluster one can also determine the brightness difference between stars at the tip of the RGB and the RR Lyrae stars, and compare that quantity to stellar evolution results. Finally, one can employ the R method, the comparison between theoretical lifetimes of model stars on the RGB and HB with the observed number ratio, $R=N_{HB}/N_{RGB}=t_{HB}/t_{RGB}$. The resulting equations used in Ref. 11, together with modifications we will discuss below, are

$$M_{RR}-0.16Z^*+7.3\delta M_c=0.59-3.5Y_e^*-\Delta_{RR}+\varepsilon\,, \tag{4a}$$

$$\Delta M_{RR}^{tip}-0.39Z^*-4.0\delta M_c=4.13-4.4Y_e^*-\Delta_{RR}\,, \tag{4b}$$

$$\log_{10}R-0.029Z^*+0.70\delta M_c+0.43\delta_{ax}=0.151+2.29Y_e^*+0.33\Delta_{RR}\,, \tag{4c}$$

where $Y_e^*=Y_e-0.25$, Y_e is the envelope helium abundance, and $Z^*=3+\log_{10}Z$. Thus solar metallicity is $Z_\odot^*=1.3$. A fixed value of $Y_e\sim0.27$ is thought to characterize globular clusters in our galaxy.[11] Measured metallicities for the stars we will consider lie in the range $-0.9\lesssim Z^*\lesssim1.1$. In (4a) ε represents a (possibly large) systematic error[11] in determining the absolute magnitudes M_{RR}: Our analysis only requires a knowledge of *relative* magnitudes.

The quantity δM_c represents the excess helium core mass, in units of $M_\odot$, that would result from anomalous cooling during RG evolution: Cooling delays helium ignition. Raffelt evaluated δM_c for plasmon decay into $\nu\bar\nu$ via a neutrino magnetic moment μ_{12}. His result (in terms of μ_{12}) can be rewritten in terms of the anomalous cooling

$$\delta M_c=0.015(\delta E_{RG})^{1/2}\,, \tag{5}$$

where δE_{RG} is in units of erg/gs, and is given in the present case by Eq. (3), evaluated at $\langle T_8^{RG}\rangle\sim0.78$.[16]

We now depart from Ref. 11 both in the physics and the analysis. As Raffelt's $\gamma\to\nu\bar\nu$ cooling is $\sim\rho$ for $\rho\lesssim10^6$ g/cm^3, the anomalous losses cease when the star enters the HB. In contrast, the nuclear axion emission depends only on temperature and is therefore somewhat stronger on the HB [$\langle T_8^{HB}\rangle\sim1.25$ (Ref. 16)]. The HB lifetime shortens in proportion to the axion losses,[17]

$$\delta\ln t_{HB}=\ln\left(1+\frac{\delta t_{HB}}{t_{HB}^0}\right)\sim\frac{\delta t_{HB}}{t_{HB}^0}=-\frac{\delta E_{HB}}{E_{3\alpha}}\sim-7.1\times10^{-3}\delta E_{HB}\equiv-\delta_{ax}\,, \tag{6}$$

where t_{HB}^0 is the HB lifetime in the absence of anomalous energy loss, $E_{3\alpha}$ is the 3α burning rate, and δE_{HB} is given by Eq. (3). Thus this term appears in Eq. (4c). Equations (4) will also be affected by associated HB luminosity changes. Frieman, Dimopoulos, and Turner[17] argue that the fractional increase in the helium-burning contribution to the HB luminosity is $\sim\delta_{ax}/\nu$, where $E_{3\alpha}\sim T^\nu$ with $\nu\sim30$–40. The energy-generation rate in the hydrogen-burning shell will also increase as the shell radius shrinks in response to core axion losses, increasing the gravitational potential, but again[17] $\delta R/R\sim\delta_{ax}/\nu$. Thus we tentatively assume both effects are sufficiently small to be ignored, with the understanding that future calculations are needed to verify this.

In the analysis of Ref. 11 the metallicity dependences of Eqs. (4) were eliminated by averaging over various globular clusters, and the resulting equations combined to eliminate Y_e^* and Δ_{RR}, thereby isolating δM_c. Because nuclear axion emission is explicitly dependent on metallicity, our analysis is very much simplified: We will obtain three independent constraints on m_a by considering *only the relative metal dependence* for each of the variables in Eqs. (4). For example, given some hypothesis about S and m_a, stellar evolution then predicts that $M_{RR}-0.16Z^*+7.3\delta M_c$ is a random variable independent of metallicity. We can subject this hypothesis to a standard variance test: Using N measurements of M_{RR} with specified (statistical) errors, we evaluate

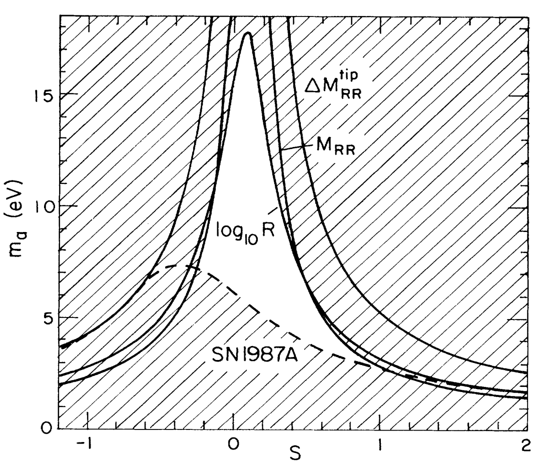

FIG. 2. Mass limits on hadronic axions from the metallicity dependence of M_{RR}, $\log_{10}R$, and ΔM_{RR}^{tip} as a function of S. The shaded regions above the solid lines are excluded at 99% C.L. The region below the dashed line is excluded by SN 1987A, according to the parametrization of Ref. 11.

the weighted mean $\langle M_{\rm RR} - 0.16Z^* + 7.38\delta M_c \rangle$ and the weighted variance from the scatter of the data about the mean. This is an unbiased estimate of the known statistical (i.e., experimental) variance and should be distributed as χ^2 with $N-1$ degrees of freedom. Thus one can increase m_a, for a given choice of S, until the calculated variance allows us to reject our hypothesis with a confidence of 99%. More details of this analysis, which remains valid even if Primakoff production off He or other metal-independent anomalous cooling mechanisms are operating, will be given elsewhere.[14]

Using the data of Tables II, IV, and V of Refs. 11 and 15, we obtain the hadronic axion mass limits shown in Fig. 2. The most restrictive of the constraints are those from $\log_{10}R$ and $M_{\rm RR}$. For the HB cooling leads to a linear dependence or energy loss ($\propto Zm_a^2$) through $\delta_{\rm ax}$, compared to the gentle $\log_{10}Z$ dependence that arises from stellar evolution. The variable $M_{\rm RR}$ tests $m_a\sqrt{Z}$ and the corresponding limit is impressive in view of the limited data set (seven RR Lyrae stars over a limited metallicity range). Because our analysis is statistical, all three limits can be improved by expanding the data bases. Our 99%-C.L. limits merge with those from the SN 1987A, as parametrized in Ref. 11, for large and small S at points, coincidentally, that roughly correspond to the NQM and EMC values for S.

As the ^{57}Fe transition is approximately neutron in character, the coupling vanishes for $S \sim 0.1$. Thus, for a band in S, a window appears between our limits and that from SN 1987A. The $\log_{10}R$ curves closes at $m_a \sim 18$ eV because ^{55}Mn, with its odd-proton transition, is of some importance at HB temperatures. At only slightly higher temperatures, $T_8 \sim 1.5$, both ^{55}Mn and ^{57}Fe would contribute equally, eliminating the window. Thus it would be interesting to identify stellar conditions where the much stronger transitions in ^{55}Mn or ^{23}Na could be exploited, and where similar metallicity arguments could be formulated. One possibility would be a comparison of stars on the HB and asymptotic giant branch.

We are indebted to G. Raffelt, A. Burrows, C. Hogan, and G. Wallerstein for a number of helpful comments. This work was supported in part by the U.S. Department of Energy, the National Science Foundation (No. PHY8912689), and NASA (No. ATP-90-103).

[1]R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. **38**, 1440 (1977).

[2]S. Weinberg, Phys. Rev. Lett. **40**, 223 (1978); F. Wilczek, *ibid.* **40**, 279 (1978).

[3]G. G. Raffelt, Phys. Rep. **198**, 1 (1990).

[4]M. S. Turner, Phys. Rep. **197**, 67 (1990).

[5]J. E. Kim, Phys. Rev. Lett. **43**, 103 (1979); M. S. Shifman, A. J. Vainstein, and V. I. Zakharov, Nucl. Phys. **B166**, 493 (1980).

[6]A. Burrows, M. T. Ressell, and M. S. Turner, Phys. Rev. D **42**, 3297 (1990).

[7]R. Mayle *et al.*, Phys. Lett. B **203**, 108 (1988); **219**, 515 (1989); G. Raffelt and D. Seckel, Phys. Rev. Lett. **60**, 1793 (1988); M. S. Turner, *ibid.* **60**, 1797 (1988).

[8]G. G. Raffelt and D. S. P. Dearborn, Phys. Rev. D **36**, 2211 (1987); M. Fukugita, S. Watamura, and M. Yoshimura, *ibid.* **26**, 1840 (1982); D. A. Dicus *et al. ibid* **22**, 839 (1980).

[9]D. B. Kaplan, Nucl. Phys. **B260**, 215 (1985).

[10]This is also a virtue of Kamioka II analysis of J. Engel, D. Seckel, and A. C. Hayes, Phys. Rev. Lett. **65**, 960 (1990).

[11]G. G. Raffelt, Max Planck Institut Report No. MPI-PAE/PTH 6/90 (to be published); G. G. Raffelt, Phys. Rev. Lett. **64**, 2856 (1990).

[12]M. Srednicki, Nucl. Phys. **B260**, 689 (1985).

[13]F. T. Avignone, III, *et al.*, Phys. Rev. D **37**, 618 (1988).

[14]W. C. Haxton and K. Y. Lee (to be published). Note that the Primakoff process does produce a metallicity-dependent contribution to the cooling, which we neglect in order to simplify the particle physics (eliminating any dependence on E/N). For $Z^* \sim 1.1$, the highest metallicity we consider, 6% of the RGB and initial HB cooling would result from isotopes other than He (principally ^{16}O, ^{56}Fe, and ^{12}C). Since the criterion for unacceptable Primakoff losses is usually taken to be 100 erg/gs, our new constraints are interesting in the region where the Z^*-dependent Primakoff losses are $\lesssim 6$ erg/gs. In comparison, our 99%-C.L. limits correspond to nuclear axion emission rates of ~ 35 erg/gs for HB and RGB stars with $Z^* \sim 1.1$.

[15]R. V. Jones, B. W. Carney, and D. W. Latham, Astrophys. J. **332**, 206 (1988).

[16]A. V. Sweigart and P. G. Gross, Astrophys. J. Suppl. **36**, 405 (1978); Astrophys. J. **190**, 101 (1974).

[17]J. A. Frieman, S. Dimopoulos, and M. S. Turner, Phys. Rev. D **36**, 2201 (1987).