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AXIONS AND STARS*

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ABSTRACT

The emission of light, noninteracting particles, such as axions and majorons, modifies the structure and evolution of stars. We show that the main effects of such an energy loss are to raise the central temperature and luminosity of stars and to reduce their main sequence lifetimes. The concordance of the standard model of the sun with observations yields a self-consistent bound on the coupling of light pseudoscalars to electrons, $g_{Be} < 1.6 \times 10^{-11}$. This corresponds to a lower limit on the Peccei-Quinn scale, $F \geq 3.2 \times 10^7 GeV$, a factor of 3 better than the usually quoted solar bound. We briefly discuss the consequences of axion emission for the solar neutrino problem and for the ages of globular clusters. Since the axion luminosity scales as a lower power of temperature than the nuclear energy generation rate, axion emission has a stronger influence on cool stars. We discuss the resultant changes in the mass-luminosity relation for low mass stars and the effects on the low-mass cut-off for main-sequence stars.

I. INTRODUCTION

In recent years, a number of particle physics models have been proposed in which an approximate global symmetry is spontaneously broken at a high energy scale.[1] These models predict the existence of light, very weakly interacting Nambu-Goldstone bosons.[2] The paradigm case is that of the axion, while other candidates include the majoron and the familon. (Throughout this paper, we will use axion as a synonym for Nambu-Goldstone boson.) Although inaccessible to conventional terrestrial experiments, axions may have significant astrophysical signatures. If their mass is much less than stellar core temperatures ($m_{axion} \ll kT_c \sim 1keV$), axions can be copiously produced in stars by atomic processes.[3–5] Since they interact so weakly, axions have mean free paths much larger than stellar radii; once created in stars, they escape unscathed.

Axion emission drains energy from stars and thereby perturbs stellar structure and evolution. Since the structure of main sequence stars, and of the sun in particular, is well understood, the agreement of stellar models with observations can constrain such emissions. Thus stars play the dual roles of axion factories and axion laboratories.

In this paper, we consider in detail the perturbative effects on main sequence stars of the emission of light noninteracting particles, such as axions. (In fact, our approach is more general than this and applies to arbitrary perturbations in the stellar energy generation rate.) Our analytic approach, based on methods developed by Stromgren and Chandrasekhar,[6] is developed in Sect.II. (Our work is complementary to that of Dearborn, Schramm and Steigman,[7] who recently incorporated axions directly into a stellar evolution code. Although less accurate, the analytic approach makes the physics of the problem transparent.) The usually quoted solar axion bound,[3] obtained by setting the axion luminosity equal to the photon luminosity of the sun, is arbitrary and, it turns out, inconsistent. In Sect.III, we derive a stronger, self-consistent, yet conservative bound by directly taking into account the effects of axion emission on the sun. In Sect.IV, we extend the discussion to low mass stars. For these stars, the effects of axion emission are more dramatic, but the prospects for observational constraints appear to be more remote, as the theoretical uncertainties are greater and the observational data more sparse.

II. HOMOLOGOUS STARS

We first briefly review the basic ideas of stellar structure and evolution.[8] Stars spend most of their lifetime in the main sequence phase, during which they burn hydrogen to helium. We focus here on lower main sequence stars, with mass $0.1M_\odot \leq M \leq 1.2M_\odot$, which burn predominantly by the proton-proton chain. These stars spend of order 10^{10} years on the main sequence; the long evolution time implies that their structure can be taken to be static to good approximation. In lower main sequence stars, the stellar core is in radiative equilibrium (i.e., energy transport is predominantly by photon diffusion), surrounded by an outer convective envelope. For stars with mass $M \lesssim 0.3M_\odot$, the convective envelope is thought to extend to the stellar center: these stars are fully convective. We will see that the structure of stars which are convective throughout the interior is qualitatively different from stars in which radiative diffusion carries a significant portion of the energy flux.

The basic equations of stellar structure are the conservation of mass,

$$\frac{dM_r}{dr} = 4\pi r^2 \rho \tag{1}$$

the condition for hydrostatic equilibrium,

$$\frac{dp}{dr} = -\frac{GM_r \rho}{r^2} \tag{2}$$

and the condition for thermal equilibrium,

$$\frac{dL_r}{dr} = 4\pi r^2 \varepsilon \tag{3}$$

Here, M_r is the mass interior to radius r, L_r is the net flux of radiation through a shell of radius r, and ε is the energy generation rate (erg gm^{-1} sec^{-1}). For the rest of this section, we shall consider stars which, like the sun, are in radiative equilibrium throughout most of the interior. In this case, energy transport through the star is governed by the equation of radiative transfer,

$$\frac{dT}{dr} = -\frac{3\kappa\rho L_r}{16\pi r^2 acT^3} \tag{4}$$

where κ is the opacity ($(\kappa\rho)^{-1}$ is the photon mean free path), and a is the Stefan-Boltzmann constant.

The condition of hydrostatic equilibrium implies that large central pressures are required to support stars against gravitational collapse. From the equation of state for matter and radiation, this implies high central temperatures, $T_c \simeq 10^6 - 10^7 K$, and thus large temperature gradients to the relatively cool stellar surface. By Eqn.(4), this gradient drives a net flux of radiation through the star and from its surface, regardless of whether the energy loss is compensated by nuclear energy generation. However, the condition of thermal equilibrium (3) can only be violated for of order the Kelvin-Helmholtz time ($\simeq 10^7 yrs$ for a solar mass star), the timescale for the star to radiate its thermal reserves. Note that the free fall time ($\simeq$ 50 min) is much shorter, so hydrostatic equilibrium cannot be violated for an appreciable time.

This simple picture allows us to understand the qualitative effects of incorporating axions into stars. Axion energy production at a rate ε_{ax} acts as a local heat sink, reducing the effective value of the stellar energy generation rate to $\varepsilon = \varepsilon_{nuc} - \varepsilon_{ax}$, where ε_{nuc} is the energy generation rate from nuclear reactions. Consider a star, initially in thermal equilibrium, in which axion production is adiabatically switched on. As the axion rate builds up, the star suffers a net energy loss and goes out of thermal equilibrium. To restore equilibrium, the star gains gravitational energy by slowly contracting on the Kelvin-Helmholtz timescale. By the virial theorem, however, half of the energy gained in contraction must go into thermal energy, so the temperature rises. Since ε_{nuc} scales as a positive power of T, hydrogen is burned at a faster rate, and the stellar lifetime is reduced. (The fact that heat loss leads to an increase in temperature reflects the negative specific heat of gravitationally bound systems.)

Clearly, this gravitational contraction will eventually cease only if ε_{ax} has a weaker temperature dependence than ε_{nuc}; otherwise, contraction reduces the net energy generation rate, and the star moves further out of equilibrium. For the proton-proton chain in lower main sequence stars, the nuclear rate is approximately[8]

$$\varepsilon_{nuc} \simeq \varepsilon_o \rho T^{\nu} \; ; \; \nu = 4 - 6 \tag{5}$$

with fixed ν in a narrow range of temperature, while for the dominant axion processes, ε_{ax} scales as $\rho T^{2.5}$ (bremsstrahlung)[4] or $\rho T^{1.5}$ (axiorecombination).[5] The Compton and Primakoff processes, while important, never dominate the axion emission rate. Thus a main sequence star with axions can relax to a new stable configuration.

To study the effects of axions in detail, we need an additional simplifying assumption. To compare a star with axions to a hypothetical uncontaminated star, we again imagine that the system is allowed to relax quasi-statically to a new configuration as the axions are 'turned on'.[♯1] In addition, we assume that the resulting contraction of the star is approximately uniform, or *homologous*, that is, "the distance between any two points is altered in the same way as the radius of the configuration".[9] In fact, chemically homogeneous stars are believed to contract in roughly homologous fashion when they go out of thermal equilibrium (e.g., when hydrogen is exhausted at the end of the main sequence phase).

It is straightforward to work out how physical quantities scale under a homologous contraction.[9] Suppose a star contracts from its initial radius R_o to a final radius R_1, with $y = R_1/R_o$. Then a volume element initially at radius r_o is moved to a new position $r_1 = yr_o$; the points r_o and r_1 are said to be homologous. It is convenient to define the homology invariant

$$x = (r_o/R_o) = (r_1/R_1) \tag{6}$$

From Eqn.(1), if the initial density at r_o is ρ_o, then the final density at the

♯1 Of course, axions will be generated *ab initio*, so the star will not, in fact, relax in this way. However, by the Vogt-Russell theorem, the structure of a static main sequence star is uniquely determined by the input parameters, i.e., it is 'path independent', so this picture should give an accurate representation of the effects of axions even if they are nonlinear.

homologous point r_1 is

$$\rho_1(x) = y^{-3}\rho_o(x) \tag{7}$$

Since $M_o(r_o) = M_1(r_1)$, from Eqn.(2) the pressure increases to

$$p_1(x) = y^{-4}p_o(x) \tag{8}$$

Also, since the equation of state for lower main sequence stellar interiors is to good approximation given by the ideal gas law, $p \sim \rho T$, the temperature scales as

$$T_1(x) = y^{-1}T_o(x) \tag{9}$$

and the temperature gradient as

$$\left(\frac{dT}{dr}\right)_1(x) = y^{-2}\left(\frac{dT}{dr}\right)_o(x) \tag{10}$$

Thus, under homologous contraction, the density, pressure and temperature profiles are unchanged aside from a global rescaling.

To proceed further, we recognize that the assumption of homologous contraction imposes restrictions on the constitutive relations for the stellar energy generation rate and the opacity. In particular, for chemically homogeneous stars one can show that[10] these quantities must be homogeneous functions of the density and temperature,

$$\varepsilon = \varepsilon_o\rho^n T^\nu \;;\; \kappa = \kappa_o\rho^s T^p \tag{11}$$

For chemically homogeneous stars in radiative equilibrium, Eqn.(11) turns out to be a good fit to the detailed calculations of the nuclear energy generation rate and opacity. Following Eqn.(5), we will take $n = 1$; for the opacity, we will use the Kramers law, with $s = 1$, $p = -3.5$, which is found to be an accurate interpolation formula throughout most lower main sequence interiors.[8]

From the discussion following Eqn.(5), the axion rate ε_{ax} does not scale in the same way as ε_{nuc}; strictly speaking, this implies a breakdown of homology. However, the temperature dependencies of the two rates are neither very steep nor very different from each other; since the temperature is fairly uniform over the energy producing core of the star, we can to good approximation replace the position-dependent ratio $\varepsilon_{ax}/\varepsilon_{nuc}$ by its average over the stellar core, $\delta_{ax} = \langle\varepsilon_{ax}/\varepsilon_{nuc}\rangle_{core}$. If the axion effects are small, $\delta_{ax} \ll 1$, the departure from homology neglected here should only lead to small higher order corrections to the results below. (Note that in this approximation, the star has neutral stability with respect to axion perturbations.)

From Eqns.(4-11), we find the luminosity scales as

$$L_1(x) = y^{-1/2} L_o(x) \tag{12}$$

On the other hand, with $\varepsilon_1 = \varepsilon_{nuc} - \varepsilon_{ax} = \varepsilon_o(1 - \delta_{ax})\rho T^\nu$, Eqn.(3) becomes

$$L_1(x) = y^{-(3+\nu)}(1 - \delta_{ax}) L_o(x) \tag{13}$$

Combining Eqns.(12) and (13) and linearizing in $y - 1 = (\delta R/R_o)$, we can relate the reduction in stellar radius to the relative axion power,

$$\frac{\delta R}{R} = -\frac{\delta_{ax}}{\nu + 2.5} \tag{14}$$

Using Eqn.(12) and the temperature scaling $T \sim y^{-1}$, we find the increase in stellar luminosity and central temperature,

$$\frac{\delta L}{L} = \frac{\delta_{ax}}{2(\nu + 2.5)} \tag{15}$$

$$\frac{\delta T_c}{T_c} = \frac{\delta_{ax}}{\nu + 2.5} \tag{16}$$

In Eqn.(16), we have used the fact that the stellar center is a homologous point. From Eqn.(15), we see that the luminosity is relatively insensitive to a small perturbation in the energy generation rate ε because the nuclear reaction rates are moderately sensitive to the temperature; this reflects our earlier conclusion that the luminosity is primarily determined by the condition of hydrostatic equilibrium and the equation of radiative transfer.[10]

The evolution time of a star is given by the rate of hydrogen burning, $\tau \sim C/\varepsilon_{nuc}$, where C is the energy released per gram of hydrogen consumed. From Eqns.(6-11) we find $\tau \sim T_c^{-(\nu+3)}$, and thus

$$\frac{\delta \tau}{\tau} = -\frac{(\nu + 3)\delta_{ax}}{(\nu + 2.5)} \tag{17}$$

Eqn.(17) gives the fractional change in the main-sequence lifetime and in the time required to burn to a given core value of the hydrogen mass fraction X_H.

III. THE SUN

In this section, we apply the analysis of Sect.II to obtain bounds on axion emission from the sun and solar mass stars. From Eqns.(14)-(17), we see that the most sensitive parameter is the stellar evolution timescale τ. From meteor dating, the solar lifetime is known to be $\tau_\odot = 4.5 \times 10^9$ years, with an estimated uncertainty $\lesssim 10^8$ years.[11] Current evolutionary models of the sun are able to reproduce quite well the observed luminosity and radius at an age corresponding to $\tau_\odot$.

If axions were included in these models, they would evolve more rapidly, leading to a luminosity and radius at $\tau_\odot$ in conflict with observations. To regain agreement with observations of the sun, the evolution time of these models could be partially compensated by changes in the initial solar composition and in the convective mixing length parameter. To be conservative, we assume that the largest timescale change which could be reasonably accounted for in this way, i.e., the uncertainty in evolutionary lifetime of solar *models*, is of order 4×10^8 years, or

$$\frac{\delta\tau}{\tau} \leq 0.1 \tag{18}$$

A good fit to the solar energy generation rate is obtained with $\nu = 4$; then Eqns.(17) and (18) imply a bound on the axion rate

$$\delta_{ax}(T_{c,\odot}) \leq 0.1 \tag{19}$$

The usually quoted bound on the axion rate arises from the requirement that the axion luminosity is less than the observed solar luminosity, $L_{ax} \leq L_\odot$. This constraint on L_{ax} has no physical basis, however, since L_{ax} is not an observable quantity. More to the point, our bound (19) is an order of magnitude stronger than this. In addition, Eqns. (15) and (19) yield a relatively strong constraint on the allowed change in solar luminosity,

$$\left(\frac{\delta L}{L}\right)_\odot \leq 0.01 \tag{20}$$

From Eqn.(20), it is clear that the solar luminosity is not a useful observable for placing limits on the axions.

The coupling of Nambu-Goldstone bosons to electrons is

$$\mathcal{L} = g_{Be} \bar{e} i \gamma_5 e \varphi \ .$$

From bremsstrahlung and Compton emission[3,4] of Nambu-Goldstone bosons, our self-consistent limit, $\delta_{ax} \leq 0.1$, corresponds to a solar bound on the pseudoscalar coupling constant,

$$g_{Be} < 1.6 \times 10^{-11} \ . \tag{21}$$

For axion models, $g_{ae} = (2X_e'/F)\, m_e$, which yields a lower bound on the Peccei-Quinn scale,

$$\frac{F}{2X_e'} > 3.2 \times 10^7 GeV \ , \tag{22}$$

a factor of 3 stronger than the usually quoted limit. (Here, $2X_e'$ is a model-dependent parameter of order unity.[12]) Assuming the axion mass is[1,12] $m_{ax} \approx 7.2\ eV\ (10^7 GeV/F)$, this corresponds to the upper limit $m_{ax} \lesssim 2.3\, eV$ (for $2X_e' = 1$). For the majoron model of Gelmini and Roncadelli[1] , the coupling constant is $g_{Me} = (1.6 \times 10^{-11} MeV^{-1})\, V_T$, where V_T is the triplet majoron vacuum expectation value. The bound (21) then corresponds to the limit

$$V_T \leq 1\ MeV \ . \tag{23}$$

If the coupling constant is near the upper bound of Eqn.(21), Nambu-Goldstone boson emission from the sun will have additional consequences. For example, the solar neutrino luminosity per unit mass from ^8B decay is approximately

$\varepsilon_\nu = A\rho T^c$, with a steep temperature dependence, $c \simeq 13$.[8,11] By analogy to Eqn.(13), the perturbed neutrino luminosity is $L_{1,\nu}(x) = y^{-(3+c)} L_{o,\nu}(x)$, and from Eqn.(14) we find

$$\frac{\delta L_\nu}{L_\nu} = \frac{(c+3)\delta_{ax}}{(\nu+2.5)} \leq 0.25 \tag{24}$$

Thus, axion emission at the solar bound would exacerbate the solar neutrino problem. Until the solar neutrino problem is understood, however, we cannot use (24) to obtain a better axion limit.

Saturation of the bound (18) would also reduce the main sequence lifetime of solar mass stars by of order 10%. The age of the oldest globular clusters corresponds roughly to the main sequence lifetime of solar mass stars; that is, from isochrone fitting of cluster Hertzsprung-Russell diagrams, the main-sequence turn-off point for old globular clusters is found to occur at a mass near $M \simeq M_\odot$. Since axion emission reduces stellar lifetimes, isochrone fitting using stellar evolutionary models without axions taken into account will overestimate the ages of globular clusters. This is good news for advocates of inflation, since some estimates for globular cluster ages, of order $(16-18) \times 10^9$ years, in conjunction with the inflationary prediction $\Omega = 1$ (assuming $\Lambda = 0$), imply an uncomfortably low value of the Hubble parameter, $H_o \simeq 35 - 40$ $km/sec/Mpc$ (the observed range for the Hubble constant is generally taken to be 75 ± 25 $km/sec/Mpc$).

It has recently been recognized that axions emitted by the sun may be detected in deep underground detectors by the axioelectric effect.[13] The lack of an observed signal at present detector sensitivities yields a bound on the axion scale,[14]

$$\frac{F}{2X_e'} > 10^7 \; GeV \qquad \text{(experiment)} \tag{25}$$

Comparison with Eqn. (22) shows that this limit, like the usually quoted solar bound, is not self-consistent: Eqn. (25) was derived using a standard solar model, but axion emission saturating this bound could severely perturb the evolution of the sun ($\delta_{ax} \sim 1$). A proper treatment would have to include the nonlinear effects of axions on solar evolution, and the consequent changes in the solar axion flux. This would alter the limit (25) slightly, but clearly cannot raise it above the solar bound (22). As detectors are improved, however, the experimental bound will become competitive with the solar limit.

III. LOW MASS STARS

Given the approximations embodied in Eqn.(11), the homology treatment given above should provide an adequate description of the response of solar-type stars to axion perturbations. These stars have thin convective envelopes, and their structure is relatively insensitive to surface boundary conditions; to first approximation, they can be treated as being in radiative equilibrium throughout. This approach is further justified by the fact that the boundary of the convective envelope is a homologous point.

For lower mass stars, the situation is different. Low mass stars have sufficiently low surface temperatures for neutral hydrogen to form in abundance in the envelope. As a result, molecular contributions to the opacity become important, and κ *increases* sharply with temperature (instead of decreasing, as in the Kramers law for fully ionized atoms). In this case, near the surface the opacity increases rapidly with depth, and radiative transfer becomes unstable to the onset of convection. This instability is responsible for the convective envelope of the sun as well; however, for the cooler low-mass stars, the convective envelope is deep and plays an important role in determining internal structure. We will consider the transition to deep convection to occur at a stellar mass $M_{tr} \simeq (0.3-0.4)M_{\odot}$, i.e., for the late M dwarfs.

For simplicity, we will treat stars with mass $M \leq M_{tr}$ as completely convective[♯2] (computer models indicate full convection for $M \lesssim 0.3M_{\odot}$[15]). The difficulty is that the luminosity of such stars is sensitive to surface boundary conditions. For a star in which most of the energy flux is transported by bulk convective motion of the fluid, the equation of radiative transfer (4) must be replaced by the mixing length theory of convection, which gives

$$L_{conv} \sim (\triangle \bigtriangledown T)^{3/2} \tag{26}$$

where

$$\triangle \bigtriangledown T = \left(\frac{dT}{dr}\right)_{ad} - \left(\frac{dT}{dr}\right) \tag{27}$$

is the superadiabatic excess, i.e., the excess of the temperature gradient over the adiabatic gradient. In contrast to the radiative luminosity (Eqn.(4)), L_{conv} is not determined by the temperature gradient itself but by the small excess (typically, $(\triangle \bigtriangledown T)/|dT/dr| \simeq 10^{-6}$). As a result, the temperature profile of the star is essentially decoupled from the convective energy transfer; this makes Eqn.(25) impractical for use in calculating the luminosity.

To compute the luminosity of completely convective stars, we must instead impose boundary conditions at the photosphere, the layer from which the stellar

♯2 The structure of low mass stars is further complicated by different contributions to the opacity in different regions (which we model crudely), as well as by the onset of partial electron degeneracy (and hence deviations from the ideal gas law) at high density (very low mass stars). As a result, the homology assumption for such stars may be less justified, but it should still give a first approximation to their response to perturbations.

energy is radiated into space. By definition, the region above the photosphere must be in radiative equilibrium, since energy is radiated, not convected, into space. In practical terms, we can define the photosphere as the region in which the stellar temperature is approximately equal to the effective temperature, $T_{ph} \simeq T_e$, which is defined as the equivalent black body temperature for a star with given luminosity,

$$L = 4\pi R^2 \sigma T_e^4 \tag{28}$$

Generally, T_e is within a factor unity of the surface temperature. From our definition above, the depth of the photosphere is roughly the photon mean free path, $\Delta r_{ph} \simeq (\kappa_{ph}\rho)^{-1}$.

We can solve for the pressure at the photosphere by using the condition of hydrostatic equilibrium (2), evaluated at the stellar surface,

$$\frac{dp}{dr} \simeq -\frac{GM\rho}{R^2} \tag{29}$$

Thus

$$p_{ph} \simeq \frac{GM\rho(\Delta r_{ph})}{R^2} \simeq \frac{GM}{\kappa_{ph} R^2} \tag{30}$$

where we have neglected radiation pressure at the surface. We have also assumed the photosphere is sufficiently thin to take the opacity there as roughly constant; in low mass stars, this is a reasonable approximation. (We could use a more sophisticated model of the photosphere than the simple thin atmosphere treatment given here, but the results do not differ substantially.) Equation (29) furnishes the requisite boundary condition to find the luminosity. In the photosphere, the opacity can again be taken to have the homogeneous form

$$\kappa_{ph} = \kappa_o \rho_{ph}^n T_e^{-s} \tag{31}$$

where now the exponent $s < 0$ reflects the influence of partial ionization. For numerical estimates, we will take $n = 1$, $s = -4$, appropriate to envelopes of low-mass stars with $T_e < 5 \times 10^3 K$.[10] Assuming the ideal gas approximation, Eqn.(29) becomes

$$p_{ph} \sim \frac{M^{1/2}}{RT_e^{3/2}} \tag{32}$$

In the convective interior, bulk motion of the fluid establishes nearly adiabatic conditions. As a result, the interior is well described by the polytropic equation of state with polytropic index $n = 1.5$,

$$p = K(M, R)T^{2.5} \tag{33}$$

where $K \sim M^{-1/2}R^{-3/2}$ is a constant for the star. (Note that (32) is invariant under a homology transformation (6-9)). In our approximation, the outer boundary of the convective zone is the base of the photosphere; imposing the boundary condition (31) on the pressure (32), we find the effective temperature scales as

$$T_e^4 \sim MR^{1/2} \tag{34}$$

and from (27) the luminosity as

$$L_{conv} \sim MR^{5/2} \tag{35}$$

Under a homology transformation, this gives

$$L_{conv} \sim y^{5/2} L_o(x) \tag{36}$$

On the other hand, the condition of thermal equilibrium still implies Eqn.(13). Comparing Eqns.(13) and (35) and linearizing as before, we find the resultant changes in the structure of convective stars:

$$-\frac{\delta R}{R} = \frac{\delta T_c}{T_c} = \frac{\delta_{ax}}{\nu + 5.5} \tag{37}$$

$$\frac{\delta L}{L} = \frac{20\delta T_e}{T_e} = \frac{2.5\delta\tau}{(\nu+3)\tau} = -\frac{2.5\delta_{ax}}{(\nu+5.5)} \,. \tag{38}$$

Two features of these relations are worth noting. First, the radius of the convective star is less sensitive to the axion perturbation than for the radiative star (this is compounded by the growth of the exponent ν at low temperature). Second, since $s < 0$, a convective star grows dim under the influence of axions, while a radiative star brightens(15). Unfortunately, even the axion-shortened lifetime of low mass stars is much longer than the age of the universe, so stellar ages cannot be used to constrain axion emission.

For our purposes, the best diagnostic to use for studying axion perturbations of low mass stars is the mass-luminosity relation. (Because axions do not affect the stellar mass, this test is preferable to the usual Hertzsprung-Russell diagram, i.e., the $L - T_e$ plane.) The results are shown in Fig.1, where we have used the main-sequence models of Refs.15; to calculate δ_{ax}, we have taken ε_{ax} from refs. 3-5 (including Compton emission) and have used the interpolation formulae for ε_{nuc} and ν from Clayton.[8] (This will result in a slight overestimate of ε_{nuc} for $M \lesssim 0.5 M_\odot$ due to the assumption of He^3 equilibration.[15])

In general, one expects axion effects to become more important at low stellar mass. However, in the mass range $(0.2-0.3)M_{\odot}$ bremsstrahlung emission, which is the dominant axion process down to $0.2M_{\odot}$, becomes severely damped by plasma screening[16] (see Fig. 3 of Ref. 5). Thus, δ_{ax} declines as the mass is lowered from $0.3M_{\odot}$ to $0.2M_{\odot}$.[17] In addition, since the central temperature only drops by a factor of 2 between $1M_{\odot}$ and $0.3M_{\odot}$, δ_{ax} does not grow appreciably in this range. For a pseudoscalar coupling at the solar bound (21), the axion effects remain linear down to the lowest stellar masses.

The luminosity changes produced by axions may also be partially compensated for by uncertainties in stellar input parameters and stellar physics, in particular in the surface (molecular) opacities and in the metallicity. For example, for a star of mass $0.3M_{\odot}$, we find that axions saturating the solar bound (21) reduce the luminosity by

$$\frac{\delta L}{L} \simeq \delta \log\left(\frac{L}{L_{\odot}}\right) = -0.04 \tag{39}$$

This is comparable to the change in luminosity induced by a factor 10 reduction in the He^3 abundance[15] or by a corresponding increase in the metallicity. The results of refs. 15 are in reasonable agreement with the observations, perhaps slightly subluminous if high metallicity (e.g., Pop. I, $Z = 0.02$) is assumed for most of the sample. Axions would worsen the agreement on the low mass end, but this could be accounted for by assuming a reduced metallicity for the low mass stars of the sample. Thus, it appears difficult to obtain a significantly stronger independent bound on the axion scale from low mass stars.

One possibility is to study the pre-main-sequence contraction of very low mass stars. Stars with mass less than $M_{cr} \simeq 0.08M_{\odot}$ are believed to never reach a stable hydrogen-burning main sequence phase.[18] This is because the interior temperature of a star contracting in the convective pre-main sequence phase reaches a maximum when the electron gas becomes degenerate. If the nuclear generation rate ϵ_{nuc} is too low to establish thermal equilibrium at $T_{\max}$, the star continues to contract. The degenerate configuration, however, now *cools* as it contracts, and the star moves further out of equilibrium (recall $\epsilon_{nuc} \sim T^{\nu}$). Stars with mass $M < M_{crit}$ end up as degenerate brown dwarfs.

Energy loss due to axion emission reduces the net heat output for given interior conditions. As a result, stars above the critical mass will have to contract to higher densities to achieve thermal equilibrium; some of these will become degenerate and never reach equilibrium. The effect of axions is thus to raise the main sequence cut-off mass M_{crit}. A simple analysis using the results of refs. 5 and 15 (here, only the axiorecombination process is important) indicates the critical mass is raised to $M'_{crit} \simeq 0.15 M_{\odot}$, again assuming the solar bound (21) is saturated. A number of stars have been observed in the mass range $M_{crit} \leq M \leq M'_{crit}$. If any of these can be demonstrated to be on the main sequence, this would yield an independent bound on the axion scale, slightly stronger than the solar limit. This avenue appears particularly promising since recent theoretical determinations of M_{crit} appear to be relatively insensitive to uncertainties in opacities and metal abundances.

IV. CONCLUSION

We have discussed the effect of Nambu-Goldstone boson emission on the structure and evolution of main sequence stars. Our primary aim in so doing was to clarify previous work on the subject; although the usually quoted solar constraint on axion emission, $L_{ax} \leq L_{\odot}$, is suggestive, it is not empirically verifiable, for L_{ax} is not an observable. As a result of our more careful treatment, we obtained a factor 3 improvement on the upper limit to the coupling of light pseudoscalars to electrons and have obtained a bound which is self-consistent. There are stronger, but relatively more speculative, constraints on axions which arise by considering axion emission from red giants,[3,7] white dwarfs[20], and neutron stars[21]; the strongest of these corresponds to $g_{Be} \lesssim 2 \times 10^{-9}$. The solar limit, though less restrictive, is, however, based on a better understood, nearby star.

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17. Actually, for a central density and temperature corresponding to a stellar mass slightly above 0.2 $M_{\odot}$, the effect of plasma damping is to make

the effective temperature dependence of the bremsstrahlung energy loss rate steeper than the nuclear generation rate. As a result, the star must contract to higher temperature and density (and reduced bremsstrahlung temperature dependence) to restore equilibrium. This increased relaxation enhances the value of δ_{ax} in this mass range.

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FIGURE CAPTIONS

1. Mass-luminosity relation for low mass stars. Filled circles represent data compiled by Popper.[19] Solid lines are theoretical results of Vandenberg, et al.[15] and open circles are theoretical points of Grossman, et al.[15] The dashed line indicates the modified M-L relation including the effects of axion emission if the solar bound (Eqn. 20) is saturated; here, the unperturbed models were taken to be those of Vandenberg, et al., with $Z = 0.02$.

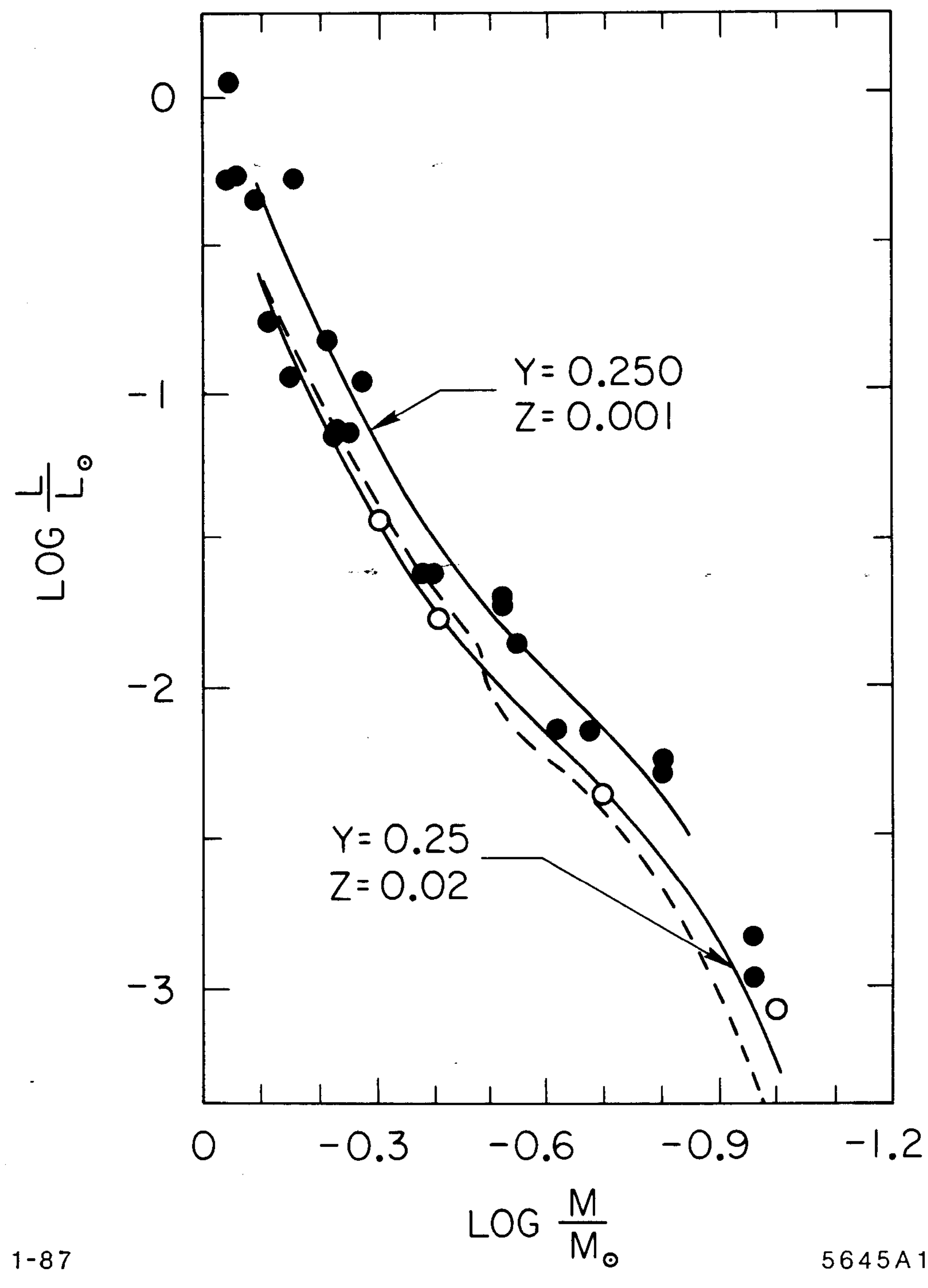
0
-1
-2
-3
LOG L/L⊙
Y= 0.250
Z= 0.001
Y= 0.25
Z= 0.02
0
-0.3
-0.6
-0.9
-1.2
LOG M/M⊙
1-87
5645A1

Fig. 1