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Astrophysics probes of particle physics

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Abstract

Astrophysical arguments constrain the properties of various elementary particles in ways which are often complementary to cosmological arguments and to laboratory experiments. This pertains, in particular, to neutrinos, axions, other Nambu–Goldstone bosons, and gravitons, which are light so that their production in the hot and dense interior of stars is not impeded by threshold effects. This review provides an update to the most important stellar-evolution limits and discusses them in the context of other information from cosmology and laboratory experiments. © 2000 Elsevier Science B.V. All rights reserved.

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1. Introduction

Astrophysical and cosmological arguments and observations have become part of the mainstream methodology to obtain empirical information on existing or hypothetical elementary particles and their interactions. Conversely, novel particle-physics hypotheses are invoked to explain several long-standing astrophysical and cosmological mysteries, notably the nature of dark matter, the solar and atmospheric neutrino anomalies, and the nature and origin of the highest-energy cosmic rays, topics which are covered by other contributions to this volume. Likewise, the role of big-bang nucleosynthesis as a cosmological particle laboratory is explored by other authors in this volume. Therefore, it is natural for my review to focus on the role of stars as particle-physics laboratories, one of David Schramm's many research interests.

The prime argument to be exploited is that a hot and dense stellar plasma is a powerful source for low-mass weakly interacting particles, notably neutrinos, axions, and gravitons. These particles subsequently escape from the star, without further interactions, and thus provide a local energy

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sink for the stellar medium. The astronomically observable impact of this phenomenon provides some of the most powerful limits on the properties low-mass weakly interacting particles. Once they have escaped from the star, they can decay on their long way to Earth, allowing one to derive interesting limits on radiative decay channels from the absence of unexpected X- or γ-ray fluxes from the Sun or other stars. Finally, the weakly interacting particles can be directly detected at Earth, thus far only the neutrinos from the Sun and supernova (SN) 1987A, allowing one to extract important information on their properties.

Space constraints or a lack of expertise prevent me from discussing a number of other interesting ways in which stars can be used to probe particle physics. The solar neutrino problem and its oscillation interpretation is a topic unto itself and has been extensively reviewed [1–3]. The high densities encountered in neutron stars make them ideal for studies and speculations concerning novel phases of nuclear matter (e.g. meson condensates or quark matter), an area covered by two recent books [4,5]. Quark stars are also the subject of an older review [6] and are covered in the proceedings of two topical conferences [7,8]. Certain grand unified theories predict the existence of primordial magnetic monopoles. They would get trapped in stars and then catalyze the decay of nucleons by the Rubakov–Callan effect. The ensuing anomalous energy release is constrained by the properties of stars, in particular neutron stars and white dwarfs [9]. These limits were improved in the wake of the discovery of the faintest white dwarf ever detected [10]. Finally, weakly interacting massive particles (WIMPs) are prime candidates for the cosmic dark matter. Some of them would get trapped in stars, annihilate with each other, and produce a secondary flux of high-energy neutrinos. The search for such fluxes from the Sun and the center of the Earth by neutrino telescopes is the "indirect method" to detect galactic particle dark matter [11].

With David Schramm as an editor, stars as probes for particle physics have been reviewed in 1990 in Physics Reports by Turner [12] and by Raffelt [13], focusing on axion limits. Solicited by David Schramm, I cast this material into a book *Stars as Laboratories for Fundamental Physics* (1996) [15] which expanded and updated the previous works. A very brief "Mini-Review" was included in the 1998 edition of the Review of Particle Physics [14]. Finally, I will heavily draw on an updated review which I have recently prepared [16]. Still, it is more than justified to represent "Particle Physics from Stars" in this memorial volume because David Schramm took a keen interest in this topic, both by active research and by encouraging others.

This review begins, in Sections 2–4, with a discussion of the main stellar objects that have been used to constrain low-mass particles, viz. the Sun, globular-cluster stars, compact stars, and SN 1987A. As an application, the main constraints on neutrinos and axions are summarized in Sections 5 and 6, while Section 7 is given over to brief concluding remarks.

2. The Sun

2.1. Basic energy-loss argument

The Sun is the best-known star and thus a natural starting point for our survey of astrophysical particle laboratories. It is powered by hydrogen burning which amounts to the net reaction $4p + 2e^- \rightarrow {}^4\mathrm{He} + 2\nu_e + 26.73\,\mathrm{MeV}$, giving rise to the measured solar ν_e flux. However, instead of nuclear processes we focus on particle fluxes which are produced in thermal plasma reactions. The

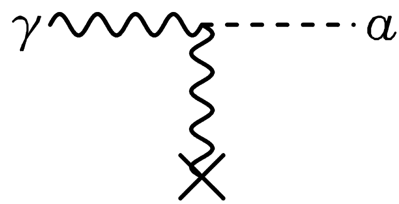

Fig. 1. Primakoff production of axions in the Sun.

solar energy loss from thermal neutrino or graviton emission is small, but it may be large for new particles. As an example we consider axions (Section 6); they can be produced by the Primakoff process in which thermal photons mutate into axions in the electric field of the medium's charged particles (Fig. 1). What is the backreaction of this new energy loss on the properties of the Sun?

The Sun is a normal star which supports itself against gravity by thermal pressure, as opposed to degenerate stars like white dwarfs which are supported by electron degeneracy pressure. According to the virial theorem, a normal star has a "negative heat capactiy" where an energy loss leads to contraction and heating. The nuclear energy generation rate scales with a high power of the temperature. Therefore, the heating implied by the new energy loss causes increased nuclear burning – the star finds a new equilibrium configuration where the new losses are compensated by an increased rate of energy generation. The main lesson is that the new energy loss does not "cool" the star; it leads to heating and an increased consumption of nuclear fuel. The Sun, where energy is transported from the central nuclear furnace by radiation, actually overcompensates the losses and brightens. This behavior is understood by a powerful "homology argument" where the nonlinear interplay of the equations of stellar structure is represented in a simple analytic fashion [17].

The solar luminosity is well measured, yet this brightening effect is not observable. The present-day luminosity of the Sun depends on its unknown initial helium mass fraction Y; in a solar model Y has to be adjusted such that $L_\odot = 3.85\times10^{33}\,\mathrm{erg\,s^{-1}}$ is reproduced after 4.6×10^9 years of nuclear burning. Even axion losses as large as $L_\odot$ can be accommodated by reducing the presolar helium mass fraction from about 27% to something like 23% [18,29]. The "standard Sun" has completed about half of its hydrogen-burning phase. Therefore, the anomalous energy losses cannot exceed approximately $L_\odot$ or else the Sun could not have reached its observed age.

2.2. Solar neutrino measurements and helioseismology

This crude limit is improved by the solar neutrino flux which has been measured in five different observatories with three different spectral response characteristics. The axionic solar models are hotter and thus produce larger neutrino fluxes. For axion losses below a few tenths of $L_\odot$, one can still find oscillation solutions to the observed ν_e deficit, but larger energy-loss rates appear to be excluded [18]. Once the neutrino oscillation hypothesis has been more firmly established and the mixing parameters are better known, the neutrino measurements may be used to pin down the central solar temperature, allowing one to constrain novel energy losses with greater precision.

The recent precision measurements of the solar p-mode frequencies have provided a more reliable way to study the solar interior. For example, the helium content of the convective surface layers is found to exceed 0.238 [19]. Gravitational settling reduces the surface helium abundance by about 0.03 so that the presolar value must have been at least 0.268, in good agreement with standard solar models. The reduced helium content required of the axionic models disagrees with this lower limit if the axion luminosity exceeds about $0.2L_{\odot}$.

One may invert the p-mode measurements to construct a "seismic model" of the solar sound-speed profile. All modern standard solar models agree well with the seismic model within its uncertainties. When the axion luminosity exceeds 10–20% of $L_{\odot}$, the difference exceeds the uncertainties of the seismic model, implying a limit on the axion–photon coupling constant of [18]

$$g_{a\gamma} \lesssim 1.0 \times 10^{-9}\,\mathrm{GeV}^{-1}\ . \tag{1}$$

Other cases may be different in detail, but it is safe to assume that any new energy-loss channel must not exceed something like $0.1L_{\odot}$.

3. Limits on stellar energy losses

3.1. Globular-cluster stars

The previous discussion suggests that the emission of new weakly interacting particles from stars primarily modifies the time scale of evolution. For the Sun, this effect is less useful to constrain particle emission than the modified p-mode frequencies or the direct measurement of the neutrino fluxes. However, the observed properties of other stars provide restrictive limits on certain evolutionary time scales so that anomalous modes of energy loss can be tightly constrained. We begin with globular-cluster stars which, together with SN 1987A, are the most successful example of astronomical observations that provide nontrivial limits on the properties of elementary particles.

Our galaxy has about 150 globular clusters which are gravitationally bound systems of up to a million stars. The stars in a cluster all formed at the same time with essentially the same chemical composition, differing primarily in their mass. Globular clusters are nearly as old as the universe, implying that stars more massive than about $1\mathscr{M}_{\odot}$ (solar mass) have already completed their evolution. For most of their lives, these low-mass stars burn hydrogen at their center. When central hydrogen is exhausted, they develop a degenerate helium core, with hydrogen burning in a shell. The envelope expands, leading to a large surface area and thus a low surface temperature – they become "red giants". The luminosity is governed by the gravitational potential at the edge of the growing helium core so that these stars become ever brighter: they ascend the red-giant branch (RGB) in the color–magnitude diagram. The higher a star on the RGB, the more massive its helium core, which grows to about $0.5\mathscr{M}_{\odot}$ when it ignites helium. The ensuing core expansion reduces the gravitational potential and thus lowers the energy production rate in the hydrogen shell source. After helium ignition, these stars occupy the horizontal branch (HB) at a much lower total luminosity than they had at the tip of the RGB. Finally, when helium is exhausted, a degenerate carbon–oxygen core develops, leading to an ascent on the asymptotic giant branch (AGB).

Anomalous energy losses modify this picture in measurable ways. We first consider an energy-loss mechanism which is more effective in the degenerate core of a red giant before helium ignition than on the HB so that the post-RGB evolution is standard. Since an RGB-star's helium core is supported by degeneracy pressure there is no feedback between energy loss and pressure: the core is

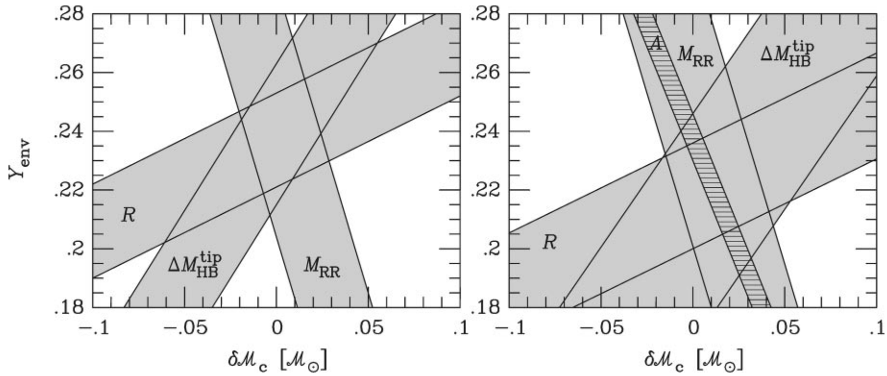

Fig. 2. Allowed values for a core-mass excess at helium ignition $\delta\mathcal{M}_c$ and the envelope helium mass fraction Y_{env} of evolved globular-cluster stars. Left panel after [15], right panel after [22]. The observables are the brightness difference ΔM_{HB}^{tip} between the HB and the RGB tip, the RR Lyrae mass-to-light ratio A, their absolute brightness M_{RR}, and the number ratio R between HB and RGB stars.

actually *cooled*. This delays the ignition of helium, leading to a larger core mass $\mathcal{M}_c$, with several observable consequences. First, the brightness of a red giant depends on its core mass so that the RGB would extend to larger luminosities, causing an increased brightness difference ΔM_{HB}^{tip} between the HB and the RGB tip. Second, an increased $\mathcal{M}_c$ implies an increased helium-burning core on the HB. For a certain range of colors these stars are pulsationally unstable and are then called RR Lyrae stars. Their measured luminosity and pulsation period implies $\mathcal{M}_c$ on the basis of their "mass-to-light ratio" A. Third, the increased $\mathcal{M}_c$ increases the luminosity of RR Lyrae stars so that absolute determinations of their brightness M_{RR} allow one to constrain the range of possible core masses. Fourth, the number ratio R of HB stars vs. RGB stars brighter than the HB is modified.

These observables depend on the measured cluster metallicity as well as the unknown helium content, which is usually expressed in terms of Y_{env}, the envelope helium mass fraction. The initial globular-cluster helium content must be close to the primordial value of 22–25%. Y_{env} should be close to this number because the initial value is somewhat depleted by gravitational settling, and somewhat increased by convective dredge-up of helium-rich material. An estimate of $\mathcal{M}_c$ from a global analysis of these observables except A was performed in [20] and re-analysed in [15], A was used in [21], and an independent analysis using all four observables in [22]. In Fig. 2 we show the allowed core mass excess $\delta\mathcal{M}_c$ and envelope helium mass fraction Y_{env} from the analyses [15,22].

Fig. 2 suggests that, within the given uncertainties, the observations overlap at the standard core mass ($\delta\mathcal{M}_c = 0$) and at a value for Y_{env} which is compatible with the primordial helium abundance. Of course, the error bands do not have a simple interpretation because they combine observational and estimated systematic errors, which involve some subjective judgement by the authors. The difference between the two panels of Fig. 2 gives one a sense of how sensitive the conclusions are to these more arbitrary aspects of the analysis. As a nominal limit it appears safe to adopt $|\delta\mathcal{M}_c| \lesssim 0.025$ or $|\delta\mathcal{M}_c|/\mathcal{M}_c \lesssim 5\%$.

In [15] it was shown that this limit can be translated into an approximate limit on the average anomalous energy-loss rate ε_x of a helium plasma,

$$\varepsilon_x \lesssim 10\,\mathrm{erg\,g^{-1}\,s^{-1}} \quad \text{at } T \approx 10^8\,\mathrm{K}, \quad \rho \approx 2\times 10^5\,\mathrm{g\,cm^{-3}}\,. \tag{2}$$

The density represents the approximate average of a red-giant core before helium ignition; the value at its center is about $10^6\,\mathrm{g\,cm^{-3}}$. The main standard-model neutrino emission process is plasmon decay $\gamma \to \nu\bar{\nu}$ with a core average of about $4\,\mathrm{erg\,g^{-1}\,s^{-1}}$. Therefore, Eq. (2) means that a new energy-loss channel must be less effective than a few times the standard neutrino losses.

We now turn to an energy-loss mechanism which becomes effective in a nondegenerate medium, i.e. the core expansion after helium ignition "switches on" an energy-loss channel that was negligible on the RGB. As for the Sun (Section 2.1), there will be little change in the HB stars' brightness, rather they will consume their nuclear fuel faster and thus begin to ascend the AGB sooner. The net observable effect is a reduction of the number of HB relative to RGB stars. From the measured HB/RGB number ratios in 15 globular clusters [23] one concludes that the duration of helium burning agrees with stellar-evolution theory to within about 10% [5]. Thus, the new energy loss of the helium core should not exceed about 10% of its standard energy production rate, implying a constraint at average core conditions of [15]

$$\varepsilon_x \lesssim 10\,\mathrm{erg\,g^{-1}\,s^{-1}} \quad \text{at } T \approx 0.7\times 10^8\,\mathrm{K}, \quad \rho \approx 0.6\times 10^4\,\mathrm{g\,cm^{-3}}\,. \tag{3}$$

This limit is slightly more restrictive than the often-quoted "red-giant bound", corresponding to $\varepsilon_x \lesssim 100\,\mathrm{erg\,g^{-1}\,s^{-1}}$ at $T = 10^8\,\mathrm{K}$ and $\rho = 10^4\,\mathrm{g\,cm^{-3}}$. It was based on the helium-burning lifetime of the "clump giants" in open clusters [24]. They have fewer stars, leading to statistically less significant limits. "Clump giants" are the open-cluster equivalent of HB stars.

Evolutionary sequences including new energy losses have been calculated by several authors. Comparing the results from such studies with what one finds from Eqs. (2) and (3) reveals that, in view of the overall theoretical and observational uncertainties, it is indeed enough to use these simple criteria [15]. They can then be applied almost mechanically to a variety of cases. The main task is to identify the dominant emission process and to calculate the energy-loss rate ε_x for a helium plasma at the conditions specified in Eqs. (2) or (3). The most important limits will be discussed in the context of specific particle-physics hypotheses in Sections 5 and 6. Here we just mention that these and similar arguments were used to constrain neutrino electromagnetic properties [20,21,24,25], axions [26–31], paraphotons [32], the photo production cross section on ^{4}He of new bosons [33], the Yukawa couplings of new bosons to baryons or electrons [34,35], and supersymmetric particles [36–38].

3.2. *White dwarfs*

White dwarfs are another case where astronomical observations provide useful limits on new stellar energy losses. These compact objects are the remnants of stars with initial masses of up to several $\mathscr{M}_\odot$. When they ascend the asymptotic giant branch they shed most of their envelope mass. The degenerate carbon–oxygen core, having reached something like $0.6\mathscr{M}_\odot$, never ignites; it subsequently simply cools. The cooling speed is infered from the white-dwarf number density per brightness interval, i.e. the "luminosity function". Its sharp drop at the faint end indicates how far the oldest white dwarfs have cooled, implying that they were born 8–12 Gyr ago, in good agreement with the estimated age of the galaxy. Therefore, a novel cooling agent cannot be much more effective than the surface photon emission. The shape of the luminosity function can also be used as an observable because it would be deformed for an appropriate temperature dependence of the particle emission rate.

White dwarfs were used to constrain the axion-electron coupling [39–41]. It was also noted that the somewhat large period decrease of the ZZ Ceti star G117-B15A, a pulsationally unstable white dwarf, could be ascribed to axion cooling [42]. Moreover, a limit on the neutrino magnetic dipole moment was derived [41]. A detailed review of these limits is provided in [15]; they are somewhat weaker than those from globular-cluster stars.

3.3. Old neutron stars

Neutron stars are the compact remnants of stars with initial masses beyond about $8\mathcal{M}_\odot$. After their formation in a core-collapse supernova (Section 4) they evolve by cooling, a process that speeds up by a new energy-loss channel. Neutron-star cooling can now be observed by satellite-borne X-ray measurements of the thermal surface emission of several old pulsars [43]. Limits on axions were derived in [44,45], on neutrino magnetic dipole moments in [46]. These bounds are much weaker than those from SN 1987A or globular clusters. Anomalous cooling by particle emission is probably not important in old neutron stars, leaving them as laboratories for other uncertain bits of input physics such as the existence of new phases of nuclear matter [4,5,43].

4. Supernovae

4.1. SN 1987A neutrino observations

When the explosion of the star Sanduleak $-69\,202$ was detected on 23 February 1987 in the Large Magellanic Cloud, a satellite galaxy of our Milky Way at a distance of about 50 kpc (165,000 lyr), it became possible for the first time to measure the neutrino emission from a nascent neutron star, turning this supernova (SN 1987A) into one of the most important stellar particle-physics laboratories [47–49]. A type II supernova explosion [50–55] is physically the implosion of an evolved massive star ($\mathcal{M} \gtrsim 8\mathcal{M}_\odot$). Its degenerate iron core becomes unstable when it has reached its Chandrasekhar limit of 1–$2\mathcal{M}_\odot$. The ensuing collapse is intercepted when the equation of state stiffens at around nuclear density ($3\times10^{14}\,\mathrm{g\,cm^{-3}}$), corresponding to a core size of a few tens of kilometers. At temperatures of tens of MeV this compact object is opaque to neutrinos. The gravitational binding energy of the newborn neutron star ("proto neutron star") of about 3×10^{53} erg is thus radiated over several seconds from the "neutrino sphere". Crudely put, the collapsed SN core cools by thermal neutrino emission in all flavors from its surface. The neutrino signal from SN 1987A was observed by the $\bar{\nu}_e p \to n e^+$ reaction in several detectors [49]. The number of events, their energies, and the distribution over several seconds corresponds well to theoretical expectations. Detailed statistical analyses of the data were performed in [56–58].

4.2. Signal dispersion

A dispersion of the neutrino burst can be caused by a time-of-flight delay from a nonvanishing neutrino mass [59]. The arrival time from SN 1987A at a distance D would be delayed by

$$\Delta t = 2.57\,\mathrm{s}\left(\frac{D}{50\,\mathrm{kpc}}\right)\left(\frac{10\,\mathrm{MeV}}{E_\nu}\right)^2\left(\frac{m_\nu}{10\,\mathrm{eV}}\right)^2 . \tag{4}$$

As the $\bar\nu_e$ were registered within a few seconds and had energies in the 10 MeV range, m_{ν_e} is limited to less than around 10 eV. Detailed analyses reveal that the pulse duration is consistently explained by the intrinsic SN cooling time and that $m_{\nu_e} \lesssim 20\,\mathrm{eV}$ is implied at something like a 95% CL limit [56,60].

The apparent absence of a time-of-flight dispersion effect of the $\bar\nu_e$ burst was also used to constrain a "millicharge" of these particles (they would be deflected in the galactic magnetic field) [1,61], a quantum field theory with a fundamental length scale [62], and deviations from the Lorentzian rule of adding velocities [63]. Limits on new long-range forces acting on the neutrinos seem to be invalidated in the most interesting case of a long-range leptonic force by screening from the cosmic background neutrinos [64].

The SN 1987A observations confirm that the visual SN explosion occurs several hours after the core-collapse and thus after the neutrino burst. Again, there is no apparent time-of-flight delay of the relative arrival times between the neutrino burst and the onset of the optical light curve, allowing one to confirm the equality of the relativistic limiting velocity for these particle types to within 2×10^{-9} [65,66]. Moreover, the Shapiro time delay in the gravitational field of the galaxy of neutrinos agrees with that of photons to within about 4×10^{-3} [67], constraining certain alternative theories of gravity [68,69].

4.3. Energy-loss argument

The late events in Kamiokande and IMB reveal that the signal duration was not anomalously short. Very weakly interacting particles would freely stream from the inner core, removing energy which otherwise would power the late-time neutrino signal. Therefore, its observed duration can be taken as evidence against such novel cooling effects. This argument has been advanced to constrain axion–nucleon couplings [70–77], majorons [78–82], supersymmetric particles [83–90], and graviton emission in quantum-gravity theories with higher dimensions [91,92]. It has also been used to constrain right-handed neutrinos interacting by a Dirac mass term [71,93–99], mixed with active neutrinos [100,101], interacting through right-handed currents [71,102–105], or a magnetic dipole moment [106,107]. Many of these results will be reviewed in Sections 5 and 6 in the context of specific particle-physics hypotheses.

Here we illustrate the general argument with axions (Section 6) which are produced by nucleon bremsstrahlung $NN \to NNa$ so that the energy-loss rate depends on the axion–nucleon Yukawa coupling g_{aN}. In Fig. 3 we show the neutrino-signal duration as a function of g_{aN}. With increasing g_{aN}, corresponding to an increasing energy-loss rate, the signal duration drops sharply. For a sufficiently large g_{aN}, however, axions no longer escape freely; they are trapped and thermally emitted from the "axion sphere". Beyond some coupling strength axions cannot be excluded.

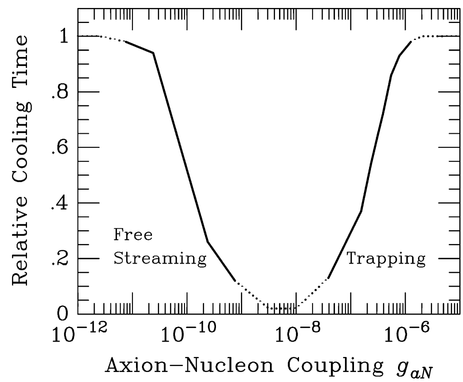

Fig. 3. Relative duration of SN neutrino cooling as a function of the axion–nucleon coupling. Freely streaming axions are emitted from the entire core volume, trapped ones from the "axion sphere". The solid line follows from the numerical calculations [74,75]; the dotted line is an arbitrary continuation.

However, particles which are on the "strong interaction" side of this argument need not be allowed. They could be important for the energy-transfer during the infall phase and they could produce events in the neutrino detectors. "Strongly coupled" axions in a large range of g_{aN} are actually excluded because they would have produced too many events by their absorption on ^{16}O [108].

Likewise, particles on the free-streaming side can cause excess events in the neutrino detectors. For example, right-handed neutrinos escaping from the inner core could become "visible" by decaying into left-handed states [109] or by spin-precessing in the galactic magnetic field if they have a dipole moment.

Returning to the general argument, one can estimate a limit on the energy-loss rate on the free-streaming side by the simple criterion that the new channel should be less effective than the standard neutrino losses, corresponding to [15]

$$\varepsilon_x \lesssim 10^{19}\,\mathrm{erg\,g^{-1}\,s^{-1}} \quad \text{at } \rho = 3\times 10^{14}\,\mathrm{g\,cm^{-3}}, \quad T = 30\,\mathrm{MeV}\,. \tag{5}$$

The density is the core average, the temperature an average during the first few seconds. Some authors find higher temperatures, but for a conservative limit it is preferable to stick to a value at the lower end of the plausible range. At these conditions the nucleons are partially degenerate while the electrons are highly degenerate. Several detailed numerical studies reveal that this simple criterion corresponds to approximately halving the neutrino signal duration [15].

4.4. Radiative neutrino decays

If neutrinos have masses one expects that the heavier ones are unstable and decay radiatively as $\nu \to \nu'\gamma$. SN 1987A is thought to have emitted similar fluxes of neutrinos and antineutrinos of all flavors so that one would have expected a burst of γ-rays in coincidence with the neutrinos. No excess counts were observed in the gamma-ray spectrometer (GRS) on the solar maximum mission (SMM) satellite [110,111], leading to restrictive limits on neutrino decays [110–114]. The GRS happened to go into calibration mode about 223 s after the neutrino burst, but for low-mass neutrinos ($m_\nu \lesssim 40\,\mathrm{eV}$) the entire γ-ray burst would have been captured, leading to a radiative decay limit of [15]

$$\tau_\gamma/m_\nu \gtrsim 0.8 \times 10^{15}\,\mathrm{s/eV}\ . \tag{6}$$

For higher-mass neutrinos, the photon burst would have been stretched beyond the GRS window. As a further complication, such higher-mass neutrinos violate the cosmological mass limit unless they decay sufficiently fast and thus nonradiatively. Comparable limits in the higher-mass range arise from γ-ray data of the Pioneer Venus Orbiter [115]. For $m_\nu \gtrsim 0.1\,\mathrm{MeV}$, decay photons still arrive years after SN 1987A. In 1991, the COMPTEL instrument aboard the Compton Gamma Ray Observatory looked at the SN 1987A remnant for about $0.68 \times 10^6\,\mathrm{s}$, providing the most restrictive limits in this mass range [116,117].

4.5. Explosion energetics

The standard scenario of a type II SN explosion has it that a shock wave forms near the edge of the core and that this shock wave ejects the mantle of the progenitor star. However, in typical numerical calculations the shock wave stalls so that this "prompt explosion" scenario does not seem to work. In the "delayed explosion" picture the shock wave is revived by neutrino heating, perhaps in conjunction with convection, but even then it appears difficult to obtain a successful or sufficiently energetic explosion. Therefore, one may speculate that nonstandard modes of energy transfer play an important role.

An example are Dirac neutrinos with a magnetic dipole moment of order $10^{-12}\,\mu_\mathrm{B}$ (Bohr magnetons). The right-handed (sterile) components would arise in the deep inner core by helicity-flipping collisions and escape. They precess back into interacting states in the large magnetic fields outside the SN core and heat the shock region; their interaction cross section would be relatively large because of their large inner-core energies [118–123].

Certainly it is important not to deposit *too much* energy in the mantle and envelope of the star. 99% of the gravitational binding energy of the neutron star goes into neutrinos, about 1% into the kinetic energy of the explosion, and about 0.01% into the optical supernova. Therefore, neutrinos or other particles emitted from the core must not decay radiatively within the progenitor's envelope radius of about 100 s, or else too much energy lights up [124,125].

4.6. Neutrino spectra and neutrino oscillations

Neutrino oscillations can have several interesting ramifications in SN physics because the temporal and spectral characteristics of the emission process depend on the neutrino flavor [52–54]. The simplest case is that of the "prompt ν_e burst" which represents the deleptonization of the outer core layers at about 100 ms after bounce when the shock wave breaks through the edge of the collapsed iron core. This "deleptonization burst" propagates through the mantle and envelope of the progenitor star so that resonant oscillations take place for a large range of mixing parameters between ν_e and some other flavor, notably for most of those values where the MSW effect operates in the Sun [126–135]. In a water Cherenkov detector this burst is visible as ν_e–e scattering, which is forward peaked. The first event in Kamiokande may be attributed to this signal, but this interpretation is statistically insignificant.

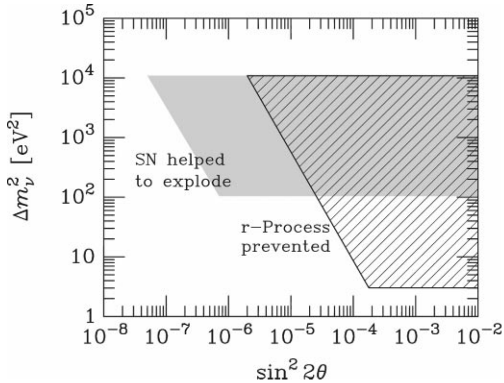

Fig. 4. Mass difference and mixing between ν_e and ν_μ or ν_τ where a spectral swap would occur to help explode supernovae, schematically after [136], and where it would prevent r-process nucleosynthesis, schematically after [137–139].

During the next few hundred milliseconds the shock wave stalls at a few hundred kilometers above the core and needs rejuvenating. The efficiency of neutrino heating is increased by resonant flavor oscillations which swap the ν_e flux with, say, the ν_τ one. Therefore, what passes through the shock wave as a ν_e was born as a ν_τ and has on average higher energies. In Fig. 4 the shaded range of mixing parameters is where supernovae are helped to explode, assuming a "normal" neutrino mass spectrum with $m_{\nu_e} < m_{\nu_\tau}$ [136].

The logic of this scenario depends on deviations from strictly thermal neutrino emission. The neutrino cross sections depend sensitively on energy and flavor so that the concept of a neutrino sphere is rather crude – the spectra are neither thermal nor equal for the different flavors [54,55]. The dominant opacity source for ν_e is the process $\nu_e + n \to p + e^-$, for $\bar\nu_e$ it is $\bar\nu_e + p \to n + e^+$, while for $\nu_{\mu,\tau}$ and $\bar\nu_{\mu,\tau}$ it is neutral-current scattering on nucleons. Therefore, unit optical depth is at the largest radius (lowest medium temperature) for ν_e, and deepest (highest temperature) for $\nu_{\mu,\tau}$ and $\bar\nu_{\mu,\tau}$. In typical calculations one finds $\langle E_{\nu_e}\rangle : \langle E_{\bar\nu_e}\rangle : \langle E_{\rm others}\rangle \approx \frac{2}{3} : 1 : \frac{5}{3}$ with $\langle E_{\bar\nu_e}\rangle =$ 14–17 MeV [53]. The SN 1987A observations imply a somewhat lower range of $\langle E_{\bar\nu_e}\rangle \approx$ 7–14 MeV [56–58]. It should be noted that, pending a more detailed numerical confirmation [140], the difference between the $\bar\nu_e$ and $\nu_{\mu,\tau}$ or $\bar\nu_{\mu,\tau}$ average energies appears to be smaller than commonly assumed [76,141,142].

A few seconds after core bounce the shock wave has long since taken off, leaving behind a relatively dilute "hot bubble" above the neutron-star surface. This region is one suspected site for r-process heavy-element synthesis, which requires a neutron-rich environment. The neutron-to-proton ratio, which is governed by β reactions, is shifted to a neutron-rich phase if $\langle E_{\nu_e}\rangle < \langle E_{\bar\nu_e}\rangle$. Resonant oscillations can again swap the ν_e flux with another one, inverting this hierarchy of energies. In the hatched range of mixing parameters shown in Fig. 4 the r-process would be disturbed [137–139]. On the other hand, $\nu_e \to \nu_s$ oscillations into a sterile neutrino could actually help the r-process by removing some of the neutron-stealing ν_e [143,144].

If the mixing angle between ν_e and some other flavor is large, the $\bar\nu_e$ flux from a SN contains a significant fraction of oscillated states that were born as $\bar\nu_\mu$ or $\bar\nu_\tau$ and thus should have higher average energies. The measured SN 1987A event energies are already somewhat low, so that a large-mixing-angle solution of the solar neutrino deficit poses a problem [58,60,145]. This conclusion, however, depends on the standard predictions for the average neutrino energies which may not hold up to closer scrutiny as mentioned above.

5. Limits on neutrino properties

5.1. Masses and mixing

After this survey of the most important stellar-evolution arguments we illustrate their use in the context of specific particle-physics cases. Beginning with neutrinos, the current discourse centers on the solar and atmospheric neutrino anomalies and the LSND experiment, which all provide suggestive evidence for neutrino oscillations. Solar neutrinos imply a Δm_ν^2 of about $10^{-5}\,\mathrm{eV}^2$ (MSW solutions) or $10^{-10}\,\mathrm{eV}^2$ (vacuum oscillations), atmospheric neutrinos 10^{-3}–$10^{-2}\,\mathrm{eV}^2$, and the LSND experiment 0.3–8 eV^2. Taken together, these results require a fourth flavor, a sterile neutrino, which is perhaps the most spectacular implication of these experiments, but also the least secure.

Core-collapse SNe are the one case in stellar astrophysics, apart from the Sun, where neutrino oscillations can be important. However, Fig. 4 reveals that the experimentally favored mass differences negate a role of neutrino oscillations for the explosion mechanism or r-process nucleosynthesis, except when sterile neutrinos exist [143,144]. Oscillations affect the interpretation of the SN 1987A signal [58,60,145] and that of a future galactic SN [152–154]. However, the main challenge at present is to develop a quantitatively more accurate understanding of neutrino spectra formation (Section 4.6).

Oscillation experiments reveal only neutrino mass differences, leaving the overall mass scale undetermined. The absence of anomalous SN 1987A signal dispersion (Section 4.2) gives a limit [56,60] $m_{\nu_e} \lesssim 20\,\mathrm{eV}$, somewhat weaker than current laboratory bounds. Observing a galactic SN with a detector like Superkamiokande could improve this limit to about 3 eV [146]. If the neutrino mass differences are as small as indicated by the current evidence for oscillations, this limit carries over to the other flavors. One can derive an independent mass limit on ν_μ and ν_τ in the range of a few 10 eV if one identifies a neutral-current signature in a water Cherenkov detector [147–149], or if a future neutral-current detector provides an additional measurement [150,151].

5.2. Dipole and transition moments

5.2.1. Plasmon decay in stars

Neutrino electromagnetic interactions would imply multifarious astrophysical consequences. The most interesting case are magnetic and electric dipole and transition moments. If the standard model is extended to include neutrino Dirac masses, the magnetic dipole moment is $\mu_\nu = 3.20 \times 10^{-19}\,\mu_\mathrm{B}\, m_\nu/\mathrm{eV}$ where $\mu_\mathrm{B} = e/2m_e$ is the Bohr magneton [155,156]. An electric dipole moment ε_ν violates CP, and both are forbidden for Majorana neutrinos. Flavor mixing implies electric and magnetic transition moments for both Dirac and Majorana neutrinos, but they are even smaller. Significant neutrino electromagnetic form factors require a more radical extension of the standard model, for example the existence of right-handed currents.

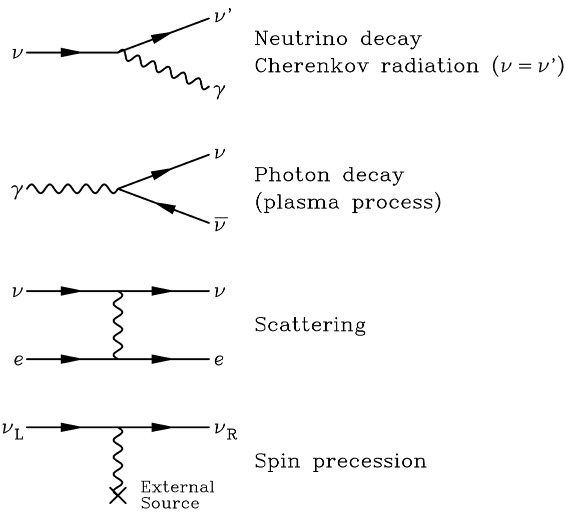

Fig. 5. Processes with neutrino electromagnetic dipole or transition moments.

Dipole or transition moments allow for many interesting processes (Fig. 5). For deriving limits, the most important case is $\gamma \to \nu\bar{\nu}$ which is kinematically possible in a plasma because the photon acquires a dispersion relation which roughly amounts to an effective mass. Even without anomalous couplings, the plasmon decay proceeds because the charged particles of the medium induce an effective neutrino–photon interaction [157,158]. The standard plasma process [159–161] dominates the neutrino production in white dwarfs or the cores of globular-cluster red giants.

The plasma process was first used in [162] to constrain neutrino electromagnetic couplings. The helium-ignition argument in globular clusters (Section 3.1), equivalent to Eq. (2), implies a limit [15,20,25]

$$\mu_\nu \lesssim 3 \times 10^{-12} \mu_{\mathrm{B}} \ , \tag{7}$$

applicable to magnetic and electric dipole and transition moments for Dirac and Majorana neutrinos. Of course, the final-state neutrinos must be lighter than the photon plasma mass, around 10 keV for the relevant conditions. The most restrictive laboratory bound is $\mu_{\nu_e} < 1.8 \times 10^{-10} \mu_{\mathrm{B}}$ at 90% CL from a measurement of the $\bar{\nu}_e$-e-scattering cross section [14]. A significant improvement should become possible with the MUNU experiment [163], but it is unlikely that the globular-cluster limit can be reached anytime soon.

5.2.2. Radiative decay

A neutrino mass eigenstate ν_i may decay to another one ν_j by the emission of a photon, where the only contributing form factors are the magnetic and electric transition moments. The inverse radiative lifetime is found to be [155,156]

$$\tau_\gamma^{-1} = \frac{|\mu_{ij}|^2 + |\varepsilon_{ij}|^2}{8\pi} \left(\frac{m_i^2 - m_j^2}{m_i} \right)^3 , \tag{8}$$

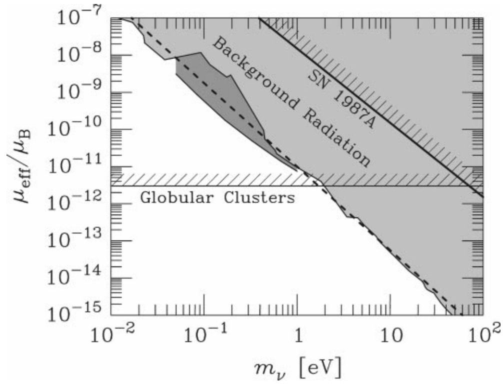

Fig. 6. Astrophysical limits on neutrino dipole moments. The light-shaded background-radiation limits are from [167], the dark-shaded ones from [168,169], the dashed line is the approximation formula in Eq. (9), bottom line.

where μ_{ij} and ε_{ij} are the transition moments while $|\mu_{\rm eff}|^2 \equiv |\mu_{ij}|^2 + |\varepsilon_{ij}|^2$. Radiative neutrino decays have been constrained from the absence of decay photons of reactor $\bar\nu_e$ fluxes [164], the solar ν_e flux [165,166], and the SN 1987A neutrino burst [110–114]. For $m_\nu \equiv m_i \gg m_j$ these limits can be expressed as

$$\frac{\mu_{\rm eff}}{\mu_{\rm B}} \lesssim \begin{cases} 0.9\times 10^{-1}\ ({\rm eV}/m_\nu)^2 & \text{Reactor } (\bar\nu_e)\ , \\ 0.5\times 10^{-5}\ ({\rm eV}/m_\nu)^2 & \text{Sun } (\nu_e)\ , \\ 1.5\times 10^{-8}\ ({\rm eV}/m_\nu)^2 & \text{SN 1987A (all flavors)}\ , \\ 1.0\times 10^{-11}\ ({\rm eV}/m_\nu)^{9/4} & \text{Cosmic background (all flavors)}\ . \end{cases} \tag{9}$$

The SN 1987A limit applies for $m_\nu \lesssim 40\,{\rm eV}$ as explained in Section 4.4. The decay of cosmic background neutrinos would contribute to the diffuse photon backgrounds, excluding the shaded areas in Fig. 6. They are approximately delineated by the dashed line, corresponding to the bottom line in Eq. (9). For low-mass neutrinos, the m_ν^3 phase-space factor in Eq. (8) is so punishing that the globular-cluster limit is the most restrictive one for m_ν below a few eV. This is precisely the mass range which today is favored from neutrino oscillation experiments.

5.2.3. Spin-flip scattering

The magnetic or electric dipole interaction couples neutrino fields of opposite chirality. In the relativistic limit this implies that a neutrino flips its helicity in an "electromagnetic collision", which in the Dirac case produces the sterile component. The active states are trapped in a SN core so that spin-flip collisions open an energy-loss channel in the form of sterile states. The SN 1987A energy-loss argument (Section 4.3) thus implies a limit $\mu_\nu({\rm Dirac}) \lesssim 3\times 10^{-12}\,\mu_{\rm B}$ for both electric and magnetic dipole and transition moments [106,107]. It is the same as Eq. (7), which however includes the Majorana case.

5.2.4. Spin and spin-flavor precession

Neutrinos with magnetic or electric dipole moments precess in external magnetic fields [171,172]. For example, solar neutrinos can precess into sterile and thus undetectable states in the

Sun's magnetic field [173,174]. The same for SN neutrinos in the galactic magnetic field where an important effect obtains for $\mu_\nu \gtrsim 10^{-12}\,\mu_B$. Moreover, the high-energy sterile states produced in the inner SN core could precess back into active ones and cause events with anomalously high energies in SN neutrino detectors, an effect which probably requires $\mu_\nu(\mathrm{Dirac}) \lesssim 10^{-12}\,\mu_B$ from the SN 1987A signal [106,175].

In a medium the refractive energy shift for active neutrinos relative to sterile ones creates a barrier to the spin precession [177]. The mass difference has the same effect if the precession is between different flavors through a transition moment [176]. However, the mass and refractive terms may cancel, leading to resonant spin-flavor oscillations [178–180]. This mechanism can explain all solar neutrino data [181,182], but requires rather large toroidal magnetic fields in the Sun. For Majorana neutrinos, the spin-flavor precession amounts to transitions between neutrinos and antineutrinos so that the observation of anti-neutrinos from the Sun would be a diagnostic for this effect [183–185].

Large magnetic fields exist in SN cores so that spin-flavor precession could play an important role, with possible consequences for the explosion mechanism, r-process nucleosynthesis, or the measurable neutrino signal [186–190]. The downside of this richness of phenomena is that there are so many unknown parameters (electromagnetic neutrino properties, masses, mixing angles) as well as the unknown magnetic field strength and distribution that it is difficult to come up with reliable limits or requirements on neutrino properties.

5.3. Right-handed currents

Right-handed (r.h.) weak interactions may exist on some level, e.g. in left–right symmetric models where the r.h. gauge bosons differ from the standard ones by their mass. In the low-energy limit relevant for stars one may account for the new couplings by a r.h. Fermi constant εG_F where ε is a small dimensionless parameter. In left–right symmetric models, one finds explicitly for charged-current processes $\varepsilon_{CC}^2 = \zeta^2 + [m(W_L)/m(W_R)]^2$ where $m(W_{L,R})$ are the l.h. and r.h. gauge boson masses and ζ is the left–right mixing parameter [102].

Assuming that neutrinos are Dirac particles, a SN core loses energy into r.h. states as an "invisible channel" by the process $e + p \to n + \nu_{e,R}$. The SN 1987A energy-loss argument (Section 4.3) then requires $\varepsilon_{CC} \lesssim 10^{-5}$ [15,71,102]. Laboratory experiments yield only $\varepsilon_{CC} \lesssim 3\times10^{-2}$ [191], but do not depend on the assumed existence of r.h. neutrinos. For neutral currents, the dominant emission process is $NN \to NN\nu_R\bar\nu_R$ which is subject to saturation effects as in the case of axion emission [76]. One then finds $\varepsilon_{NC} \lesssim 3\times10^{-3}$ [15], somewhat less restrictive than the original limits of [71,102]. This bound is somewhat less restrictive than $\varepsilon_{NC} \lesssim 10^{-3}$ found from big-bang nucleosynthesis [192].

6. Axions and other pseudoscalars

6.1. Interaction structure

New spontaneously broken global symmetries imply the existence of Nambu–Goldstone bosons that are massless and as such present the most natural case (besides neutrinos) for using stars as particle-physics laboratories. Massless scalars would lead to new long-range forces so that we may focus here on pseudoscalars. The most prominent example are axions which were proposed more than twenty years ago as a solution to the strong CP problem [193–195]; for reviews see [196,197] and for the latest developments the proceedings of a topical conference [198]. We use axions as a generic example – it will be obvious how to extend the following results and discussions to other cases.

Actually, axions are only "pseudo Nambu–Goldstone bosons" in that the spontaneously broken chiral Peccei–Quinn symmetry $U_{\rm PQ}(1)$ is also explicitly broken, endowing these particles with a small mass $m_a = 0.60\,\mathrm{eV}\ 10^7\,\mathrm{GeV}/f_a$. Here, f_a is the Peccei–Quinn scale, an energy scale which is related to the vacuum expectation value of the field that breaks $U_{\rm PQ}(1)$. The properties of Nambu–Goldstone bosons are always related to such a scale which is the main quantity to be constrained by astrophysical arguments, while the (m_a–f_a)-relationship is specific to axions.

In order to calculate the axionic energy-loss rate from stellar plasmas one needs to specify the interaction with the medium constituents. The interaction with a fermion j (mass m_j) is generically $\mathscr{L}_{\rm int} = (C_j/2f_a)\,\overline{\Psi}_j\gamma^\mu\gamma_5\Psi_j\partial_\mu a$ or $-\mathrm{i}(C_j m_j/f_a)\,\overline{\Psi}_j\gamma_5\Psi_j a$, where Ψ_j is the fermion and a the axion field and C_j is a model-dependent coefficient of order unity. The combination $g_{aj} \equiv C_j m_j/f_a$ plays the role of a Yukawa coupling and $\alpha_{aj} \equiv g_{aj}^2/4\pi$ acts as an "axionic fine structure constant". The derivative form of the interaction is more fundamental in that it is invariant under $a \to a + a_0$ and thus respects the Nambu–Goldstone nature of these particles. The pseudoscalar form is usually equivalent, but one has to be careful when calculating processes where two Nambu–Goldstone bosons are attached to one fermion line, for example an axion and a pion attached to a nucleon [71,199–201].

The dimensionless couplings C_i depend on the detailed implementation of the Peccei–Quinn mechanism. Limiting our discussion to "invisible axion models", where f_a is much larger than the scale of electroweak symmetry breaking, it is conventional to distinguish between models of the DFSZ type [202,203] and of the KSVZ type [204,205]. In KSVZ models, axions have no tree-level couplings to the standard quarks or leptons, yet axions couple to nucleons by their generic mixing with the neutral pion. The latest analysis gives numerically [77] $C_p = -0.34$ and $C_n = 0.01$, with a statistical uncertainty of about ± 0.04 and an estimated systematic uncertainty of roughly the same magnitude. The tree-level couplings to standard quarks and leptons in the DFSZ model depend on an angle β which measures the ratio of vacuum expectation values of two Higgs fields. One finds [77] $C_e = \frac{1}{3}\cos^2\beta$, $C_p = -0.07 - 0.46\cos^2\beta$, and $C_n = -0.15 + 0.38\cos^2\beta$, with similar uncertainties as in the KSVZ case.

The CP-conserving interaction between photons and pseudoscalars is commonly expressed in terms of an inverse energy scale $g_{a\gamma}$ according to $\mathscr{L}_{\rm int} = \frac{1}{4}g_{a\gamma}F_{\mu\nu}\widetilde{F}^{\mu\nu}a = -g_{a\gamma}\mathbf{E}\cdot\mathbf{B}a$, where F is the electromagnetic field-strength tensor and $\widetilde{F}$ its dual. For axions

$$g_{a\gamma} = \frac{\alpha}{2\pi f_a}\,C_\gamma, \quad C_\gamma = \frac{E}{N} - 1.92 \pm 0.08\ , \tag{10}$$

where E/N is the ratio of the electromagnetic and over color anomalies, a model-dependent ratio of small integers. In the DFSZ model or grand unified models one has $E/N = 8/3$, for which $C_\gamma \approx 0.75$, but one can also construct models with $E/N = 2$, which significantly reduces the axion–photon coupling [206].

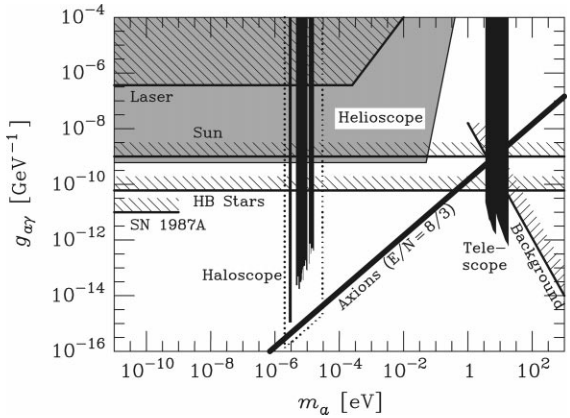

Fig. 7. Limits to the axion–photon coupling $g_{a\gamma}$ after [14,15,208,209]. The "haloscope" search assumes that these particles are the galactic dark matter; the dotted region marks the sensitivity range of the ongoing dark-matter axion searches.

6.2. Limits on the interaction strength

6.2.1. Photons

The axion interaction with fermions or photons allows for numerous reactions which can produce axions in stars, which may imply limits on the axion coupling strength. Beginning with photons, pseudoscalars have a two-photon coupling which allows for the decay $a \to 2\gamma$ and also, in stellar plasmas, for the Primakoff conversion $\gamma \leftrightarrow a$ in the electric fields of electrons and nuclei (Fig. 1). The helioseismological constraint on solar energy losses leads to Eq. (1) as a bound on $g_{a\gamma}$. Fig. 7 shows this constraint ("Sun") in the context of other bounds. For axions the relationship between $g_{a\gamma}$ and m_a is indicated by the heavy solid line, assuming $E/N = 8/3$.

One may also search directly for solar axions. One method ("helioscope") is to direct a dipole magnet toward the Sun, allowing solar axions to mutate into X-rays by the inverse Primakoff process [210,211]. A first pilot experiment was not sensitive enough [212], but the exposure time was significantly increased in a new experiment in Tokyo where a dipole magnet was gimballed like a telescope so that it could follow the Sun [213,214]. The resulting limit $g_{a\gamma} \lesssim 6 \times 10^{-10}\,\mathrm{GeV}^{-1}$ is more restrictive than Eq. (1). An intruiging project (SATAN) at CERN would use a decommissioned LHC test magnet that could be mounted on a turning platform to achieve reasonable periods of alignment with the Sun [215]. This setup could begin to compete with the globular-cluster limit of Eq. (11).

The axion–photon transition in a macroscopic magnetic field is analogous to neutrino oscillations and thus depends on the particle masses [216]. For a large mass difference the transition is suppressed by the momentum mismatch of particles with equal energies. Therefore, the Tokyo limit applies only for $m_a \lesssim 0.03\,\mathrm{eV}$. In a next step one will fill the helioscope with a pressurized gas, giving the photon a dispersive mass to overcome the momentum mismatch. An alternative method is "Bragg diffraction", which uses the strong electric field of a crystal lattice which has large Fourier components for the required momentum transfer [217,218].

The Primakoff conversion of stellar axions can also proceed in the magnetic fields of Sun spots or in the galactic magnetic field so that one might expect anomalous X- or γ-ray fluxes from the Sun [219], the red supergiant Betelgeuse [220], or SN 1987A [221,222]. Observations of SN 1987A yield $g_{a\gamma} \lesssim 0.1 \times 10^{-10}\,\mathrm{GeV}^{-1}$ for nearly massless pseudoscalars with $m_a \lesssim 10^{-9}\,\mathrm{eV}$. A similar limit is obtained from the isotropy of the cosmic X-ray background which would be modified by its Primakoff conversion in the galactic B field [223].

The most important limit derives from the helium-burning lifetime of HB stars in globular clusters, i.e. from Eq. (3),

$$g_{a\gamma} \lesssim 0.6 \times 10^{-10}\,\mathrm{GeV}^{-1}\ . \tag{11}$$

For $m_a \gtrsim 10$ keV this limit quickly degrades as the emission is suppressed when the particle mass exceeds the stellar temperature. Eq. (11) was first stated in [15], superseding the slightly less restrictive but often-quoted "red-giant bound" of [29]. The axion relation Eq. (10) leads to

$$m_a C_\gamma \lesssim 0.3\,\mathrm{eV} \quad \text{and} \quad f_a/C_\gamma \gtrsim 2 \times 10^7\,\mathrm{GeV}\ . \tag{12}$$

In the DFSZ and grand unified models, $C_\gamma \approx 0.75$ so that $m_a \lesssim 0.4\,\mathrm{eV}$ (Fig. 8). For models with $E/N = 2$, the bounds are much weaker.

On the basis of their photon coupling alone, pseudoscalars can reach thermal equilibrium in the early universe. Their subsequent $a \to 2\gamma$ decays would contribute to the cosmic photon backgrounds [208], excluding a non-trivial m_a–$g_{a\gamma}$-range. Some of the pseudoscalars would end up in galaxies and clusters of galaxies. Their decay would produce an optical line feature that was not found [170,224,225], leading to the "telescope" limits in Fig. 7. For axions, the telescope limits exclude an approximate mass range 4–14 eV even for small C_γ.

Axions with a mass in the μeV (10^{-6} eV) range could be the dark matter of the universe (Section 6.3). The Primakoff conversion in a microwave cavity placed in a strong magnetic field ("haloscope") allows one to search for galactic dark-matter axions [210]. Two pilot experiments [226,227] and first results from a full-scale search [228] already exclude a range of coupling strength shown in Fig. 7. The new generation of experiments [198,228–230] should cover the dotted area in Fig. 7, perhaps leading to the discovery of axion dark matter.

6.2.2. Electrons

Pseudoscalars which couple to electrons are produced by the Compton process $\gamma + e^- \to e^- + a$ [27,28,31] and by the electron bremsstrahlung process $e^- + (A,Z) \to (A,Z) + e^- + a$ [28,40,44]. A standard solar model yields an axion luminosity of [28] $L_a = \alpha_{\mathrm{ae}}\, 6.0 \times 10^{21} L_\odot$ where α_{ae} is the axion electron "fine-structure constant". The helioseismological constraint $L_a \lesssim 0.1\, L_\odot$ (Section 2.2) implies $\alpha_{\mathrm{ae}} \lesssim 2 \times 10^{-23}$. White-dwarf cooling gives [15,39] $\alpha_{\mathrm{ae}} \lesssim 1.0 \times 10^{-26}$, while the most restrictive limit is from the delay of helium ignition in low-mass red-giants [31] in the spirit of Eq. (2), $\alpha_{\mathrm{ae}} \lesssim 0.5 \times 10^{-26}$ or $g_{\mathrm{ae}} \lesssim 2.5 \times 10^{-13}$. For $m_a \gtrsim T \approx 10\,\mathrm{keV}$ this limit quickly degrades because the emission from a thermal plasma is suppressed. For axions one finds

$$m_a C_e \lesssim 0.003\,\mathrm{eV} \quad \text{and} \quad f_a/C_e \gtrsim 2 \times 10^9\,\mathrm{GeV}\ . \tag{13}$$

In KSVZ-type models $C_e = 0$ at tree level so that no interesting limit obtains. In the DFSZ model $m_a \cos^2\beta \lesssim 0.01\,\mathrm{eV}$ and $f_a/\cos^2\beta \gtrsim 0.7 \times 10^9\,\mathrm{GeV}$. Since $\cos^2\beta$ can be very small, there is no generic limit on m_a.

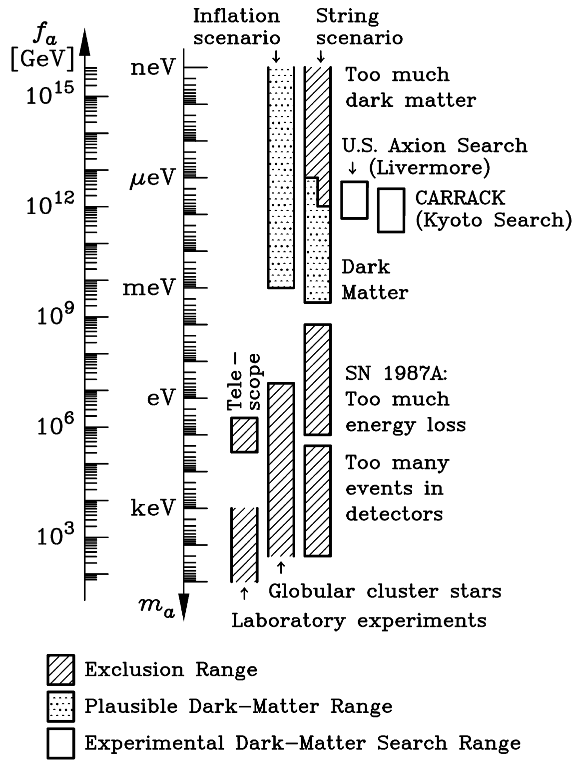

Fig. 8. Astrophysical and cosmological exclusion regions (hatched) for the axion mass m_a, or equivalently the Peccei–Quinn scale f_a. The globular-cluster limit depends on the axion–photon coupling; it was assumed that $E/N = 8/3$ as in GUT models or the DFSZ model. The SN 1987A limits depend on the axion–nucleon couplings; the shown case corresponds to the KSVZ model and approximately to the DFSZ model. The dotted "inclusion regions" indicate where axions could plausibly be the cosmic dark matter. Most of the allowed range in the inflation scenario requires fine-tuned initial conditions. In the string scenario the plausible dark-matter range is somewhat controversial as indicated by the step in the low-mass end of the "inclusion bar". Also shown is the projected sensitivity range of the search experiments for galactic dark-matter axions.

6.2.3. Nucleons

The axion–nucleon coupling strength is primarily constrained by the SN 1987A energy-loss argument [70–77]. The main problem is to estimate the axion emission rate reliably. In the early papers it was based on a somewhat naive calculation of the bremsstrahlung process $NN \to NNa$, using quasi-free nucleons that interact perturbatively through a one-pion exchange potential. Assuming an equal axion coupling g_{aN} to protons and neutrons this treatment leads to the g_{aN}-dependent shortening of the SN 1987A neutrino burst of Fig. 3. However, in a dense medium the bremsstrahlung process likely saturates, reducing the naive emission rate by as much as an order of magnitude [76]. With this correction, and assuming that the neutrino burst was not shortened by more than half, one reads from Fig. 3 an excluded range $3 \times 10^{-10} \lesssim g_{aN} \lesssim 3 \times 10^{-7}$. This implies an exclusion range

$$0.002\,\mathrm{eV} \lesssim m_a C_N \lesssim 2\,\mathrm{eV}, \quad 3 \times 10^6\,\mathrm{GeV} \lesssim f_a/C_N \lesssim 3 \times 10^9\,\mathrm{GeV}\ . \tag{14}$$

For KSVZ axions the coupling to neutrons disappears while $C_p \approx -0.34$. With a proton fraction of about 0.3 one estimates an effective $C_N \approx 0.2$ so that [15,76]

$$0.01\,\mathrm{eV} \lesssim m_a \lesssim 10\,\mathrm{eV}, \quad 0.6\times 10^6\,\mathrm{GeV} \lesssim f_a \lesssim 0.6\times 10^9\,\mathrm{GeV} \tag{15}$$

is excluded. In a detailed numerical study the values for C_n and C_p appropriate for the KSVZ model and for the DFSZ model with different choices of $\cos^2\beta$ were implemented [77]. For KSVZ axions one finds a limit $m_a \lesssim 0.008\,\mathrm{eV}$, while it varies between about 0.004 and 0.012 eV for DFSZ axions, depending on $\cos^2\beta$. In view of the large overall uncertainties it is probably good enough to remember $m_a \lesssim 0.01\,\mathrm{eV}$ as a generic limit (Fig. 8).

Axions on the "strong interaction side" of the exclusion range would have produced excess counts in the neutrino detectors by their absorption on oxygen if $1\times 10^{-6} \lesssim g_{aN} \lesssim 1\times 10^{-3}$ [108]. For KSVZ axions this crudely translates into $20\,\mathrm{eV} \lesssim m_a \lesssim 20\,\mathrm{keV}$ as an exclusion range (Fig. 8).

6.3. Cosmological limits

The astrophysical axion mass limits are particularly interesting when juxtaposed with the cosmological ones. For $f_a \gtrsim 10^8\,\mathrm{GeV}$ cosmic axions never reach thermal equilibrium in the early universe. They are produced by a nonthermal mechanism that is intimately intertwined with their Nambu–Goldstone nature and that implies that their contribution to the cosmic density is proportional to $f_a^{1.175}$ and thus to $m_a^{-1.175}$. The requirement not to "overclose" the universe with axions thus leads to a *lower* mass limit.

One must distinguish between two generic cosmological scenarios. If inflation occurred after the Peccei–Quinn symmetry breaking or if $T_{\mathrm{reheat}} < f_a$, the initial axion field takes on a constant value $a_\mathrm{i} = f_a \Theta_\mathrm{i}$ throughout the universe, where $0 \le \Theta_\mathrm{i} < \pi$ is the initial "misalignment" of the Θ parameter [231–234]. If $\Theta_\mathrm{i} \sim 1$ one obtains a critical density in axions for $m_a \sim 1\,\mu\mathrm{eV}$, but since Θ_i is unknown there is no strict cosmological limit on m_a. However, the possibility to fine-tune Θ_i is limited by inflation-induced quantum fluctuations which in turn lead to temperature fluctuations of the cosmic microwave background [235–238]. In a broad class of inflationary models one thus finds an upper limit to m_a where axions could be the dark matter. According to the most recent discussion [238] it is about $10^{-3}\,\mathrm{eV}$ (Fig. 8).

If inflation did not occur at all or if it occurred before the Peccei–Quinn symmetry breaking with $T_{\mathrm{reheat}} > f_a$, cosmic axion strings form by the Kibble mechanism [239,240]. Their motion is damped primarily by axion emission rather than gravitational waves. After axions acquire a mass at the QCD phase transition they quickly become nonrelativistic and thus form a cold dark matter component. Unknown initial conditions no longer enter, but details of the string mechanism are sufficiently complicated to prevent an exact prediction of the axion density. On the basis of Battye and Shellard's treatment [241,242] and assuming that axions are the cold dark matter of the universe one finds a plausible mass range of $m_a = 6$–$2500\,\mu\mathrm{eV}$ [14]. Sikivie et al. [207,243,244] predict somewhat fewer axions, allowing for somewhat smaller masses if axions are the dark matter. Either way, the ongoing search experiments for galactic dark matter axions in Livermore (U.S. Axion Search [228]) and in Kyoto (CARRACK [229,230]) aim at a cosmologically well-motivated range of axion masses (Fig. 8).

7. Conclusion

Stellar-evolution theory in conjunction with astronomical observations, the SN 1987A neutrino burst, and certain X- and γ-ray observations provide a number of well-developed arguments to constrain the properties of low-mass particles. The most successful examples are globular-cluster stars where the "energy-loss argument" was condensed into the simple criteria of Eqs. (2) and (3) and SN 1987A where it was summarized by Eq. (5). New particle-physics conjectures must first pass these and other simple astrophysical standard tests before being taken too seriously.

A showcase example for the interplay between astrophysical limits with laboratory experiments and cosmological arguments is provided by the axion hypothesis. The laboratory and astrophysical limits push the Peccei–Quinn scale to such high values that it appears almost inevitable that axions, if they exist at all, play an important role as a cold dark matter component. This makes the direct search for galactic axion dark matter a well-motivated effort. Other important standard limits pertain to neutrino electromagnetic form factors – laboratory experiments will have a difficult time catching up.

Most of the theoretical background relevant to this field could not be touched upon in this brief overview. The physics of weakly coupled particles in stars is a nice playing field for "particle physics in media" which involves field theory at finite temperature and density (FTD), many-body effects, particle dispersion and reactions in magnetic fields and media, oscillations of trapped neutrinos, and so forth. In the context of SN theory such issues are naturally of particular interest, but even the plasmon decay $\gamma \to \nu\bar{\nu}$ in normal stars or the MSW effect in the Sun are interesting cases. Particle physics in media and its astrophysical and cosmological applications is a fascinating topic in its own right which well deserves a dedicated review.

Much more information of particle-physics interest may be written in the sky than has been deciphered as yet. Other objects or phenomena should be considered, perhaps other kinds of conventional stars, perhaps more exotic phenomena such as γ-ray bursts. The particle-physics lessons to be learned from them are left to be reviewed in a future report!

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