Physics Letters B 751 (2015) 209–214

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Production of unstable heavy neutrinos in proto-neutron stars

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ARTICLE INFO

Article history:
Received 4 September 2015
Accepted 11 October 2015
Available online 19 October 2015
Editor: J. Hisano

ABSTRACT

We discuss the production of a class of heavy sterile neutrinos ν_h in proto-neutron stars. The neutrinos, of mass around 50 MeV, have a negligible mixing with the active species but relatively large dimension-5 electromagnetic couplings. In particular, a magnetic dipole moment $\mu \approx 10^{-6}$ GeV^{-1} implies that they are thermally produced through $e^+e^- \to \bar{\nu}_h \nu_h$ in the early phase of the core collapse, whereas a heavy–light transition moment $\mu_{\rm tr} \approx 10^{-8}$ GeV^{-1} allows their decay $\nu_h \to \nu_i \gamma$ with a lifetime around 10^{-3} s. This type of electromagnetic couplings has been recently proposed to explain the excess of electron-like events in baseline experiments. We show that the production and decay of these heavy neutrinos would transport energy from the central regions of the star to distances $d \approx 400$ km, providing a very efficient mechanism to enhance the supernova shock front and heat the material behind it.

1. Introduction

Neutrinos define a sector of the Standard Model that still presents some important unknowns. The current scheme of mass differences and mixings seems able to explain most of the existing data [1], but the absolute value of their masses, their Dirac or Majorana nature [2] or the presence of additional sterile modes [3,4] are yet to be determined. In particular, the production of sterile neutrinos ν_s, through collisions with standard matter or through flavor oscillations has important implications both in particle physics and astrophysics [5–8]. The mixing with an active neutrino ν may provide sterile modes with small couplings to the W and Z gauge bosons that translate into dimension-6 operators of type

$$-\mathcal{L}_{\rm eff} = \frac{G_F \sin\theta}{\sqrt{2}} \bar{f} \gamma_\mu (C_V - C_A \gamma_5) f \, \bar{\nu}_s \gamma^\mu (1-\gamma_5) \nu + \text{h.c.} \qquad (1)$$

In addition, the low-energy effective Lagrangian may also include dimension-5 operators from loops involving heavy particles. Although these operators are usually overlooked, they could mediate the dominant reactions of sterile neutrinos in a star under favorable thermodynamical conditions. Here we will study this possibility in the context of supernova explosions.

When a supernova goes off a proto-neutron star can be formed having a typical initial radius of (20–60) km and a mass of (1–1.5) $M_\odot$. It is believed that most of the gravitational binding energy ($E_{\rm grav} \approx 3 \times 10^{53}$ erg) is released in a ∼20 second neutrino burst [9]. The neutrino spectrum from supernova SN1987A detected at SuperK and IMB indicated a weak decoupling from baryonic matter, confirming that neutrino transparency sets in as their temperature falls below a few MeV [10] in the dense core. At earlier phases of the collapse, however, computational simulations [11,12] reveal internal peak temperatures exceeding 20 MeV in the central high density regions of the star. At such temperatures and densities the evolution of these astrophysical objects becomes sensitive to the fundamental properties of neutrinos and to the presence of hypothetical weakly-coupled particles. In this context, a lot of effort has been devoted to the cooling through neutrino emission in the nuclear medium [13], to the matter opacity [14] and revival of the stalled shock that arises in the standard paradigm of supernova core collapse [15], or to the synthesis of heavy nuclei taking place in the hot bubble behind the shock [16,17].

In this work we will focus on the astrophysical consequences of the production of a heavy sterile neutrino ν_h whose dominant interactions are *not* the weak ones in Eq. (1) but of electromagnetic kind. This type of particles have been proposed as a possible explanation [18] for the excess of electron-like events in baseline experiments [19]. Let us briefly show how the required couplings could be generated. Consider an $SU(2)_L$-singlet Dirac neu-

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http://dx.doi.org/10.1016/j.physletb.2015.10.026

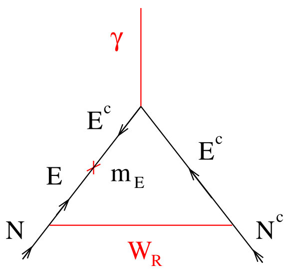

Fig. 1. Diagram contributing to the magnetic dipole moment μ of ν_h.

trino, ν_h, of mass $m_h = 50$ MeV. We will denote by N and N^c the (2-component) neutrino and antineutrino spinors defining ν_h,

$$\nu_h = \begin{pmatrix} N \\ \bar{N}^c \end{pmatrix}. \tag{2}$$

Let us also suppose that at TeV energies the gauge symmetry is $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ and that ν_h is accommodated within two $SU(2)_R$ doublets together with a charged lepton,

$$L = \begin{pmatrix} N \\ E \end{pmatrix} \quad L^c = \begin{pmatrix} E^c \\ N^c \end{pmatrix}. \tag{3}$$

In order to avoid collider bounds [20], the breaking of the left–right symmetry must be such that the charged lepton (E, E^c) gets a mass $m_E \geq 300$ GeV while ν_h remains light. Loop diagrams of heavy gauge bosons and fermions (see Fig. 1) will then generate the operator

$$-\mathcal{L}_{\text{eff}} = \mu\, \bar{\nu}_h \sigma_{\mu\nu} \nu_h\, \partial^\mu A^\nu\,, \tag{4}$$

where A^ν is the electromagnetic field and μ is a magnetic dipole moment of order [21]

$$\mu \approx e\, \frac{g_R^2}{16\pi^2}\, \frac{m_E}{M_R^2} \approx 10^{-6}\ \text{GeV}^{-1}\,. \tag{5}$$

In addition, the possible mixing of the sterile and the active neutrinos will be parametrized by an angle θ, so that the mass eigenstates read $N' = \cos\theta\, N + \sin\theta\, \nu$ and $\nu' = -\sin\theta\, N + \cos\theta\, \nu$. This mixing will generate electromagnetic transitions through the same type of diagrams (we drop the prime to indicate mass eigenstates):

$$\mathcal{L}_{\text{eff}} = \frac{1}{2}\, \mu_{\text{tr}}\, \bar{\nu}_h\, \sigma_{\mu\nu}\, (1-\gamma_5)\, \nu\, \partial^\mu A^\nu + \text{h.c.}\,, \tag{6}$$

with $\mu_{\text{tr}} \approx \sin\theta\mu$ being the transition dipole moment. This operator may imply that the dominant decay mode of the heavy neutrino is $\nu_h \to \nu\gamma$. Notice also that the presence of additional heavy singlets ($\nu_{h'}$ with $m_{h'} \approx m_E$) mixed both with ν and ν_h will give additional contributions to μ_{tr}. Therefore, at this point we will treat μ and μ_{tr} as independent parameters.

This type of sterile neutrinos could change substantially the evolution of a proto-neutron star. We will show that sterile pairs can be produced abundantly during the ∼20 second neutrino burst, escape the star core more easily than standard neutrinos, and finally decay within a few hundred km from the core. The very energetic photons from the decay could deposit energy, helping revive the stalled accretion shock formed during the collapse and change the thermal environment in the vicinity of the star. Our scenario could be considered a different realization of the *eosphoric* neutrino hypothesis proposed in [22].

2. Decay rate, production and scattering cross sections

Let us first describe the dominant decay and production channels for the heavy neutrino ν_h in vacuum. Later we will discuss how the hot and dense medium in a proto-neutron star (including populations of neutrons, protons, electrons and muons) affects these processes.

To be definite in our calculation we will take as reference values $m_h = 50$ MeV, $\mu = 10^{-6}$ GeV^{-1} $= 3.3 \times 10^{-9}\mu_B$, and $\mu_{\text{tr}} \approx 10^{-8}$ GeV^{-1}. We use $\hbar = c = 1$. For these values of the mass and the transition moment, the heavy neutrino will decay into $\nu\gamma$ with a lifetime

$$\tau_h = \frac{16\pi}{\mu_{\text{tr}}^2\, m_h^3} = \frac{(50\ \text{MeV})^3}{m_h^3} \times \frac{(10^{-8}\ \text{GeV}^{-1})^2}{\mu_{\text{tr}}^2} \times 0.0026\ \text{s}\,. \tag{7}$$

We will also assume that the mass mixing, i.e., the active component in ν_h, is smaller than $\sin^2\theta < 10^{-3}$ and only along the muon and/or the tau flavors. In that case, the radiative decay will dominate over the weak processes $\nu_h \to \nu e^+ e^-$, $\nu\nu_i\bar{\nu}_i$, which appear with a branching ratio

$$\text{BR}(\nu_h \to \nu e^+ e^-) \approx \frac{\sin^2\theta}{10^{-3}} \times \frac{(10^{-8}\ \text{GeV}^{-1})^2}{\mu_{\text{tr}}^2} \times \frac{m_h^2}{(50\ \text{MeV})^2} \times 0.05\%\,. \tag{8}$$

This type of sterile neutrino avoids cosmological bounds since it decays before primordial nucleosynthesis. At colliders it is hardly detectable: even if it were produced in 1% of kaon or muon decays, ν_h is too long lived to decay inside the detectors and too light to change significantly the kinematics of the decay [23]. Actually, more elaborate setups with two sterile modes have been proposed to explain the excess of electron-like events at MiniBooNE in terms of the photon that results from its decay [18].

The dominant production channels of ν_h will also be electromagnetic. In particular, electron–positron annihilation into ν_h pairs, $e^+e^- \to \bar{\nu}_h\nu_h$, will be mediated by a photon through the magnetic dipole moment coupling in Eq. (4). The differential cross section is given by

$$\frac{d\sigma}{dt} = \frac{\alpha\mu^2}{s^2 - 4sm_e^2} \times \left(-t + 2m_h^2 + m_e^2 - \frac{t^2 - 2(m_h^2 + m_e^2)t + (m_h^2 - m_e^2)^2}{s} \right), \tag{9}$$

where α is the fine structure constant, m_e is the electron mass, and s and t are the usual Mandelstam variables ($\sqrt{s}$ is the center-of-mass energy). In Fig. 2 (left panel) we plot the total cross section for this process. Muon pair annihilation will give an analogous but subleading contribution, since muons are less abundant than electrons in the star core.

The active to sterile transition mediated by a photon can be catalyzed by the presence of charged particles $X = p, e$: $\nu X \to \nu_h X$ (see right panel in Fig. 2). This contribution, however, can be neglected here due to the smaller value of the transition coupling that we have assumed, $\mu_{\text{tr}} \approx 10^{-2}\mu$. The weak channels which dominate the production of active neutrinos [24] give also a subleading contribution due to the small mixing $\sin^2\theta < 10^{-3}$ of our sterile, whereas other processes like plasmon decay [25] are irrelevant for heavy neutrino masses around 50 MeV (i.e., much larger than the electron mass).

In addition to its production and decay, the collisions of ν_h with charged matter will be essential in order to understand its propagation in the dense medium and estimate how efficiently these

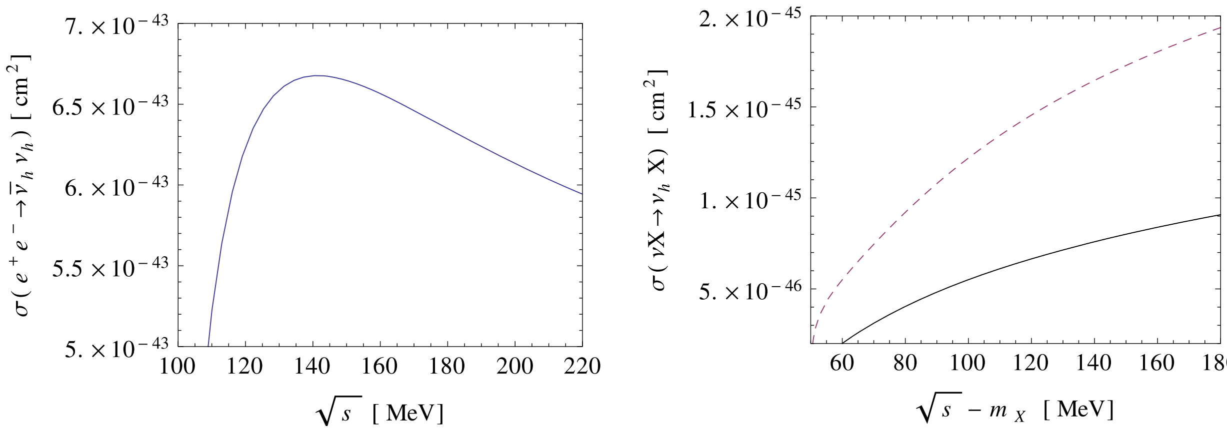

Fig. 2. Left: Total cross section $\sigma(e^+e^- \to \bar{\nu}_h \nu_h)$ for $m_h = 50$ MeV and $\mu = 10^{-6}$ GeV^{-1}. Right: $\sigma(\nu X \to \nu_h X)$ for $\mu_{\rm tr} = 10^{-8}$ GeV^{-1} (solid: $X = p$, dashes: $X = e$).

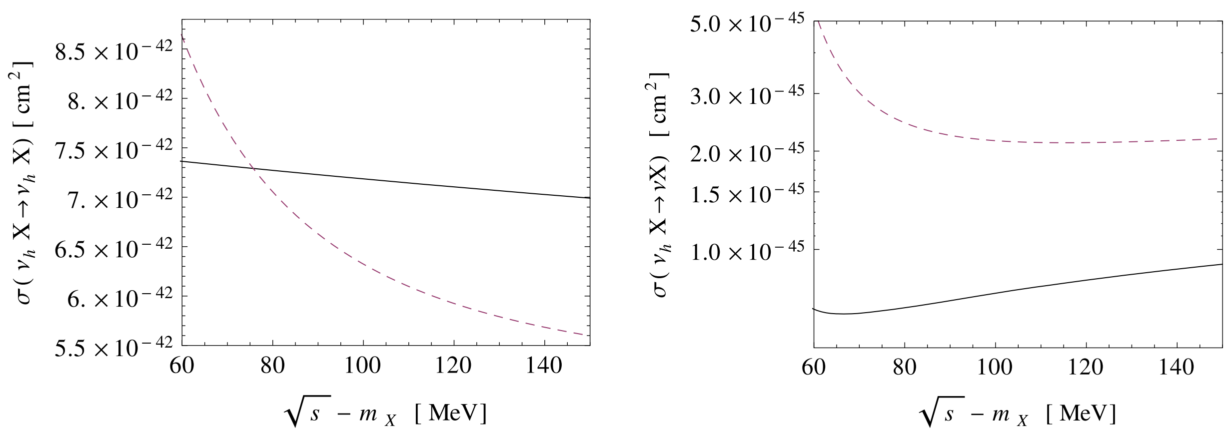

Fig. 3. Left: $\sigma(\nu_h X \to \nu_h X)$ for $m_h = 50$ MeV, $\mu = 10^{-6}$ GeV^{-1} and a scattering angle $\theta > 30°$ in the center-of-mass frame. Right: $\sigma(\nu_h X \to \nu X)$ for $\mu_{\rm tr} = 10^{-8}$ GeV^{-1} (solid: $X = p$, dashes: $X = e$).

neutrinos escape the proto-neutron star. We need to distinguish between elastic scatterings

$$\nu_h\, X \to \nu_h\, X \tag{10}$$

and absorption reactions of type

$$\nu_h\, X \to \nu\, X\,. \tag{11}$$

The differential cross section for the first process reads

$$\frac{{\rm d}\sigma}{{\rm d}t} = \frac{\alpha\mu^2}{s^2 - 2s\left(m_X^2 + m_h^2\right) + \left(m_h^2 - m_X^2\right)^2} \times \left(-s + 2m_h^2 + m_X^2 - \frac{s^2 - 2(m_h^2 + m_X^2)s + (m_h^2 - m_X^2)^2}{t}\right). \tag{12}$$

This is a long distance (photon-mediated) process with a divergent total cross section; if we restrict to collisions substantially changing the direction of the incident ν_h (e.g., a scattering angle $\theta > 30°$ in the center-of-mass frame) we obtain the cross section depicted in Fig. 3 (left panel). For the inelastic process that transforms the heavy neutrino into an active one we obtain

$$\frac{{\rm d}\sigma}{{\rm d}t} = \frac{\alpha\mu_{\rm tr}^2}{2\left(s^2 - 2s\left(m_X^2 + m_h^2\right) + \left(m_h^2 - m_X^2\right)^2\right)} \times \left(-s + \frac{1}{2}(m_h^2 + 2m_X^2) - \frac{s^2 - (m_h^2 + 2m_X^2)s + \frac{1}{2}m_h^4 + m_X^4}{t} - \frac{m_X^2 m_h^4}{t^2}\right), \tag{13}$$

and its total cross section is also depicted in Fig. 3 (right panel). These cross sections are, in both cases, much smaller than the ones for active neutrinos off nucleons mediated by weak bosons. For example, a neutrino in a gas at $T \simeq 20$ MeV has an average energy $\langle E_\nu \rangle \simeq \pi T \approx 60$ MeV, and its lowest order elastic cross section with a neutron is $\sigma(\nu_i n \to \nu_i n) \approx G_F^2 E_\nu^2 (3C_A^2 + C_V^2)/\pi \simeq 3 \times 10^{-40}$ cm^2, with $i = e, \mu, \tau$ (the cross section with a proton is approximately hundred times smaller). As for the absorption of a ν_e through a charged current interaction, we have $\sigma(\nu_e n \to ep) \approx 2 \times 10^{-39}$ cm^2.

3. Production in a proto-neutron star

Let us now calculate the production rate of heavy neutrinos at the astrophysical site. Nucleon and lepton densities in the medium are constrained by electric charge neutrality and baryon and lepton number conservation. Typical baryonic densities in the core of a proto-neutron star are well above nuclear saturation density $n_B \simeq$ (2–3)n_0 with $n_0 = 0.17$ fm^{-3} [9], whereas the lepton and electron fraction evolve dynamically from $Y_L \approx 0.31$, $Y_e \approx 0.27$ at $t = 0.1$ s to $Y_L \sim 0.18$, $Y_e \sim 0.17$ at $t = 10$ s [26]. The extreme conditions in the core are such that quantum effects will be important. The population of baryons and leptons is described by the Fermi–Dirac distribution

$$f_i(E) = \frac{1}{e^{(E-\mu_i)/T} + 1}\,, \tag{14}$$

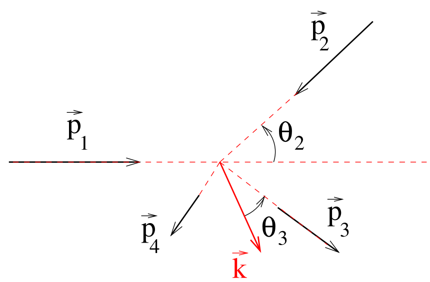

Fig. 4. Kinematical variables in the reaction $12 \to 34$ used in this work.

where μ_i $(i = n, p, e^{\pm})$ denotes the chemical potential of the considered species. Both T and μ_i evolve within the star, in particular, the chemical potentials take care of the conservation of charges and quantum numbers in a self consistent way [27].

For a given value of the temperature and the electron chemical potential ($\mu_{e^+} = -\mu_{e^-}$), the *total energy emissivity* (energy produced per unit volume and unit time) of heavy neutrino pairs through the dominant process $e^+e^- \to \bar{\nu}_h\nu_h$, $Q_E(e^+e^- \to \bar{\nu}_h\nu_h) = \frac{dE}{dtdV}$ is given by [28,29]

$$Q_E = \frac{4}{(2\pi)^8} \int \frac{d^3p_1}{2E_1}\frac{d^3p_2}{2E_2}\frac{d^3p_3}{2E_3}\frac{d^3p_4}{2E_4}\,\delta^4(p_1+p_2-p_3-p_4) \times (E_1+E_2)\,|\bar{\mathcal{M}}|^2\, f(f_1,f_2,f_3,f_4) \quad (15)$$

where the factor $f(f_1,f_2,f_3,f_4) = f_1f_2(1-f_3)(1-f_4)$ includes the Pauli blocking factor in the generic reaction $12 \to 34$ and $p_i = (E_i, \vec{p}_i)$ are the 4-momenta. We will consider that the reaction is not affected by the quenching of outgoing sterile states, i.e., $(1-f_3) \simeq 1 \simeq (1-f_4)$, since heavy neutrinos do not achieve chemical equilibrium and their number density inside the star is always small. The squared matrix element for the interaction defined in Eq. (4) is given by

$$|\bar{\mathcal{M}}(e^+e^- \to \bar{\nu}_h\nu)|^2 = 4e^2\mu_h^2\left(-t + 2m_h^2 + m_e^2 - \frac{t^2 - 2(m_h^2+m_e^2)t + (m_h^2-m_e^2)^2}{s}\right) \quad (16)$$

with e the electron charge, $s = (p_1+p_2)^2$ and $t = (p_1-p_3)^2$, see Fig. 4.

To perform the phase space integral in Eq. (15) we note that there are only four non-trivial independent variables: the initial energies E_1 and E_2, the angle θ_2, as defined in Fig. 4, and E_3. Any other kinematical variables can be derived from these four or can be trivially integrated. It is convenient to define the 4-vector $k = p_1 + p_2$ (notice that $k^2 = s$ and $|\vec{k}|^2 = (E_1+E_2)^2 - s$) and the angle θ_3 of $\vec{k}$ with $\vec{p}_3$:

$$\cos\theta_3 = \frac{2\,(E_1+E_2)\,E_3 - s}{2\,|\vec{k}|\sqrt{E_3^2 - m_h^2}}\,. \quad (17)$$

After integrating the 4-dimensional Dirac delta that enforces energy and momentum conservation, we obtain

$$Q_E = \frac{1}{64\pi^5}\int_{m_e}^{\infty} dE_1 \int_{E_2^{\min}}^{\infty} dE_2 \int_{c_2^{\min}}^{1} d\cos\theta_2 \int_{E_-}^{E_+} dE_3 \times \frac{p_1p_2}{|\vec{k}|}\,(E_1+E_2)\,|\bar{\mathcal{M}}|^2\, f_{e^+}(E_1)\, f_{e^-}(E_2)\,, \quad (18)$$

where the minimum values of E_2 and $\cos\theta_2$, $E_2^{\min}$ and $c_2^{\min}$ respectively, result from the kinematical restriction $s > 4m_h^2$

$$E_2^{\min}(E_1) \approx \sqrt{\frac{m_h^4 + m_e^2E_1^2}{E_1^2 - m_e^2}}\,,$$
$$c_2^{\min}(E_1,E_2) = \mathrm{Max}\left[-1, \frac{2m_h^2 - m_e^2 - E_1E_2}{p_1p_2}\right]. \quad (19)$$

We also define

$$E_\pm(E_1,E_2,\cos\theta_2) = \frac{E_1+E_2}{2} \pm \frac{|\vec{k}|}{2}\sqrt{1 - \frac{4m_h^2}{s}}\,. \quad (20)$$

The Mandelstam variables in terms of these four quantities read

$$s = 2\left(m_e^2 + E_1E_2 + p_1p_2\cos\theta_2\right) \quad (21)$$
$$t = -2E_1E_3 + 2p_1p_3\cos(\theta_3 - \alpha) + m_e^2 + m_h^2 \quad (22)$$

where α $(0 \le \alpha \le \pi)$ is the angle between $\vec{k}$ and $\vec{p}_1$,

$$\alpha = \arctan\left(\frac{p_2\sin\theta_2}{p_1 - p_2\cos\theta_2}\right). \quad (23)$$

4. Transport of energy out of the star core

The possible impact of the heavy neutrino ν_h on the evolution of the proto-neutron star will depend on its ability to take a significant amount of energy out of the core. If the sterile neutrinos are abundant inside the star core but unable to reach the surface before decaying into a photon plus an active neutrino, then they become just a state mediating interactions of electrons with neutrinos and photons. We will show that this is not the case and that they could play an interesting role in supernova explosions.

Let us take a temperature $T_0 = 25$ MeV and an electron chemical potential $\mu_{e0} = 100$ MeV, which are typical values at the inner central region of a proto-neutron star (see [11,12]). Although these quantities are time and density dependent, the chosen values can be used to estimate the possibilities of our scenario.

Varying the mass m_h, the magnetic dipole moment μ, the electron chemical potential μ_e and the temperature T and performing a fit of Q_E in Eq. (18) we obtain

$$Q_E \approx 1.5\times10^{36}\left(\frac{\mu}{10^{-6}\ \mathrm{GeV}^{-1}}\right)^2\left(\frac{T}{25\ \mathrm{MeV}}\right)^{7.4} e^{-\frac{m_h+\mu_e}{3T}}\ \frac{\mathrm{erg}}{\mathrm{s\,cm^3}}\,. \quad (24)$$

For the reference values of all the parameters, the expression above yields $Q_E \approx 2\times10^{35}$ erg/s cm^3, with an average ν_h energy of $\langle E_h\rangle \approx 103$ MeV. This is a very large production rate, $\sim10^2$–10^3 times larger than the one obtained in [22] using heavier sterile neutrinos mixed with the active ones. Our neutrinos, however, will not leave the star core unscattered.

We can also compare this production rate with the one of standard neutrinos in early cooling of proto-neutron stars. For example, in the central core the direct URCA process $n \to pe\bar{\nu}_e$ provides $Q_E^{DURCA} \approx 2.4\times10^{41}R$ erg/s cm^3 at $T = 25$ MeV [28,31], being R a factor of order unity [30]. This is five decades over the ν_h production rate that we have found. The direct URCA process requires a high proton fraction, $Y_p \gtrsim 11\%$, only accessible to large mass objects, however the less demanding modified URCA cooling also gives a much faster rate than for steriles, $Q_E^{MURCA} \approx 1.5\times10^{40}R$ erg/s cm^3. These active neutrinos will be, to a large extent, trapped (before transparency sets in) inside the star core, whereas ours have weaker interactions with the protons and electrons in the medium.

It is then apparent that we need to consider propagation effects of ν_h. Although a precise calculation would require a complete multidimensional simulation that is beyond the scope of this work, we will discuss the qualitative picture and show that the model has enough parameter space to realize it.

The first important effect is the diffusion from the center to the surface of the star core, with a radius $r \approx 20$ km. Active muon and tau neutrinos scatter there mainly off neutrons with a mean free path $\lambda_\nu^S \sim 1$ m [11,32]. Analogously, ν_h will scatter elastically with protons with a cross section ∼40 times smaller, which suggests an interaction length λ_h^S inside the star longer by the same factor. This implies a larger diffusion coefficient $D \approx \lambda_h^S c/3$ and a more efficient transport from the core to the outer parts of the star. The typical diffusion time for that process will be $\tau_D \approx r^2/(2D) \approx 5 \times 10^{-2}$ s. The second crucial effect in the propagation of the heavy neutrinos is their absorption: the star will capture a fraction of them through the inelastic collisions $\nu_h p \to \nu p$. For $\mu_{\rm tr} = 10^{-8}$ GeV^{-1} the absorption length is approximately $\sigma(\nu_h X \to \nu_h X)/\sigma(\nu_h X \to \nu X) \approx 3000$ times larger than λ_h^S, i.e., $\lambda_h^A \approx 120$ km. A final effect, analogous to absorptions, is their decay. Since the values of $\mu_{\rm tr}$ and m_h that we have assumed imply a sterile neutrino lifetime $\tau_h \approx 0.003$ s, 16 times smaller than their diffusion time, a large fraction of the heavy neutrinos produced in the core will decay into $\gamma\nu$ before they have diffused to the outer layers. For a time window of ∼ 20 s energy can be transported to distances $d \simeq \sqrt{2D\tau_D} \simeq 400$ km.

We estimate that, even if absorptions and decays reduced the number of neutrinos leaving the star from our estimate to 1% of the ones produced (some of them closer to the surface), ν_h may still carry a total of 10^{51}–10^{52} erg and deposit this energy outside the star during the 20 second neutrino burst. Notice that a reduction in the magnetic moment μ would also reduce the production rate of heavy neutrinos, but it would increase the mean free path between elastic scatterings and then the fraction of neutrinos that reach the surface. In addition to the transition moment $\mu_{\rm tr}$, the mass m_h is another parameter that could impact the production rate or the decay length of the heavy neutrino.

An interesting variation of the model that depends only on two parameters (m_h and $\mu_{\rm tr}$) would be obtained by suppressing the magnetic moment μ and slightly increasing the transition one, e.g. $\mu_{\rm tr} \approx 5 \times 10^{-8}$ GeV^{-1}. In that case the dominant production channel would be

$$\nu X \to \nu_h X\,; \qquad \bar\nu X \to \bar\nu_h X\,, \tag{25}$$

where X is any charged particle and the active ν may be any linear combination of ν_μ and ν_τ. In the star core the heavy neutrinos would be partially absorbed through the inverse reaction (we estimate an absorption length $\lambda_h^A \approx 4$ km), and the ones reaching the surface would decay with $c\tau_h \approx 30$ km. Since the production would not be so abundant as in the case discussed in the previous section, this possibility provides a *safer* scenario still able to transport energy to the region near the star surface.

Although the optimal value of these parameters would require a full Monte Carlo simulation, the scenario seems flexible enough to introduce acceptable changes in the dynamics of supernova explosions, with the decay into photon plus active neutrino playing an important role in the enhancement of the supernova shock.

5. Conclusions

The collapse of a very massive star defines an astrophysical object with extreme conditions where neutrinos determine the thermodynamics. These proto-neutron stars are suitable laboratories to probe the properties of any weakly coupled particles of mass $\lesssim 100$ MeV. Here we have proposed a sterile neutrino ν_h much heavier than the standard ones ($m_h \approx 50$ MeV) and with sizable electromagnetic couplings: a magnetic dipole moment μ and dipole sterile–active transition $\mu_{\rm tr}$ that mediates its decay $\nu_h \to \nu\gamma$ with a lifetime $\tau_h \approx 10^{-3}$ s. This simple 3-parameter model seems to have interesting implications in supernova explosions. The heavy neutrino is produced in the core at a high rate through $e^+e^- \to \bar\nu_h \nu_h$, it may escape the star more efficiently than active neutrinos and decays depositing a large amount of energy in the outer layers of the star.

We believe that the type of heavy sterile neutrino proposed here could be an essential ingredient to help the progression of the internal shock, which is responsible for the observed supernova events. Full computational simulations could shed more light into the complex energy transport that results from competing processes of scattering, interaction and decay.

Acknowledgements

We would like to thank T. Fischer and J. Pons for useful discussions. This work has been supported by MINECO of Spain (FPA2013-47836, FIS2012-30926) and Consolider-Ingenio **Multidark** CSD2009-00064) and by Junta de Andalucía (FQM225, FQM101,3048).

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