

Mini Z' burst from relic supernova neutrinos and late neutrino masses

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ABSTRACT: In models in which neutrinos are light, due to a low scale of symmetry breaking, additional light bosons are generically present. We show that the interaction between diffuse supernova relic neutrinos (SRN) and the cosmic background neutrinos, via exchange of these light scalars, can result in a dramatic change of the supernova (SN) neutrinos flux. Measurement of this effect with current or future experiments can provide a spectacular direct evidence for the low scale models. We demonstrate how the observation of neutrinos from SN1987A constrains the symmetry breaking scale of the above models. We also discuss how current and future experiments may confirm or further constrain the above models, either by detecting the “accumulative resonance” that diffuse SRN go through or via a large suppression of the flux of neutrinos from nearby $\lesssim \mathcal{O}$ (Mpc) SN bursts.

KEYWORDS: Cosmology of Theories beyond the SM, Neutrino Physics.

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1. Introduction

The observation of neutrino flavor changing from solar [1–3], atmospheric [4] and terrestrial [5, 6] neutrino data has provided firm evidence for neutrino flavor conversion. The recent new Super-Kamiokande (SK) data on the L/E -dependence of atmospheric neutrino events [7], L being the distance traveled by neutrinos of energy E , and the new spectrum data from terrestrial experiments [8, 9], has yielded for the first time evidence of the expected oscillatory behavior. This strongly favors non-vanishing sub-eV neutrino masses. These outstanding developments on the experimental side of neutrino physics have placed a distinct burden on theorists—to understand what is the origin of these tiny neutrino masses.

The most elegant and popular solution to this puzzle is the seesaw mechanism [10]. In this scenario one assumes that lepton number is violated at some high scale Λ_L in the form of right-handed neutrino Majorana masses. This induces, at a lower scale, an effective operator of the form $\mathcal{O}(1) \times (LH)^2/\Lambda_L$, where L denotes a lepton doublet and H the Higgs field. The oscillation data then imply that $\Lambda_L \sim 10^{14}$ GeV. While the seesaw mechanism is very appealing from the theoretical side, it is unlikely to be subject to direct experimental test sometime in the near future. An additional virtue of the seesaw mechanism is that it can naturally provide a platform for generating the observed baryon asymmetry of the universe through leptogenesis [11]. Introduction of such a high scale, however, requires a mechanism for electroweak-symmetry-breaking-scale stabilization which typically leads to various moduli/gravitino problems in the context of cosmology. Thus it is important to explore alternate origins for neutrino masses.

One such alternative is the late neutrino mass framework that induces small neutrino masses due to a low scale of symmetry breaking [12]. This idea points to a completely different understanding for the origin of neutrino masses. The neutrino masses are protected by some flavor symmetry different from the one related to the charged fermion masses. When this symmetry is (say spontaneously) broken by a set of flavor symmetry breaking vevs, f , of fields ϕ , the neutrinos acquire masses from $\left(\frac{\phi}{M_F}\right)^n LNH$ for Dirac neutrinos, or $\left(\frac{\phi}{M_F}\right)^n LNH + M_R NN$ for Majorana neutrinos, where N denotes a right handed neutrino, L, H stand for the SM lepton doublet and Higgs fields respectively and M_F is a scale in which flavor dynamics takes place [13]. We want to stress that these textures do not depend on the details of the symmetry mechanisms, whether global [12] or gauge [14]. Furthermore a similar scenario can be realized via strongly coupled dynamics where the compositeness scale is given by f [15, 16].

With this alternate scenario, it is then of immediate import to delimit the allowed range for the symmetry breaking scale, f at which new physics (NP) appears. Since the principal consequences of the symmetry breaking are neutrino masses and the relevant new degrees of freedom couple only to neutrinos, direct experimental limits on the parameters of this model are unlikely to be attained. In fact, the strongest limits on f come from cosmology and astrophysics rather than from laboratory data [12]. As will be discussed shortly, there are generically associated with this mechanism some extra light degrees of freedom. In the case that the number of these exceed present bounds from big bang nucleosynthesis (BBN), the requirement that these not be in thermal equilibrium during BBN gives a limit on f of approximately [12, 17, 16, 18, 14]

$$f \gtrsim 10 \text{ keV} . \tag{1.1}$$

A similar bound is obtained by demanding that SN cooling not be modified in the presence of the above additional fields.

It is remarkable that this framework with a low NP scale, $f \lesssim \Lambda_{\text{EWB}}$ where Λ_{EWB} is the electroweak symmetry breaking (EWB) scale, cannot be excluded by direct experimental data. In many late neutrino mass models, there are degrees of freedom beyond ϕ . These additional degrees of freedom can yield indirect signals provided that standard cosmology is assumed. In the case in which neutrino masses are protected by global or approximate symmetries [12] or the case with strong dynamics (in which chiral symmetries are being broken by the condensate) [16], light pseudo-Goldstone bosons (PGB) field are typically present. Similarly, in models with gauge symmetries [14] the corresponding gauge boson masses are suppressed, relative to f , by an additional gauge coupling g , and therefore play a role similar to the one played by the PGBs. This additional light fields interact with the plasma through their coupling to the neutrinos. This happens even below the BBN phase transition and may leave a trace in the observed cosmic microwave background radiation (CMBR) [12].¹

The focus of this work is to investigate other more direct ways of testing the low scale models of neutrino masses. We find that such a possibility of a more direct probing of this

¹see also [19] for related analysis.

class of models, at present or in the not-too-distant future, does exist. The desired signal would consist of a dramatic modification of the supernova neutrino flux (diffuse or burst) through interaction between the these neutrinos and the cosmic background neutrinos (CBN). These interactions are mediated by the new scalar particles introduced by the NP.

In Section 2 we discuss the dominant processes which modify the incoming SRN fluxes. We divide this section into two: in 2.1 we describe the resonant process which happens only in a narrow range of parameter space, but leads to a spectacular signal through what we denote as accumulative resonance; in 2.2 we consider the non-resonant processes which, in conjunction with data from SN1987A, yield a lower bound on f comparable to the one from BBN. Also, at the beginning of Section 2 we summarize bounds on the model parameters imposed by BBN. These bounds are sketched in Fig.8, along with the parameter range for which the resonant process occurs. Finally, we conclude in Section 3.

2. Main idea, formalism and important processes

Our main idea in this work is to show that the presence of the additional light bosons, required in the late neutrino masses framework, can introduce a significant interaction between the SRN and the CBN. This, in some region of the model parameter space, can lead to a measurable modification of the incoming SRN flux. The typical SRN energies are above average solar neutrino energies and below the atmospheric ones. Consequently, this flux is likely to be observed by SK [20] and KamLAND [21] in the near future, or by successor experiments. Thus, there is a window (although not very wide) in which we can observe the presence of both this extra light degrees of freedom and the CBN! In this part we introduce the relevant part of the Lagrangian and discuss important processes which yield the signal.

We first discuss resonant processes and present the phenomenon of accumulative resonance which can yield a possible signal. Then we move to discuss non-resonant processes, which, through the observation of neutrinos from SN1987A yield a bound comparable with the BBN one. Other implications of these processes, related to experiments envisioned for the near future, are also discussed.

Below EWB scale and close to the neutrino flavor symmetry breaking scale the effective Lagrangian can be written as

$$\mathcal{L}_\nu^D = \mathcal{L}_{\text{kin}} + y_\nu \phi \nu N + V(\phi), \quad \mathcal{L}_\nu^M = \mathcal{L}_{\text{kin}} + y_\nu \phi \nu \nu + V(\phi), \quad (2.1)$$

where $\mathcal{L}_\nu^{D,M}$ stands for the Dirac and Majorana case respectively, \mathcal{L}_{kin} denotes the kinetic part (for the gauge case this contains interaction between ϕ and the additional gauge bosons [14]), ν represents an active neutrino, $V(\phi)$ is the scalar potential (for the global case this contains interaction between ϕ and the additional Goldstone bosons [14]), and flavor and spinor indices are suppressed for simplicity. The above implies that

$$m_\nu = y_\nu f. \quad (2.2)$$

As we shall show below our signal is similar in both the Dirac and Majorana cases. For simplicity through our discussion below we omit the effect of neutrino mixings (apart from the discussion related to the SRN flux).

In order to establish a reference point, we first pause to summarize the bounds on various models imposed by BBN constraints in terms of the Yukawa couplings y_ν . This will enable an evaluation of the feasibility of our program.

1) The minimal model is of Majorana neutrinos with Abelian symmetry. We assume that the symmetry breaking scale, f , is below the BBN temperature of about 1 MeV. Then during the BBN epoch we cannot separate the Goldstone and the scalar (higgs) as they are a single entity, a complex scalar field. The updated BBN bound on the number of neutrinos is $N = 3.24 \pm 1.2$ at 95% [22]. The complex scalar adds 8/7 (neutrino) degrees of freedom, so this additional degree of freedom can be accommodated with the BBN bound above. However, there are other cosmological bounds, such as SN cooling rate. If the BBN bound on relativistic degrees of freedom should decrease by a significant amount, then the yukawa y_ν would be subject to the upper bound obtained in the next paragraph.

2) In the non-Abelian Majorana models, typically several complex scalars are present, which are not permitted to be by BBN considerations. Thus, in this case y_ν must be bounded from above to ensure decoupling. This bound was derived in Ref. [18] and we have also calculated this bound (as a check) by considering all the processes that would produce G's. Recoupling via the $2 \rightarrow 1$ process $\nu\nu \rightarrow G$ takes place as the temperature falls to some value T_{rec} determined by equating the decay rate at T_{rec} to the Hubble expansion rate:

$$\frac{M_G}{3T_{\text{rec}}} \frac{y_\nu^2 M_G}{16\pi} = \sqrt{\frac{8\pi^5 g}{45}} \frac{T_{\text{rec}}^2}{M_{Pl}^2} , \tag{2.3}$$

where g is the number of degrees of freedom at T_{rec} . By requiring $T_{\text{rec}} < T_{\text{BBN}}$ we find

$$y_\nu \lesssim 6 \times 10^{-7} (\text{keV}/M_G) \tag{2.4}$$

3) Finally, for the Dirac case, the absence of a negligible population of right-handed (sterile) neutrinos (N) in the bath disallows the reaction $\nu N \rightarrow G$, so that G's can only be produced via $\nu_L \nu_L \rightarrow G G$ (via t channel N exchange). Requiring that this process be out of equilibrium at T_{BBN} yields a BBN bound of

$$y_\nu \lesssim 1 \times 10^{-5} . \tag{2.5}$$

The s -channel process requires a chirality flip which makes the bound weaker, as pointed out in Ref. [12]. Note that this bound is independent of the Goldstone mass.

2.1 Resonance, accumulative resonance

The simplest possible process which will modify the spectrum is the resonant production of one of the above light bosons in the collision of an SRN and a CBN. For simplicity we shall assume that the boson couples only to neutrinos with a strength y_ν (this is a good approximation for the case in which $y_\nu \lesssim 10^{-6}$ as discussed below). Thus after being produced the boson, say a scalar, will decay back yet to a pair of neutrinos, where in our frame they have an energy spectrum flatly distributed between 0 and the resonant lab energy E^* . The diagram describing this process is shown in Fig. 1. For clarity, we frame

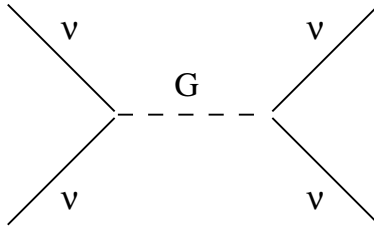


Figure 1: Diagram representing resonant scattering. G denotes a pseudo-Goldstone or a gauge boson.

the discussion which follows in terms of Dirac neutrinos. Except for differing dynamics in the high temperature environment of the supernova, the Majorana case is similar.

The resonant scattering affects the incoming SRN spectrum: the energy of the incoming SRN is now divided between the two decay products, so that we expect to observe a depletion in the expected spectrum for incoming neutrino with the appropriate energies. In addition, since this process is effectively $1 \rightarrow 2$, one may expect some depletion in the observed incoming neutrino spectrum at appropriate energies.

Two questions are in order:

1. In view of the small upper limit on the interaction strength (2.5) and the low density of the background neutrinos, will the resonance process indeed produce the depletion discussed in the preceding paragraph?

2. If the answer to (1) is affirmative, can the resultant depletion be observed by present or near future experiments?

Both these questions will be answered in the affirmative.

2.1.1 Resonance: no cosmological expansion

First, we consider the case of no expansion and begin by estimating the mean free path (m.f.p) λ_{Res} for the resonance process.

The cross section, written in Breit-Wigner form for the process in Fig. 1 is roughly given by

$$\sigma_{\text{Res}} \simeq \frac{y_\nu^4}{16\pi} \frac{s}{(M_G^2 - s)^2 + M_G^2 \Gamma_\nu^2}, \quad (2.6)$$

where G stands for the gauge/Goldstone bosons and s is the square of the center of mass energy. In addition Γ_ν is the decay width of the boson into neutrino pair,

$$\Gamma_\nu \sim \frac{y_\nu^2 M_G}{4\pi}. \quad (2.7)$$

Consider the case in which the SRN energy, E , is on resonance, E^* ,²

$$E^* \simeq \frac{M_G^2}{2m_\nu}, \quad (2.8)$$

²We shall neglect here the effect of thermal broadening, effectively assuming that the background neutrinos are at rest. We further discuss this point below.

so that

$$\sigma_{\text{Res}} \simeq \frac{\pi}{M_G^2}. \quad (2.9)$$

Consequently on the resonance³ the m.f.p is give by

$$\lambda_{\text{Res}} \approx \frac{1}{n_\nu^0 \sigma_{\text{Res}}} \sim \frac{M_G^2}{\pi n_\nu^0}, \quad (2.10)$$

where $n_\nu = 3\pi\Gamma(3)\zeta(3)T_\nu^3 \sim 56(1+z)^3 \text{ cm}^{-3}$, n_ν^0 is the present background neutrino density and T_ν is the background neutrino temperature with $T_\nu^0 \sim 1.6 \times 10^{-4} \text{ eV}$ being the present one. On carrying out the numerical substitution and using Eq. 2.8 we find⁴

$$\lambda_{\text{Res}} \approx \frac{1}{n_\nu^0 \sigma_{\text{Res}}} \sim \frac{2m_\nu E^*}{\pi n_\nu^0} \sim 5 \times 10^{-6} \text{ pc} \frac{m_\nu}{5 \times 10^{-2} \text{ eV}} \frac{E^*}{10 \text{ MeV}}. \quad (2.11)$$

Since this is much smaller than a typical distance traveled by a SRN neutrino, we find that the answer to Question 1. is positive for practically *any* value of y_ν . That is, if a neutrino is produced with the appropriate resonance energy then the process will go through.

However, this analysis also implies that the answer to the second question is negative. The extent of the dip in the SRN flux is controlled by the width of the boson G . This width is tiny (2.7), rendering it impossible at present or in the near future to detect such a narrow depletion in the SRN flux. The results are markedly different when cosmological expansion is included, so that we now turn to consider this case.

2.1.2 Resonance: case of cosmological expansion

We start by finding the conditions on the coupling for which there is sizable resonant degradation of of the *original* (not products of G -decay) flux of supernova neutrinos.

The the probability $P(E, z)$ that a neutrino, created at red shift z , with energy $(1+z)E$ arrives unscattered at the detector with energy E is given [23] by an integration over proper time (converted to an integral over intermediate red shifts \bar{z}):

$$P(E, z) = \exp \left[- \int_0^z \frac{d\bar{z}}{H(\bar{z})(1+\bar{z})} \bar{n}_\nu \sigma_{\nu\nu \rightarrow \phi}(2m_\nu(1+\bar{z})E) \right], \quad (2.12)$$

where σ is the resonant scattering cross section, a function only of $s = 2m_\nu(1+\bar{z})E$, the (c.m. energy)² at the time of scattering, \bar{n}_ν is the neutrino density (per flavor) at redshift \bar{z} and $H(\bar{z})$ is the Hubble constant at redshift \bar{z} .

For the purposes of this section, it is sufficient to employ a δ -function approximation for the cross section,

$$\sigma = \frac{\pi}{4} y_\nu^2 \delta(s - M_G^2), \quad (2.13)$$

³The result is qualitatively the same if one averages over the width of the resonance.

⁴Note that for the lighter neutrino the m.f.p is even somewhat shorter.

so that the integral in (2.12) is trivially done, with the result

$$P(E, z) = \exp \left[- \left(\frac{\pi y_\nu^2}{4M_G^2} \right) n_\nu^0 \left(\frac{E^*}{E} \right)^3 \left(\frac{1}{H(E^*/E)} \right) \right]$$

for $\frac{E^*}{(1+z)} < E < E^*$,

(2.14)

and $P = 1$ otherwise. Here E^* is the resonance energy $M_G^2/(2m_\nu)$. There will be large depletion of the initial SRN flux in this entire domain if

$$\frac{\pi y_\nu^2}{4M_G^2} \frac{n_\nu^0}{H_0} > 2 ,$$
(2.15)

which gives (for $H_0 = 70 \text{ km s}^{-1} \text{ Mpc}^{-1}$)

$$y_\nu \gtrsim 4 \times 10^{-8} \frac{M_G}{1 \text{ keV}} .$$
(2.16)

In the Dirac or Abelian Majorana cases, this permits a reasonable window of more than two orders of magnitude in y_ν (see Eq. (2.5)) for substantial depletion of the SRN flux. For the non-abelian Majorana case, the window narrows to a bit more than an order of magnitude (Eq. (2.4)). In this “strong coupling” regime we take $P = 0$ in the domain in Eq. (2.14). Thus, with the appropriate constraint on y_ν , Question 1. is answered in the affirmative: there will be substantial depletion of the original flux due to resonant scattering. We now proceed to see how the cosmological evolution permits a signal to be formulated for the resonant scattering.

2.1.3 Accumulative resonance

In the domain (2.14), there will be resonant absorption out of the original neutrino flux, but some replenishment as well, from neutrinos re-emitted in the decay of a G produced in the domain in Eq. (2.14). More specifically, suppose that a neutrino emitted with energy $\epsilon \geq E^*$ from a source at redshift z undergoes resonant scattering at redshift $\bar{z} < z$, so that

$$E^* = \epsilon \frac{1 + \bar{z}}{1 + z} .$$
(2.17)

This is followed by the emission of a decay neutrino with energy $E' = fE^*$, $0 \leq f \leq 1$ immediately following emission. The flatness of the emitted-neutrino spectrum implies that that f will vary uniformly over the region $[0,1]$.⁵ In that case the observed energy at the present era is

$$\begin{aligned} E &= \frac{fE^*}{1 + \bar{z}} \\ &= \frac{f\epsilon}{1 + z} \\ &= fE_{\text{unscattered}} \end{aligned}$$
(2.18)

⁵We thank K. Hikasa for drawing our attention for this issue.

where $E_{\text{unscattered}} \equiv \epsilon/(1+z)$ would be the observed energy of the neutrino in the absence of resonant scattering. This shows how the entire allowed region of energy below E^* is populated by rescattered neutrinos, with the energies shifted downward from the original spectrum. Especially interesting is the spectrum for $E_{\text{unscattered}}$ at or just below the limit E^* : in that case, the only replenishment of flux is from the tail of the decay distribution ($f \simeq 1$), from resonant production that has taken place only recently. (From Eq. (2.17), one can see that the condition for this is that $\bar{z} = 0$.) The restriction to $f \simeq 1$ implies very little replenishment, so that a dip at $E = E^*$ should be a universal feature of the final spectra observed. This is a qualitative response, in the affirmative, to Question 2: the spectrum with absorption will show a dip at $E = E^*$, and will be shifted downward from the spectrum absent resonant absorption. The complete effect of neutrinos emitted with non-resonant energies, passing through resonance, and then replenishing the flux at lower energies, is what we call accumulative resonance.

2.1.4 A note on thermal broadening

In the presence of the cosmological expansion, the effect of thermal broadening (because the CBN are not at rest) on our principal result (the universal dip described in the previous subsection) is negligible. The argument is as follows: after decoupling, the CBN spectrum is Fermi-Dirac, but in momentum rather than energy, even into the non-relativistic region. Thus the effect of thermal broadening is the introduction of a momentum spread in the target neutrinos of $\mathcal{O}(T_\nu)$. The principal feature of our result, the dip at neutrino energy E^* , occurs with neutrinos undergoing resonant scattering in the recent era, where the CBN is completely non-relativistic. In that case, the fractional energy shift of the target neutrinos is $\mathcal{O}(T_\nu^2/m_\nu^2) \sim 10^{-3}$ (unless the lightest neutrino has mass $\lesssim 10^{-4}$ eV). As can be seen from Eq.(2.12), the effect of this uncertainty is the introduction of a spread of $\mathcal{O}(10^{-3})$ in the value of \bar{z} , the red shift at resonant scattering. The consequence, after integrating over red shift of the source, is that the sharpness of the dip at E^* in the observed energy spectrum is softened by effects of $\mathcal{O}(10^{-3})$ rather than $\mathcal{O}(y_\nu^2) \sim 10^{-15}$, corresponding to the intrinsic width of the resonance. Thus, even with thermal broadening, the relative sharpness of the dip is preserved.

2.1.5 Event rates

Next we consider the effect of the accumulative resonance on the total SRN differential flux. We first present the standard expressions for the SRN flux. Then we shall discuss how to incorporate the accumulative resonance effect. The differential flux of Supernova Relic Neutrinos (SRN) is given by

$$\frac{dF}{dE} = \int_0^{z^{max}} R_{\text{SN}}(z) \left\langle \frac{dN(\epsilon)}{d\epsilon} \right\rangle_{\epsilon=(1+z)E} (1+z) \left| \frac{dt}{dz} \right| dz, \quad (2.19)$$

where for heuristic purposes we adopt as our standard the Fermi-Dirac distribution

$$\frac{dN(\epsilon)}{d\epsilon} = \mathcal{E} \times \frac{120}{7\pi^4} \times \frac{\epsilon^2}{(T_\nu^{\text{SN}})^4} \times \frac{1}{\exp\left(\frac{\epsilon}{T_\nu^{\text{SN}}}\right) + 1}. \quad (2.20)$$

The constant $\mathcal{E} = 0.5 \times 10^{53}$ ergs is the total energy carried by each flavor of neutrino. The temperature for the electron antineutrinos is $T_{\bar{\nu}_e}^{\text{SN}} = 5$ MeV and for the non-electron neutrinos and antineutrinos is $T_{\nu_x}^{\text{SN}} = 8$ MeV. However, a more general form

$$\frac{dN}{d\epsilon} = \frac{(1 + \alpha)^{1+\alpha} \mathcal{E}}{\Gamma(1 + \alpha) \bar{\epsilon}^2} \left(\frac{\epsilon}{\bar{\epsilon}}\right)^\alpha e^{-(1+\alpha)\epsilon/\bar{\epsilon}}, \quad (2.21)$$

has been proposed [24] which provides a good fit to simulated explosions [25] of high-mass progenitors. Here where $\bar{\epsilon}$ is the average antineutrino energy at the source and the values of the fitting parameters $\bar{\epsilon}$ and α for the $\bar{\nu}_e$ and ν_x spectra from three different groups [25, 26, 24] (designated as LL, TBP, and KRJ, respectively), summarized in Table 1 of Ref. [27]. In Fig. 3 we will show the spread in our results obtained from the different spectra.

Since we are considering the case of Dirac neutrinos, we take the CBN to consist of equal mixtures of left- and right-handed non-relativistic neutrinos and antineutrinos, with total number equal to two degrees of freedom in equilibrium during BBN. The resonant scattering will take place between right-handed SRN antineutrinos and CBN neutrinos, as well as left-handed SRN antineutrinos and CBN neutrinos. The problem becomes complex since the spectrum of SRN neutrinos and antineutrinos are different. Again, for illustrative purposes, we present results for the simplified case in which only SRN right-handed antineutrinos undergo resonant scattering, but from a CBN left-handed neutrino population given by n_0 , as defined after Eq. (2.12).

The Jacobian factor in Eq. (2.19) is given by

$$\frac{dt}{dz} = - \left[100 \frac{\text{km}}{\text{s Mpc}} h (1 + z) \sqrt{\Omega_M (1 + z)^3 + \Omega_\Lambda} \right]^{-1} \quad (2.22)$$

with $\Omega_M = 0.3$, $\Omega_\Lambda = 0.7$ and $h = 0.7$. The (comoving) rate of supernova formation R_{SN} is parameterized as

$$R_{\text{SN}}(z) = \left(\frac{0.013}{M_\odot} \right) \dot{\rho}_*(z) \quad (2.23)$$

where

$$\dot{\rho}_*(z) = (1 - 2) \times 10^{-2} M_\odot \text{yr}^{-1} \text{Mpc}^{-3} \times (1 + z)^\beta. \quad (2.24)$$

and the exponent changes at $z = 1$, from $\beta \sim 2 - 4$ (for $0 < z < 1$) to $\beta \sim 0$ (for $z > 1$) The uncertainty in parameters describing $R_{\text{SN}}(z)$ comes from the uncertainty in present knowledge of the Cosmic Star Formation Rate (CSFR) [28]. In this paper we choose “median” values for these parameters [28], $R_{\text{SN}}(0) = 2 \times 10^{-4} \text{yr}^{-1} \text{Mpc}^{-3}$, $\beta = 2$ (for $0 < z < 1$) and $\beta = 0$ (for $z > 1$).

The fact that the SN density is either constant (for far ones) or increasing with the distance (for near ones) is of great importance. It implies that most of the incoming neutrinos originate from distant SN (for far SN, the flux decreases like square of the distance while their density grows like the cube of the distance). Thus these neutrinos are redshifted and go through the resonance. The detection process is sensitive only to the flux of incoming anti-neutrino electrons. Consequently, to include contributions to the flux from the

muon/tau and electron neutrinos we use the corresponding temperature in the expression for the neutrino spectrum, and include a mixing factor which for the electron neutrinos is 0.69 and for the muon neutrinos is 0.31. As discussed in Refs. [29–31], the relation between the $\bar{\nu}_e$ spectrum observed on Earth to the various neutrino spectra at production depends critically on whether the neutrino mass hierarchy is normal or inverted. If normal ($m_3 > m_2 > m_1$), then strong matter effects cause the $\bar{\nu}_e$ at production to emerge from the stellar surface as the lightest eigenstate $\bar{\nu}_1$, with electron component $|U_{e1}|^2 \simeq 0.69$. The small mixing of the electron with the third eigenstate $|U_{e3}|^2 \ll 1$ allows an equivalent two-flavor picture, with the result that neutrinos produced in the supernova as $\bar{\nu}_x$, $x = \mu$ or τ will be received at Earth as $\bar{\nu}_e$ with probability 0.31, and with energies corresponding to the $\bar{\nu}_\mu/\bar{\nu}_\tau$ spectrum at production. For the case of the inverted hierarchy, $\bar{\nu}_e$'s produced in the supernova emerge as the lightest mass eigenstate, now $\bar{\nu}_3$. For $\sin^2 2\theta_{13} \lesssim 10^{-6}$ the resonance is non-adiabatic and there is complete conversion $\bar{\nu}_3 \rightarrow \bar{\nu}_1$. This case then is the same as for the normal hierarchy. The adiabatic case ($\sin^2 2\theta_{13} \gtrsim 10^{-4}$) is very different: the original $\bar{\nu}_e$'s remain as $\bar{\nu}_3$ when emerging from the stellar surface, contributing negligibly to the $\bar{\nu}_e$ flux at Earth. The entire $\bar{\nu}_e$ flux at Earth then corresponds to the original $\bar{\nu}_x$ produced in the supernova. For intermediate values of $\sin^2 2\theta_{13}$, the situation is of course more complicated. In this paper we will consider only the normal hierarchy. This can be regarded as conservative, since in some cases the $\bar{\nu}_x$ spectrum (which generates the $\bar{\nu}_e$ signal in the adiabatic inverted hierarchy scenario) is harder than the $\bar{\nu}_e$ spectrum, and will give rise to more events which escape the low energy cuts.

In this context, we also note that Earth matter effects have been shown to modify the observed fluxes and spectra on earth [32]. Since however the hierarchy in average energies between the $\bar{\nu}_e$ and the other flavor is milder than thought this effect is expected to be subdominant [33]. Thus for simplicity and since these effects are not expected to induce gross modification of the observed spectrum we neglected them altogether.

In order to obtain the observed spectrum, we note again that all neutrinos emitted from a source at redshift z with energies ϵ *outside* the window $E^* < \epsilon < (1+z)E^*$ will arrive at $z = 0$ without undergoing resonance, and with a flux given by Eq. (2.19) with energies $E > E^*$ and $E < E^*/(1+z)$. For source energies in the resonant absorption region, all of the original flux will undergo resonance absorption followed by decay into a flat spectrum. Since all energies, both before and after absorption, are redshifted at the same rate, one can obtain the rescattered spectrum by generating neutrino numbers according to Eq. (2.20) in energy slices $\Delta\epsilon$ at the source (for ϵ in the absorption region), and redistributing this number according to Eq. (2.18) uniformly over the observed energy region $0 < E < \epsilon/(1+z)$.

In Fig. 2 we show the resulting differential flux (with source flux given by Eq. (2.20), with and without the accumulative resonant effect, integrated over redshift up to $z = 4$, and for $M_G = 1.1$ keV. As discussed following Eq. (2.18), there is a sharp dip at $E = E^* \equiv M_G^2/2m_\nu$ for all values of z . To demonstrate the spread introduced through differing assumptions about the source flux, we show in Fig. 3 the flux with accumulative resonant effect for the three choices discussed in the context of Eq. (2.21).

To estimate the event rates for SK and GADZOOKS [34] (SK enriched by Gd) we

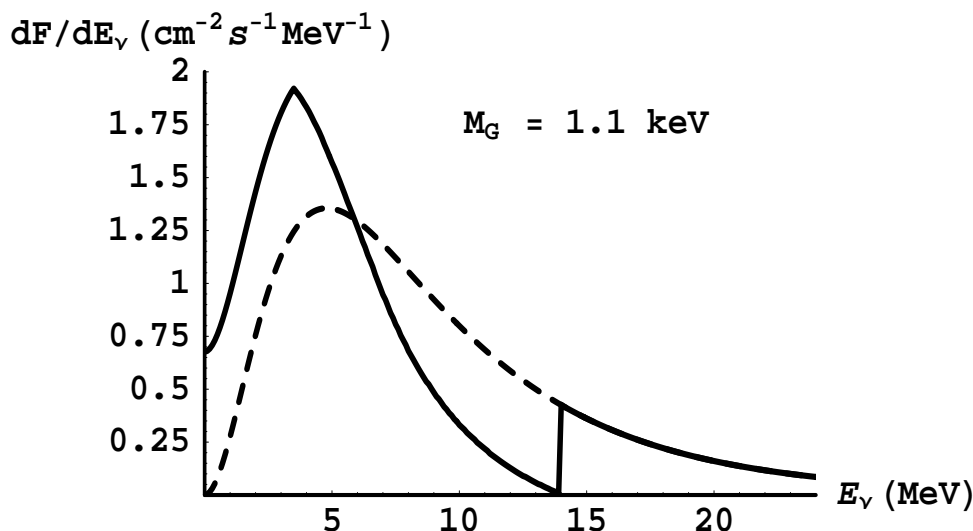


Figure 2: Depletion in the incoming SRN flux due to the resonance (solid curve) [35], compared to SRN flux without the resonance (dashed curve). The source flux is Fermi-Dirac.

show in Fig. 4 the differential neutrino flux folded with the detection cross section. This is for the inverse beta decay induced by the anti-neutrino capture in the detector [36]. (For calculating the event rate we have used the quasielastic neutrino-nucleon cross section given in Ref. [37].) The shape of the differential rate is modified due to the energy dependence of the cross section, which increases quadratically with energy. The main features of the effect due to the accumulation resonance such as the location of the dip and its width remain unmodified.

The differential rate for SK and Kamland is rather low regardless of the presence of the accumulative resonance. We therefore present in Figs. 5,6 the integrated flux for SK and KamLand, with and without accumulative resonance, respectively. Since the energy thresholds for SK and KamLand are 18 and 6 MeV respectively, we note the interesting feature that the event rate for SK is roughly unmodified (the 18 MeV threshold is well above the absorption region) while a suppression of roughly 25% is obtained for KamLand. This feature gains definition in Fig. 4. The GADZOOKS experiment has a much lower background (due to the ability of identifying the emitted neutron) and therefore may be able to provide a differential rate information.

2.2 Non-resonance

The effectiveness of the resonance process in redistributing the R_{SN} flux requires the resonance energy to be rather close to the peak energy $E_{\text{max}} \simeq 3T_{\nu}^{\text{SN}}$ of the R_{SN} (perhaps

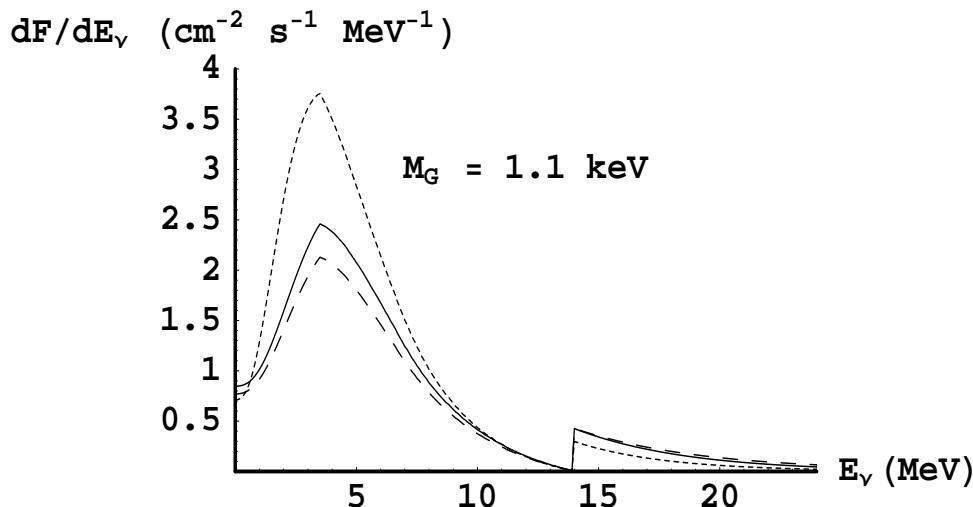


Figure 3: Depleted incoming SRN flux due to the resonance for three different assumptions about the source flux. Dotted, solid and dashed curves are fits designated as LL, TBP and KRJ, respectively, after eq. (2.21).

redshifted by a factor of 2) so that the boson mass must be in the range of 1 keV,

$$M_{\text{boson}} \simeq \sqrt{2m_\nu(1.5 T_\nu^{\text{SN}})} \sim 1 \text{ keV} \left(\frac{E}{10 \text{ MeV}} \right)^{1/2} \left(\frac{m_\nu}{0.05 \text{ eV}} \right)^{1/2}. \quad (2.25)$$

We note that the BBN constraint (1.1) implies that the mass of the symmetry breaking scalar $M_\phi \sim 10 \text{ keV}$ is typically above the resonance mass (not that far though). Thus it is more likely that the other light bosons (Goldstones or gauge) whose masses are almost unconstrained are required to provide a resonant channel for the R_{SN} scattering.

In view of this restriction, it is important to check whether other non-resonant processes can become important. One interesting possibility within the present dynamical framework is shown in Fig. 7. Again, this presumes the existence of either light Goldstone bosons [12, 17, 16, 18] or light gauge bosons [14]. For $s \ll M_\phi^2 \sim f^2$ we can estimate the cross section to be

$$\sigma_{\text{GG}} \sim \frac{y_\nu^2 f^2}{m_\phi^4} \sim \frac{y_\nu^2}{f^2} \sim \frac{y_\nu^4}{m_\nu^2}, \quad (2.26)$$

where we assume that the light bosons are produced on shell. This implies

$$M_G \leq \sqrt{2m_\nu E}/2. \quad (2.27)$$

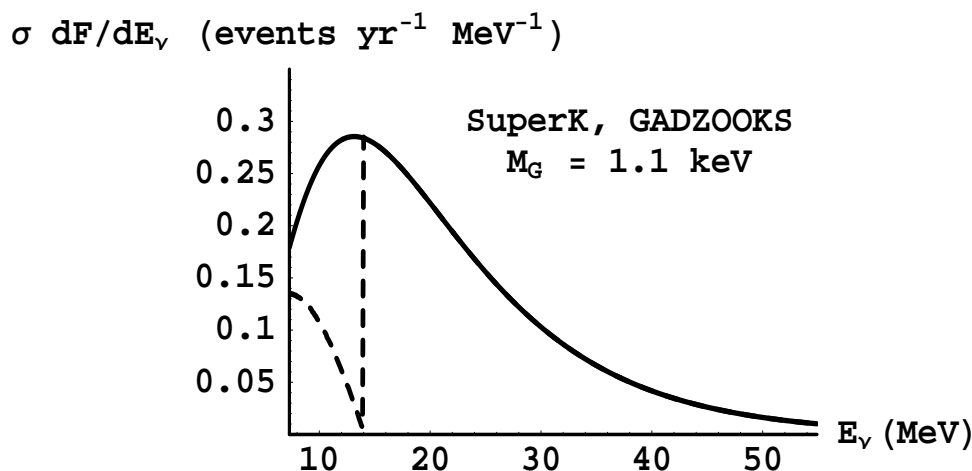


Figure 4: Depletion in the incoming SRN flux folded with the cross section for detection with SK and GADZOOKS (dashed curve) compared to no resonance case (solid curve). Source flux is Fermi-Dirac. The essential features of the accumulation resonance remain unmodified.

To have a substantial scattering we require the non-resonant mean free path $\lambda_{\text{non-res}}$ to be smaller than H^{-1} ,

$$\lambda_{\text{non-res}} H \simeq \frac{H m_\nu^2}{n_\nu y_\nu^4} \ll 1, \quad (2.28)$$

which yields a lower bound on y_ν

$$y_\nu > \left(\frac{H m_\nu^2}{n_\nu} \right)^{1/4} \sim 10^{-6} \left(\frac{m_\nu}{0.05 \text{ eV}} \right)^{1/2}. \quad (2.29)$$

This requirement is valid for any value of M_G .

The above process can have an effect on the SRN flux. If $M_G < 2m_\nu$ and there is sufficient optical depth, all the SRN will be transformed into invisible Goldstones and the signal is lost (for a related effect, see [38]). If $M_G > 2m_\nu$ then the process can effectively be characterized as $\nu \rightarrow 4\nu$, implying a substantial shifting of the entire SRN spectrum to lower energies. For a point source at distance ℓ , the condition analogous to Eq. (2.29) for sufficient optical depth is

$$y_\nu \geq 3.3 \times 10^{-6} \left(\frac{3000 \text{ Mpc}}{\ell} \right)^{1/4} \quad (2.30)$$

where l is the distance travelled by the SRN. For SN1987A, $\ell = 50,000$ pc, the fact that non-resonant scattering have not occurred, i.e. neutrinos with undegraded energy were

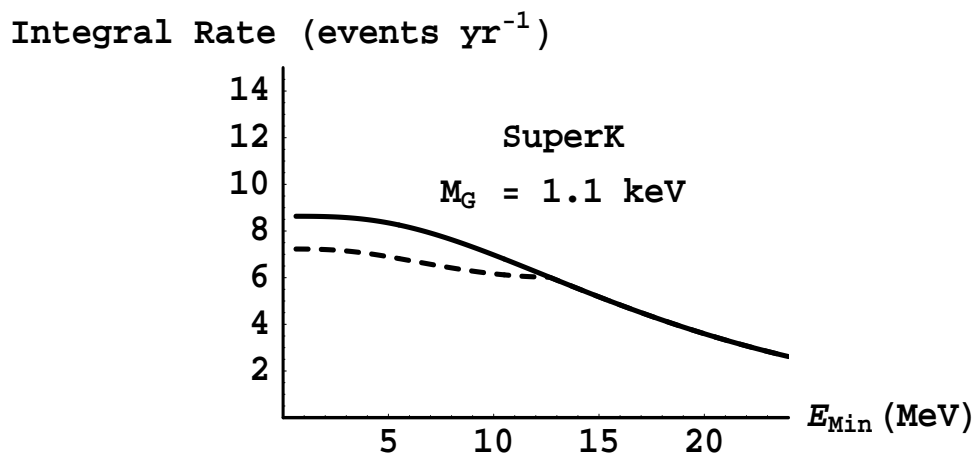


Figure 5: Integrated event rates for GADZOOKS and SK. The threshold for GADZOOKS is 10 MeV and for SK it is 18 MeV. Source flux is Fermi-Dirac.

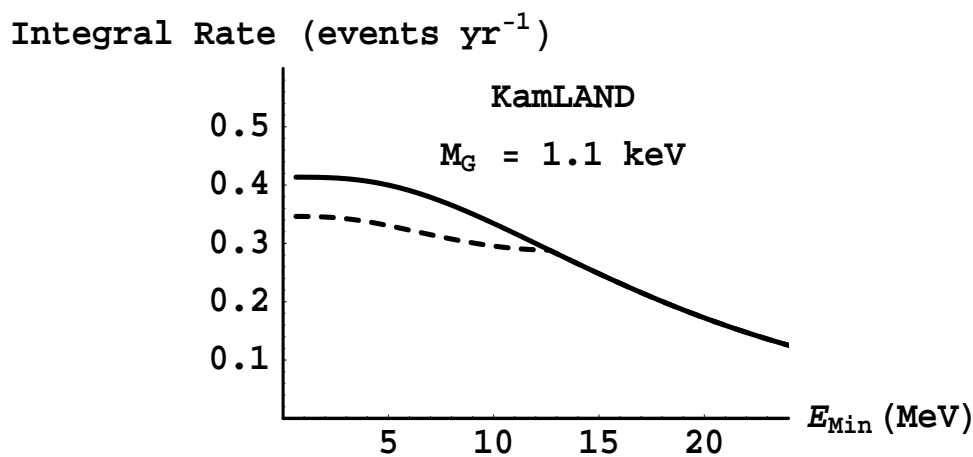


Figure 6: Integrated event rates for KamLand. The threshold for KamLand is 6 MeV.

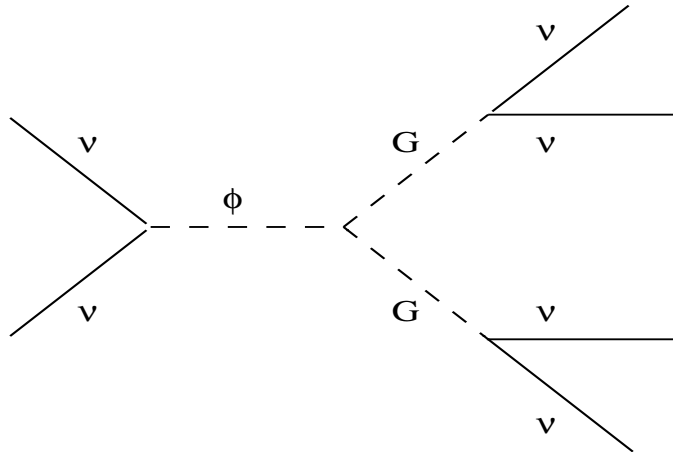


Figure 7: Two Goldstones production through Higgs exchange, off-resonance process.

observed [39], gives an independent upper bound on y_ν ,

$$y_\nu \lesssim 5.5 \times 10^{-5}. \quad (2.31)$$

This bound is comparable to the cosmological one in Eq.(1.1). However, considerably more detailed work is required in order to establish such a bound: a combined likelihood analysis in the symmetry breaking scale and the parameters describing the neutrino spectrum needs to be done in order to establish confidence levels for all variables [40], followed by marginalizing on the spectrum parameters. In the meantime, we adopt (2.31) as a rough, and preliminary, indication that this bound can be comparable to others we have mentioned.

3. Conclusions

In the models in which neutrino masses are light due to the low-energy symmetry breaking scale, extra light bosons are typically present. These light bosons couple to neutrinos with the coupling that is proportional to their masses and therefore directly related to the symmetry breaking scale. We have shown that, in principle, one can measure this coupling because SN neutrinos interact with cosmic background neutrinos via these bosons modifying the SN neutrino flux dramatically.

We have discussed two types of processes that are present due to these interactions. The first is a low energy analog of a Z-burst [41] where neutrinos interact producing an on-shell boson which subsequently decays to a pair of neutrinos. The expansion of the universe allows for a wide range of energies in which such a mini Z burst can occur. We characterize this process as accumulative resonance. The second is a non-resonance process which leads to a global degradation of energies in the supernova neutrinos flux. The observation of neutrinos from the 1987a supernova yields an important constraint on the parameter space of the above models since the observations were fully consistent with no such degradation.

As is often the case with observations related to neutrino physics, the signal we find is currently beyond the limit of each individual present experiment. The strength of our signal can improve once data from several experiments are combined. For example, our analysis reveals that a robust prediction of these models is that results from SK will be unaffected by the above processes while KamLand should observe a suppressed flux. However, in order to be convinced that depletion is observed it is desirable to actually observe the predicted dip in the flux, which requires certain amount of information on the energy dependence of the flux. This can be obtained in the future via water Cherenkov detectors enriched with added Gadolinium. (the GADZOOKS [34] proposal). This can be implemented in the very near future using an upgrade of SK, or attained via future experiments such as HyperK, UNO etc... Such and experiments can collect tens of relic supernova neutrinos per year and provide us with information on the energy dependence of SN flux. Furthermore they are more sensitive to observation of a single SN event. In principle these Megaton scale water Cherenkov detectors might detect neutrino bursts from O(Mpc)-distant SN with average rate of one burst per two years or so (see *e.g.* [42] and references therein). In this case there are two possible scenarios. The first, less exciting, is that the future burst of SN neutrinos will have differential flux consistent with the one observed from SN1987A. In this case we expect the bound on the model parameters to be improved. Alternatively, a more exciting possibility is that the incoming neutrinos are absent, or are severely degraded in energy relative to expectations. This can then be interpreted as a measurement of the coupling between the neutrino, the Higgs and the Goldstone bosons which yielded the degradation in the flux through the non-resonance process.

We note parenthetically that we have for simplicity neglected modification of the neutrino spectra due to shock wave effects [43]. These effects can change the spectral features, inducing non-adiabatic transitions, in a manner that depends on the neutrino flavor parameters [44]. However, since the exact dynamics related to the shock propagation in the SN is not well understood and since the above effects are red-shift dependent and will be smeared out once the integration over z is applied we shall not include them in our computations. In this context it is important to emphasize that the position of the universal dip will not vary with redshift.

It is also important to stress that in this work we focus on the modification of the dynamics of the SN neutrinos while they propagate in space outside the SN. Interesting effects may be induced by the presence of the new light degrees of freedom inside the SN [45]. These, however, are model dependent; for example, in the case of late Dirac neutrino masses the overall effect is expected to be miniscule [12]. The Majorana case is more involved since the extra bosons are expected to be thermalized inside the SN core [12, 45] through their reactions $\nu\nu \leftrightarrow G$. Consequently, the dynamics inside the SN might be significantly modified due to the presence of new flavor and lepton violating interactions. This requires a more detailed study which is beyond the scope of our paper. Note that it is likely that only the (pseudo) Goldstone boson will be light enough to be thermally produced inside the core [12, 13]. Thus the modification of the neutrino spectra [46] is expected to be below the 10% level expected just by counting degrees of freedom, which is probably smaller than other systematic effects which have not been included in this analysis.

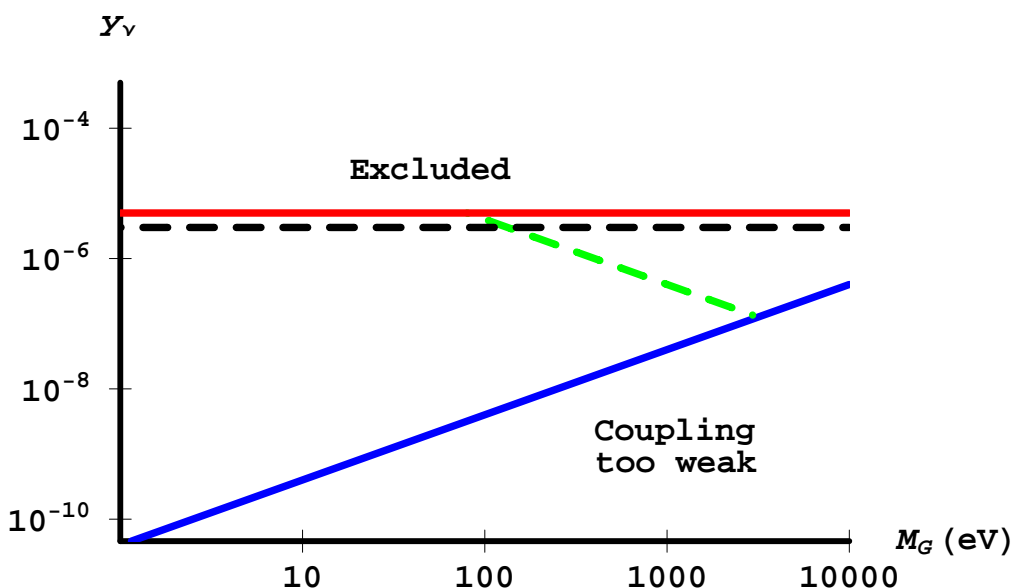


Figure 8: The cosmological bounds and the regions for the supernova neutrino spectrum distortion due to the resonance and non-resonance processes for a single Majorana (Dirac) neutrino for an abelian (non abelian) model are shown in (y_ν, M_G) plane. The region above the red (solid horizontal) line is excluded by the BBN constraint (for the Dirac case), SN cooling (for Majorana case) and due to the observation of (undegraded) SN1987A neutrinos. In the region below the blue (solid slanting) line the mean free path is too long for the resonance to occur. The region above the green (dashed slanting) line, which is relevant only for the non abelian Majorana case, is the region excluded by the BBN constraint. The region above the black (dashed horizontal) line is the region of the future experimental sensitivity to the observation/non-observation of the SRN neutrinos due to the non-resonant processes.

A summary of the available parameter space, subject to BBN bounds, for observation of the resonant process is given in Fig. 8. It is clear that BBN constraints are most constrictive in the case of Majorana neutrinos with a non-Abelian symmetry breaking sector. Least restrictive is the Dirac, Abelian, scenario; it also has the least impact on SN dynamics at the source.

We stress that observation of the above signals may not only shed light on the origin of neutrino masses but also yield an indirect observation of the elusive background relic neutrinos. In addition, we note that the above processes may be induced by other light degrees of freedom which couples to neutrinos. Consequently our signal may provide a test for other frameworks apart from the late neutrino masses one.

Acknowledgments

The authors thank the Aspen Center for Physics where this work was initiated. We wish to thank A. Burrows, G.D. Kribs, G. Steigman, H. Davoudiasl, H. Murayama, J.G. Learned,

K. Hikasa, L.J. Hall, R. Kitano, S. Oliver for helpful discussions. IS thanks J. Baker for the help with the numerical evaluations. This research was supported in part by the National Science Foundation Grant No. PHY99-0794. HG is supported by NSF under grant PHY-0244507, G.P. is supported by DOE under contract DE-AC03-76SF00098 and I.S. is supported by DOE under contracts DE-FG02-04ER41319 and DE-FG02-04ER41298 (Task C). I.S. would like to thank KITP Santa Barbara and LBNL Theory Group for hospitality while this work was being completed.

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