NEUTRINO MASS SPECTRUM FROM GRAVITATIONAL WAVES GENERATED BY DOUBLE NEUTRINO SPIN-FIP IN SUPERNOVAE

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ABSTRACT

The supernova (SN) neutronization phase produces mainly electron ($\nu_e$) neutrinos, the oscillations of which must take place within a few mean free paths of their resonance surface located nearby their neutrinosphere. The latest research on the SN dynamics suggests that a significant part of these $\nu_e$ can convert into right-handed neutrinos by virtue of the interaction of the electrons and the protons flowing with the SN outgoing plasma, whenever the Dirac neutrino magnetic moment is of strength $\mu_\nu < 10^{-11}\mu_B$, with $\mu_B$ being the Bohr magneton. In the SN envelope, some of these neutrinos can flip back to the left-handed flavors due to the interaction of the neutrino magnetic moment with the magnetic field in the SN expanding plasma (see the work by Kuznetsov & Mikheev; Kuznetsov, Mikheev, & Okrugin; Akhmedov & Kholopov; Itoh & Tsuneto; and Itoh et al.), a region where the field strength is currently accepted to be $B \approx 10^{13}$ G. This type of $\nu$ oscillation was shown to generate powerful gravitational wave (GW) bursts (see the work by Mosquera Cuesta; Mosquera Cuesta & Fiuza; and Loveridge). If such a double spin-flip mechanism does run into action inside the SN core, then the release of both the oscillation-produced $\nu_\mu$ and $\nu_\tau$ particles and the GW pulse generated by the coherent $\nu$ spin-flips provides a unique emission offset $\Delta T_{GW}^{\nu}$ for measuring the $\nu$ travel time to Earth. As massive $\nu$ particles get noticeably delayed on their journey to Earth with respect to the Einstein GW they generated during the reconversion transient, then the accurate measurement of this time-of-flight delay by SNEWS + LIGO, VIRGO, BBO, DECIGO, etc., might readily assess the absolute $\nu$ mass spectrum.

Subject headings: elementary particles — gravitational waves — methods: data analysis — neutrinos — stars: magnetic fields — supernovae: general

Online material: color figure

1. INTRODUCTION

The determination of the absolute values of neutrino masses is certainly one of the most difficult problems from the experimental point of view (Bilenky et al. 2003). One of the main difficulties of the issue of determining the $\nu$ masses from solar or atmospheric $\nu$ experiments concerns the ability of $\nu$ detectors to be sensitive to the species mass square difference instead of being sensitive to the $\nu$ mass itself. In this paper we introduce a model-independent novel nonparend method to achieve this goal. We argue that a highly accurate and largely improved assessment of the $\nu$ mass scale can be directly achieved by measurements of the delay in time of flight between the $\nu$ particles themselves and the gravitational wave (GW) burst generated by the asymmetric flux of neutrinos undergoing coherent (Pantaleone 1992) helicity (spin-flip) transitions during either the neutronization phase or the relaxation (diffusion) phase in the core of a Type II supernova (SN) explosion. Because special relativistic effects do preclude massive particles from traveling at the speed of light, while massless particles are not (the graviton in this case), the measurement of this $\nu$ time lag leads to a direct accounting of its mass. We posit from the start that two bursts of GWs can be generated during the proto–neutron star (PNS) neutronization phase through spin-flip oscillations: (1) one signal from the early conversion of active $\nu$ particles into right-handed partners, at density $\rho \sim$ few $\times 10^{12}$ g cm$^{-3}$, via the interaction of the Dirac neutrino magnetic moment [of strength $\mu_\nu < (0.7–1.5) \times 10^{-12}\mu_B$, with $\mu_B$ being the Bohr magneton] with the electrons and the protons in the SN outflowing plasma. Specifically, the neutrino chirality flip is caused by the scattering via the intermediate photon (plasmon) of the plasma electromagnetic current presented by electrons, $\nu_\mu e^- \rightarrow \nu_\mu e^-$; protons, $\nu_\mu p^+ \rightarrow \nu_\mu p^+$; etc. (2) A second signal exists by virtue of the reconversion process of these sterile $\nu$ particles back into actives some time later, at lower density, via the interaction of the neutrino magnetic moment with the magnetic field in the SN envelope (SNE). The GW characteristic amplitude, which depends directly on the luminosity and the mass square difference of the $\nu$ species partaking in the coherent transition (Pantaleone 1992), and the GW frequency of each of the bursts are computed. Finally, the time-of-flight delay $\nu \leftrightarrow$ GW that can be measured upon the arrival of both signals to Earth observatories is then estimated, and the prospective of obtaining the $\nu$ mass spectrum from such measurements is discussed.

2. DOUBLE RESONANT CONVERSION OF NEUTRINOS IN SUPERNOVAE

2.1. Interaction of $\nu_\mu$ Dirac Magnetic Moment with SN Virtual Plasmon

The neutrino chirality conversion process $\nu_\mu \leftrightarrow \nu_\tau$ in a SN has been investigated in many papers (see, for instance, Voloshin 1988; Peltoniemi 1992; Akhmedov et al. 1993; Dighe & Smirnov 2000). Next, we follow the reanalysis of the double $\nu$ spin-flip in SNe recently revisited by Kuznetsov & Mikheev (2007) and Kuznetsov et al. (2008), who obtained a more stringent limit on the neutrino magnetic moment, $\mu_\nu$, after demanding compatibility with the SN 1987A $\nu$ luminosity. The process becomes feasible in virtue of the interaction of the Dirac $\nu$ magnetic moment with a virtual plasmon, which can be produced, $\nu_\mu \rightarrow \nu_\tau + \gamma$, and absorbed, $\nu_\tau + \gamma \rightarrow \nu_\mu$, inside a SN. Our main goal here is to estimate the $\nu_\tau$ luminosity after the first resonant conversion inside the SN.
This quantity is one of the important parameters for estimating the GW amplitude of the signal generated at the transition (see § 3 below). The calculation of the spin-flip rate of creation of the $\nu_R$ in the SN core is given by (Kuznetsov & Mikheev 2007)

$$L_{\nu R} = \frac{dE_{\nu R}}{dt} = V \int_0^\infty \frac{dn_{\nu R}}{dE'} E' dE' = \frac{V}{2\pi^2} \int_0^\infty E'^3 \Gamma(E') dE',$$

where $dn_{\nu R}/dE'$ defines the number of right-handed $\nu$ particles emitted in the 1 MeV energy band of the $\nu$ energy spectrum, and per unit time, $\Gamma(E')$ defines the spectral density of the right-handed $\nu$ luminosity, and $V$ is the plasma volume. Thus, by using the SN core conditions that are currently admitted (see, for instance, Janka et al. 2007), plasma volume $V \approx 4 \times 10^{18}$ cm$^3$, temperature range $T = 30$–$60$ MeV, electron chemical potential range $\mu_e = 280$–$307$ MeV, neutrino chemical potential $\mu_\nu = 160$ MeV$^3$, one obtains

$$L_{\nu R} \approx \left(\frac{\mu_\nu}{\mu_B}\right)^2 (0.4–2) \times 10^{77} \text{ erg s}^{-1},$$

which for a $\mu_\nu = 3 \times 10^{-12} \mu_B$ compatible with SN 1987A neutrino observations and preserving causality with respect to the left-handed diffusion $\nu$ luminosity $L_{\nu L} < L_{\nu R} \lesssim 10^{53}$ erg s$^{-1}$, renders $L_{\nu R} \approx 4 \times 10^{53}$ erg s$^{-1}$. This constraint is on the order of the luminosities estimated in our earlier papers (Mosquera Cuesta 2000, 2002; Mosquera Cuesta & Fuji 2004) to compute the GW amplitude from $\nu$ flavor conversions, which were different from the one estimated by (Loveridge 2004). More remarkable, this analysis means that only $\sim 1\%$–$2\%$ of the total number of $\nu_L$ particles may resonantly convert into $\nu_R$ particles.

2.2. Conversion of $\nu_R \rightarrow \nu_L$ in the SN Magnetic Field

Kuznetsov et al. (2008) have shown that by taking into account the additional energy $C_L$, which the left-handed electron-type neutrino $\nu_L$ acquires in the medium, the equation of the helicity evolution can be written in the form (Voloshin & Vysotskii 1986; Voloshin et al. 1986a, 1986b; Okun 1986, 1988)

$$\frac{\partial}{\partial t} \left( \begin{array}{c} \nu_R \\ \nu_L \end{array} \right) = \left( \begin{array}{cc} \mathbf{E}_0 & \left( \begin{array}{c} \mu_\nu B_L \\ \mu_B \end{array} \right) \\ \mathbf{C}_L \left( \begin{array}{cc} \mu_\nu & \mu_B \\ \mu_B & \mu_B \end{array} \right) \end{array} \right) \left( \begin{array}{c} \nu_R \\ \nu_L \end{array} \right),$$

where the ratio $\rho/m_\nu = n_B$ is the nucleon density, $Y_e = n_e/n_B = n_p/n_B$, $Y_\alpha = n_{\alpha}/n_B$, and $n_{e,p,\alpha}$ are the densities of the electrons, protons, and neutrons, respectively, $B_L$ is the transverse component of the magnetic field with respect to the $\nu$ propagation direction, and the term $E_0$ is proportional to the unit matrix, however, it is not crucial for the analysis below.

As pointed out by Kuznetsov et al. (2008), the additional energy $C_L$ of left-handed $\nu$ particles deserves a special analysis. It is remarkable that the possibility exists for this value to be zero just in the region of the SNE we are interested in. And, in turn, this is the condition of the resonant transition $\nu_R \rightarrow \nu_L$. When the $\nu$ density in the SNE is low enough, one can neglect the value $Y_{\nu_L}$ in the term $C_L$, which gives the condition for the resonance in the form $Y_e = 1/3$. (Typical values of $Y_e$ in SNE are $Y_e \sim 0.4$–$5$, which are rather similar to those of the collapsing matter). However, the shock wave causes the nuclei dissociation and makes the SNE material more transparent to $\nu$ particles. This leads to the proliferation of matter deleptonization in this region and, consequently, to the so-called short $\nu$ outburst. According to the latest research on SNe, a typical gap appears along the radial distribution of the parameter $Y_e$ where it can achieve values as low as $Y_e \sim 0.1$ (see Mezzacappa et al. 2001 and also Fig. 2 in Kuznetsov et al. 2008, and references therein). Thus, a transition region unavoidably exists where $Y_e$ takes the value of $1/3$. It is remarkable that only one such point appears where the $Y_e$ radial gradient is positive, i.e., $dY_e/dr > 0$. Nonetheless, the condition $Y_e = 1/3$ is the necessary but yet not the sufficient one for the resonant conversion $\nu_R \rightarrow \nu_L$ to occur. It is also required to satisfy the so-called adiabatic condition. This means that the diagonal element $C_L$ in equation (3), at least, should not exceed the nondiagonal element $\mu_\nu B_L$, when the shift is made from the resonance point at the distance of the order of the oscillation length. This leads to the condition (Voloshin 1988)

$$C_L \leq \frac{\mu_\nu (dC_L/dr)}{\mu_B} \leq \frac{3G_\odot \rho}{\sqrt{2} m_\nu} \frac{dY_e}{dr} \leq \frac{3}{2} \frac{G_\odot \rho}{\sqrt{2} m_\nu} Y_e \frac{dY_e}{dr} \leq 10^{14} G \frac{dY_e}{dr} \leq 10^{14} G.$$

And values of these typical parameters inside the considered region are $dY_e/dr < 10^{-8}$ cm$^{-1}$ and $\rho \sim 10^{10}$ g cm$^{-3}$. Therefore, the magnetic field strength that realizes the resonance condition reads as

$$B_L \geq 2.6 \times 10^{14} G \frac{10^{-12} \mu_B}{\mu_\nu} \left(\frac{dY_e}{dr}\right)^{1/2} \left(\frac{10^{10} \text{ g cm}^{-3}}{10^{12} \text{ g cm}^{-3}}\right)^{1/2} \frac{dY_e}{dr} \left(\frac{10^{10} \text{ g cm}^{-3}}{10^{12} \text{ g cm}^{-3}}\right)^{1/2}.$$
where $D$ is the source distance, $L_\nu(t)$ is the total $\nu$ luminosity, $e_i \otimes e_i$ is the GW polarization tensor, the superscript $TT$ stands for the transverse-traceless part, and finally, $\alpha(t)$ is the instantaneous quadrupole anisotropy. Above, we estimated the $v_\perp$ luminosity; next, we estimate the degree of asymmetry of the PNS through the anisotropic parameter $\alpha$ and the timescale $\Delta T_{\nu L \rightarrow \nu_R}$ for the resonant transition to take place, as discussed above.

To estimate the star asymmetry, let us recall that the resonance condition for the transition $\nu_{\perp L} \rightarrow \nu_{\perp R}$ is given by (at the resonant $\vec{r}$)

$$V_{\nu}(\vec{r}) + B(\vec{r}) \cdot \hat{\vec{p}} - 2c_2 = 0. \quad (7)$$

Thus, the PNS magnetic field vector $B$ in equation (7) distorts the surface of resonance due to the relative orientation of $\vec{p}$ with respect to $B$ (see vector $B$ in Fig. 1). The deformed surface of resonance can be parameterized as $r(\beta) = \vec{r} + \rho \cos \beta$, where $\rho (\leq \vec{r})$ is the radial deformation and $\cos \beta = B \cdot \hat{\vec{p}}$. The deformation enforces a nonsymmetrical outgoing neutrino flux, i.e., the net flux of neutrinos emitted from the upper hemisphere is different from the one emitted from the lower hemisphere (see Fig. 1). Therefore, a geometrical definition of the quadrupole anisotropy can be $\alpha = (S_+ - S_-)/(S_+ + S_-)$, where $S_\pm$ is the area of the upper/lower hemisphere, whereas one obtains $\alpha \approx \rho /\vec{r}.^5$ The anisotropy of the outgoing neutrinos is also related to the energy flux $F_\nu$ emitted by the PNS and, in turn, to the fractional momentum asymmetry $\Delta |p|/|p|$ (Kusenko & Segré 1996; Barkovich et al. 2002; Lambiasa 2005a, 2005b; Mosquera Cuesta & Fiuza 2004). To compute $F_\nu$, one has to take into account the structure of the flux at the resonant surface, which acts as an effective emission surface, and the $\nu$ distribution in the diffusive approximation (Barkovich et al. 2002). As a result, one gets $\Delta |p|/|p| = \frac{1}{2} \int_0^\infty F_\nu \cdot u \, dS / \int_0^\infty F_\nu \cdot n \, dS \approx 2 \rho /\vec{r}^2 \hat{\vec{n}}$ ($\hat{\vec{n}}$ is a unit vector normal to the resonant surface and $\vec{u} = B /|B|$).^6 An anisotropy of \(\sim 1\)% would suffice to account for the observed pulsar kicks (Kusenko & Segré 1996; Loveridge 2004; Mosquera Cuesta 2000, 2002); hence, $\alpha \approx 0.045 \sim O(0.01) \sim O(0.1)$, which is consistent with numerical results (Burrows & Hayes 1996; Müller & Janka 1997). Finally, the conversion probability is $P_{\nu_L \rightarrow \nu_{\perp R}} = 1/2 \approx 1/2 \cos 2\theta \cos 2\theta$ (Okun 1986, 1988), where $\theta$ is defined as

$$\tan 2\theta = 2 \mu_\nu B_z / (B \cdot \hat{\vec{p}} + V_{\perp R} - 2c_2). \quad (8)$$

The quantities $\theta_L = \theta(\vec{r}_L)$ and $\theta_R = \theta(\vec{r}_R)$ are the values of the mixing angle at the initial point $\vec{r}_L$ and the final point $\vec{r}_R$ of the neutrino path.

Meanwhile, the average timescale of this first $\nu$ spin-flip conversion is (Dar 1987, Voloshin 1988)

$$\Delta T_{\nu L \rightarrow \nu_R} = \frac{(\mu_\nu \mu_\nu)}{2\pi c_2 (1 + (Z)) Y_e} \left( \frac{m_\nu}{p} \right), \quad (9)$$

where $(Z) \approx O(1-30)$ is the average electric charge of the nuclei and $\alpha_{\text{dec}}$ is the fine-structure constant. Using the current bounds on the neutrino magnetic moment $\mu_\nu \approx 3 \times 10^{-12} \mu_B$, $T_\nu \approx 1/3$, $(Z) \approx 1$, $\rho \approx 2 \times 10^{14} \text{ g cm}^{-3}$, and $\alpha \approx 0.04$, it follows that $\Delta T_{\nu L \rightarrow \nu_R} \approx (1 - 10) \times 10^{-2} \text{ s}$ (parameters have been chosen from SN simulations evolving the PNS on timescales of $\sim 3 \text{ ms}$ around core bounce; Mayle et al. 1987; Walker & Schramm 1987; Burrows & Hayes 1996; Mezzacappa et al. 2001; van Putten 2002; Arnaud et al. 2002; Beacom et al. 2001). In such a case, the above time-scale suggests that the GW burst would be as long as the expected duration of the pure neutronization phase itself, i.e., $\Delta T_{\text{neut}} \approx 10 - 100 \text{ ms}$, according to most SN analyses and models (Mayle et al. 1987; Walker & Schramm 1987; Burrows & Hayes 1996; Mezzacappa et al. 2001; van Putten 2002; Arnaud et al. 2002; Beacom et al. 2001), with the maximum GW emission taking place around $\Delta T_{\text{max}} \approx 3 \text{ ms}$ (van Putten 2002; Arnaud et al. 2002; Mosquera Cuesta 2000, 2002; Mosquera Cuesta & Fiuza 2004). Hence, the outcomeing GW signal will be the evolute (linear superposition) of all the coherent $\nu_{\perp L} \rightarrow \nu_{\perp R}$ oscillations taking place over the neutronization transient, in analogy with the GW signal with the collective motion of neutrion matter in a just-born pulsar. This implies a GW frequency of $f_{\text{GW}} \sim 1/\Delta T_{\text{neut}} \sim 100 \text{ Hz}$ for the overall GW emission, and $f_{\text{GW}} \sim 1/\Delta T_{\text{max}} \sim 330 \text{ Hz}$ at its peak. Meanwhile, according to our probability discussion above, about $1\%-2\%$ of the total $\nu$ particles released during the SN neutronization phase may oscillate (Voloshin 1988; Peltoniemi 1992; Akhmedov et al. 1993; Dighe & Smirnov 2000), carrying away an effective power $P_{\text{GW}} \approx 3 \times 10^{44} - 10^{45} \text{ erg s}^{-1}$, i.e., $0.01 \times 3 \times 10^{33} \text{ erg}$, emitted during $\Delta T_{\text{neut}} \approx 10 - 100 \text{ ms}$ (this is similar to the upper limit computed in Peltoniemi [1992], $L_{\nu} = (2/10) \times 10^{53} \mu_\nu / (10^{-12} \mu_B)$ erg s$^{-1}$). Moreover, as is evident from equation (6), the GW amplitude is a function of the helicity-changing $\nu$ luminosity, i.e., $h = h_{\text{Max}} \cos 2\theta$. The $\nu$ luminosity itself depends

\[5\] A detailed analysis of the asymmetry parameter $\alpha$ requires one to study its time evolution during the SN collapse. Such a task goes beyond the aim of this paper. Working in the stationary regime, we may assume $\alpha$ constant (see Burrows & Hayes 1996; Burrows et al. 1995; Zwerger & Müller 1997; van Putten 2002).

\[6\] To compute $\Delta |p|/|p|$ one uses the standard resonance condition $V_c = 2c_2$ (see Barkovich et al. 2002 for details). According to Mezzacappa et al. (2001), during the first $\sim 100 - 200 \text{ ms}$, $Y_e$ may assume values $\sim 0.3$ so that $F_\nu \sim (Y_e - 1)$ is suppressed by several orders of magnitude. At $\sim 10 \text{ ms}$, $\rho \approx 10^{12} \text{ g cm}^{-3}$, $r \approx 50 \text{ km}$, and $|p| \approx 10 \text{ MeV}$, the resonance condition leads to a range for $\Delta m^2 \cos 2\theta$ consistent with solar (or atmospheric) neutrino data.

\[7\] By using the typical values $\beta = 10^{12} \text{ G}$, $\mu_\nu \approx 9 \times 10^{-10} \mu_B$, and the profile $\rho \approx \rho_{\text{core}}(r)$ for $r \gg r_\text{c}$, $r_e \approx 10 \text{ km}$ is the core radius and $\rho_{\text{core}} \approx 10^{14} \text{ g cm}^{-3}$, one can easily verify that the adiabatic parameter $\gamma = 2(\mu_\nu B_z)/(8\pi c_2 |p|/|\vec{p}|) > 1$ at the resonance $\vec{r}$.\[8\]
on the probability of conversion (Peltoniemi 1992; Mosquera Cuesta 2000, 2002; Mosquera Cuesta & Fiuza 2004; Mosquera Cuesta & Lambiase 2004), i.e., $\Delta T_{GW} = \frac{P_{\nu_1 - \nu_2}}{D}$.

The characteristic GW strain $\gamma_1$ from the outgoing flux of spin-flipping (first transition) $\nu$ particles is

$$\gamma_1 = 1.1 \times 10^{-3} \left( \frac{H_{100}}{100 \text{ Hz}} \right)^{1/2} \frac{P_{\nu_1 - \nu_2}}{D} \left( \frac{10^3 \text{ erg s}^{-1}}{L_{GW}} \right),$$

for a SN exploding at a fiducial distance of 2.2 Mpc, e.g., at the Andromeda galaxy (see Table 1). The GW strain in this mechanism (see Fig. 2) is several orders of magnitude larger than in the SN $\nu$ diffusive escape (Burrows & Hayes 1996; Müller & Janka 1997; Arnaud et al. 2002; Loveridge 2004) because of the huge $\nu$ luminosity the $\nu$ oscillations provide by virtue of being a highly coherent process (Pantaleone 1992; Mosquera Cuesta 2000, 2002; Mosquera Cuesta & Fiuza 2004). This makes it detectable from very far distances. These GW signals are right in the bandwidth of the highest sensitivity ($10^{-3}$–300 Hz) of most ground-based interferometers.

Spin-flavor oscillations $\nu_{\mu L} \rightarrow \nu_{\mu R}$, which according to the latest research on SN dynamics do take place during the neutrinoization phase of core-collapse SNe (Ray et al. 1987; Walker & Schramm 1987; Voloshin 1988; Dighe & Smirnov 2000; Kuznetsov & Mikheev 2007), allow powerful GW bursts to be released from one side (according to the $\nu_{\mu L}$ dominant GW signal) and a stream of $\nu_{\mu R}$ particles to be generated from the other side, over a timescale given by equation (9). The latter would in principle escape from the PNS were it not for the appearance of several resonances that catch up with them (Voloshin 1988; Peltoniemi 1992; Akhmedov et al. 1993). If there were no such resonance, the $\nu_{\mu L} \rightarrow \nu_{\mu R}$ oscillation process would leak away all the binding energy of the star, leaving no energy at all for the left-handed $\nu_{\mu}$ particles that are said to drive the actual SN explosion and for us to have observed them during SN 1987A. A new resonance may occur at $\tilde{r} \approx 100$ km from the center, which converts $\sim 90\%$–$99\%$ of the spin-flip–produced $\nu_{\mu R}$ particles back into $\nu_{\mu L}$ ones (Voloshin 1988; Akhmedov 1988; Akhmedov & Khlopov 1988a, 1988b; Itoh & Tsuruta 1972; Itoh et al. 1996; Peltoniemi 1992; Akhmedov et al. 1993; Athar et al. 1995). As discussed in these papers, in fact, in the outer layer of the SN core, the amplitude of the coherent weak interaction of $\nu_{\mu L}$ with the PNS matter ($V_{\nu_L}$) can cross smoothly enough to ensure adiabatic resonant conversion of $\nu_{\mu R}$ into $\nu_{\mu L}$. Following Mezzacappa et al. (2001), the region where $V_{\nu_e} = 0$ and $Y_e = 1/3$ corresponds to a postbounce timescale $\sim 100$ ms and radius $\sim 150$ km at which the $\nu$ luminosity is $L_{\nu} \approx 3 \times 10^{52}$ erg s$^{-1}$, and the matter density is $\rho \sim 10^9$ g cm$^{-3}$. There, the adiabaticity condition demands $B_1 \approx 10^{10} G$ for the $\mu$ neutrino, quoted above (such a field is characteristic of young pulsars). This reverse transition (rt) should resonantly produce an important set of ordinary (muon and tau) $\nu_{\mu}$ and $\nu_{\tau}$ particles, which could be found far from their own neutrinosphere and, hence, could stream away from the PNS. Whence a second GW burst with the characteristics $h \approx 10^{-23}$ Hz$^{-1/2}$ for $D = 2.2$ Mpc and $\Delta T_{GW} \approx 1.4$ s is released in this region. Notice that this $h$ is similar to the one for the first transition despite the $\nu$ luminosity being lower. This feature makes it similar to the GW memory property of the $\nu$-driven signal, i.e., time-dependent strain amplitude with average value nearly constant (Burrows & Hayes 1996). To obtain this result, equations (9) and (10) were used. Wherefore, the GW frequency $f_{GW} \approx 1/\Delta T_{GW} \sim 0.7$ Hz falls in the low-frequency band and could be detected by the planned BBO and DECIGO GW interferometer observatories. Notice also that the time lag for the event at LIGO, VIRGO, etc., and the one at BBO and DECIGO is then about 100 ms. It is this transition which defines the offset to measure the time-of-flight delay, since both $\nu_{\mu L}$ and GW free-stream away from the PNS at this point.

4. TIME-OF-FLIGHT DELAY $\nu \leftrightarrow \nu$

The measurement of the $\nu \leftrightarrow \nu$ time delay from $\nu$ oscillations in SNe promises to be an innovative procedure to obtain the $\nu$ mass spectrum. Provided that Einstein’s GWs do propagate

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**TABLE 1**

| Source | Distance (kpc) | $\nu_{\mu L}$ | $\nu_{\mu R}$ | $\nu_{\tau}$ |
|--------|---------------|---------------|---------------|--------------|
| GC     | 10            | $5.15 \times 10^{-9}$ | $5.15 \times 10^{-3}$ | 0.32 |
| LMC    | 55            | $2.83 \times 10^{-8}$ | $2.83 \times 10^{-2}$ | 1.7 |
| M31    | $2.2 \times 10^3$ | $1.13 \times 10^{-6}$ | 1.13 | 68.8 |
| Source | $1.1 \times 10^4$ | $5.66 \times 10^{-6}$ | 5.66 | 344.0 |

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Note.—The $\nu$ masses in eV are $10^{-3}$, 1.0, and 2.5, for flavors $\nu_{\mu L}$, $\nu_{\mu R}$, and $\nu_{\tau}$, respectively.
at the speed of light, the GW burst produced by spin-flip oscillations during the neutronization phase will arrive at GW observatories earlier than its source (the massive \( \nu \) particles from the second conversion) will get to \( \nu \) telescopes.

As pointed out above, the mechanism to generate GWs at the instant at which the second transition \( \nu_\mu \rightarrow \nu_\tau \) takes place can by itself define a unique emission offset, \( \Delta T_{GW-\nu}^{\text{em}} = 0 \), which makes possible a cleaner and highly accurate determination of the \( \nu \) mass spectrum by “following” the GW and neutrino propagation to Earth observatories. The time lag in arrival is (Beacom et al. 2001)

\[
\Delta T_{GW-\nu}^{\text{isr}} \approx 0.12 \text{s} \left( \frac{D}{2.2 \text{ Mpc}} \right) \left( \frac{m_\nu}{0.2 \text{ eV}} \right)^2 \left( \frac{|p|}{10 \text{ MeV}} \right)^2. \tag{11}
\]

5. DISCUSSION

In most SN models (Burrows & Hayes 1996; Mezzacappa et al. 2001; Beacom et al. 2001; van Putten 2002), the neutronization burst is a well-characterized process of intrinsic duration \( \Delta T \approx 10 \text{ ms} \), with its maximum occurring within \( 3.5 \pm 0.5 \text{ ms} \) after core collapse (Mayle et al. 1987; Walker & Schramm 1987; van Putten 2002; Burrows & Hayes 1996). This timescale relates to the detectors’ approximate sensitivity to \( \nu \) masses beyond the mass limit

\[
m_\nu > 6.7 \times 10^{-2} \text{ eV} \left( \frac{2.2 \text{ Mpc}}{D} \right) \left( \frac{\Delta T}{10 \text{ ms}} \right)^{1/2} \left( \frac{|p|}{10 \text{ MeV}} \right). \tag{12}
\]

This threshold is in agreement with the current bounds on \( \nu \) masses (Fukuda et al. 1998).

Nearby SNe will somehow be seen. Apart from GWs and neutrinos, \( \gamma \)-rays, X-rays, visible, infrared, or radio signals will be detected. Therefore, their position on the sky and distance \( (D) \) may be determined quite accurately, including—if far from the Milky Way—their host galaxy (Ando et al. 2005). Besides, the Universal Time of arrival of the GW burst to three or more gravitational radiation interferometric observatories or resonant detectors will be precisely established (Schutz 1986; Arnaud et al. 2002). The uncertainty in the GW timing depends on the signal-to-noise ratio \( \sqrt{S/N} \) as \( \Delta T_{GW}^{\text{det}} = 1.45 \text{ s}/\sqrt{S/N} \sim 0.15 \text{ ms} \), with \( \tau \sim 1 \text{ ms} \) being the rms width of the main GW peak (Arnaud et al. 2002). Meanwhile, the type of \( \nu \) and its energy and Universal Time of arrival to \( \nu \) telescopes of the SNESW network will be highly accurately measured (Antonioli et al. 2004; Beaum & Vogel 1999). The \( \nu \) timing uncertainty is \( \Delta T^{\text{max}}_{GW-\nu} = \sigma_{\text{flash}}(N_\nu)^{-1/2} \), with \( \sigma_{\text{flash}} \sim (2.3 \pm 0.3) \text{ ms} \) and \( N_\nu \) being the event statistics (proportional to \( D \)). This leads to the SN distance-dependent uncertainty in the \( \nu \) mass, \( \Delta m_\nu^2 \propto \Delta T^{\text{max}}_{GW-\nu}/D \sim 0.5 - 0.6 \text{ eV}^2 \) (Arnaud et al. 2002), which implies \( m_\nu \sim 7 \times 10^{-1} \text{ eV} \), which is consistent with our previous estimate from equation (12). Hence, those \( \nu \) particles and their spin-flip conversion signals must be detected.

Therefore, the left-hand side of equation (11), i.e., the time-of-flight delay \( \Delta T_{GW-\nu} \), will be measured with a very high accuracy. With these quantities, a very precise and stringent assessment of the absolute \( \nu \) mass eigenstate spectrum will be readily set out by means not explored earlier in astroparticle physics: an innovative technique involving not only particle but also GW astronomy. For instance, at a 10 kpc distance, e.g., to the Galactic center (GC in Fig. 2), the resulting time delay should approximate \( \Delta T_{GW-\nu} \approx 5.2 \times 10^{-3} \text{ s} \), for a flavor of mass \( m_\nu \leq 1 \text{ eV} \) and \( |p| \sim 10 \text{ MeV} \). A SN event from the GC or Large Magellanic Cloud (LMC) would provide enough statistics in SNO, SK, etc., \( \sim 5000 - 8000 \) events, so as to allow for the definition of the \( \nu \) mass eigenstates (Beacom et al. 2001). Further out, \( \nu \) events are less promising in this perspective, but we stress that one \( \nu \) event collected by the planned megaton \( \nu \) detector, from a large-distance source, may prove sufficient (see further arguments in Ando et al. 2005).

6. SUMMARY

In this paper, it has been emphasized that knowing the \( \nu \) absolute mass scale with enough accuracy would turn out to be a fundamental test of the physics beyond the standard model of fundamental interactions. By virtue of the very important two-step mechanism of \( \nu \) spin-flavor conversions in SNe, very recently revisited by Kuznetsov et al. (2008), we suggest that by combining the detection of the GW signals generated by those oscillations and the \( \nu \) signals collected by SNESW from the same SN event, one might conclusively extract the \( \nu \) mass spectrum. In particular, sorting out the neutronization phase signal from both the \( \nu \) light curve and the second peak in the GW waveform (with its memorylike feature; Burrows & Hayes 1996) might allow one to achieve this goal in a nonpærel fashion.

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