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BURSTS OF GRAVITATIONAL WAVES DRIVEN BY NEUTRINO OSCILLATIONS

Herman J. Mosquera Cuesta
[1]Abdus Salam International Centre for Theoretical Physics, Strada Costiera 11, Miramare 34014, Trieste, Italy
[2]Centro Brasileiro de Pesquisas Físicas, Laboratório de Cosmologia e Física Experimental de Altas Energias
Rua Dr. Xavier Sigaud 150, CEP 22290-180 URCA, Rio de Janeiro, RJ, Brazil

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ABSTRACT

Active-to-sterile neutrino oscillations and non-spherical distortion of the resonance surface may trigger asymmetric emission of sterile neutrinos during the core bounce of a supernova collapse. The huge binding energy released and the proto-neutron star rapid rotation may power bursts of gravitational waves by the time the neutrino flavor conversions ensue. These bursts would be detectable upto distances of $\sim$ 55kpc. The relativistic requirement of an ellipsoidal axisymmetric core at maximum GWs emission may be explained by the neutrino-sphere geometry when the oscillations onset. Thus, neutrino oscillations might naturally be the underlying mechanism to produce either asymmetric structures, neutron star kicks and bi-polar jet ejecta in supernovae.

Subject headings: gravitation — relativity: theory — elementary particles — stars: neutron

1. ASTROPHYSICAL MOTIVATION

Core-collapse supernovae explosions are one of the most powerful sources of $\sim 10^{58}$ neutrinos of all species ([10-25]MeV): ν_e, ν_μ, ν_τ and probably ν_s and its antiparticles. $\Delta E_{total} \sim 5.2 \times 10^{53}\mathrm{erg}\left(\frac{10\mathrm{km}}{\mathrm{R_{NS}}}\right)\left(\frac{\mathrm{M_{NS}}}{1.4\mathrm{M}_\odot}\right)^2$ is released ($\Delta t \sim 12$s) from the proto-neutron star (PNS) binding energy (Burrows, Hayes & Fryxell 1995; Woosley & Weaver 1995; Janka & Müller 1996; Müller & Janka 1997). Numerical simulations have shown that high density gradients in the PNS appear as the implosion develops, creating anisotropic mass distribution regions where the matter density is too high for ν-trapping to occur (see Woosley & Weaver 1995; Janka & Müller 1996; Müller & Janka 1997), while in others they stream away freely. The time variation of density gradients suggest the ν-fluid energy-momentum tensor is somehow quadrupolar, whose evolution should induce emission of GWs (Burrows & Hayes 1995). This asymmetric ν-flow provides the conditions for GWs to be generated.

This *letter* suggests, for the first time, a new fundamental source of gravitational waves (GWs): the resonant conversion (Mikheyev & Smirnov 1985; Wolfenstein 1979) of active to sterile neutrinos in a *supernova core bounce*. The neutrino temperature asymmetry produced during $\nu_{\bar{\tau},\bar{\mu}} \longleftrightarrow \nu_s$ flavor changes and core rapid rotation couples to a non-spherical mass-energy distribution (the oscillating ν-fluid) inside the PNS during a core-collapse supernova explosion (SNE).

Generation of GWs bursts through active to sterile neutrino conversions, $\nu_{\bar{\tau},\bar{\mu}} \longleftrightarrow \nu_s$, are possible and powered by the onset of the inner core collapse and bounce of massive star explosions. Due to the dynamical conditions at core bounce we may expect the ν-oscillation process itself and its associated GWs burst to occur during the very early stages of evolution of the PNS: $\sim$ 1ms after core bounce ($\rho > 2 \times 10^{14}\mathrm{gcm}^{-3}$), i. e., when the sound wave produced when the PNS matter bounces at supranuclear densities becomes itself a shock wave (see Figure 6. in Burrows, Hayes & Fryxell 1995. Also Janka & Müller 1996; Müller & Janka 1997; Burrows & Hayes 1995; Epstein 1978, and references therein).

The issue of GWs pulses during SNEs has been subject of theoretical and extensive numerical follow-up (Müller & Janka 1997). Current theories for the generation of GWs signals during SNe core-collapse (type-II events) assume that the explosion itself is driven (revived) by a ν-flash released on a timescale $\sim$ 12s following the core bounce Müller & Janka 1997). The huge energy and momentum deposited at the base of the star mantle (Woosley & Weaver 1995; Janka & Müller 1996) finally succeed in ejecting the large (massive) envelope of the exploding red super-giant star. In all these models the GWs burst is produced due to the mass-tensor quadrupolar distorsion of the hunged-up core of the exploding supermassive star. The space-time transverse-traceless dimensionless metric strain is computed using the quadrupole formula $h^{tt}_{ij} = \frac{2G}{c^4 D}\frac{d^2 Q_{ij}}{dt^2}$, where Q_{ij} defines the mass quadrupole tensor and D corresponds to the source distance. The overall characteristics of the gravitational radiation emitted were estimated to have amplitudes of $10^{-(21-22)}$ for distances as far as the VIRGO cluster of galaxies $\sim$ 20Mpc (Woosley & Weaver 1995; Janka & Müller 1996), and frequencies around 1kHz.

However, since the binding energy of the exploding star is so powerful, we can expect the neutrino flux itself to be a strong source of GWs. This idea has been pursued in the late seventies by Epstein (1978), and more recently by Burrows & Hayes (1995). In the former study the amplitude of the GWs signal, produced when the outgoing (actives) ν-flux is absorbed and reemitted, was given as

$$|\theta^{ij(\nu)}_{TT}| \sim 2.0 \times 10^{-20} e^2 \left(\frac{10\mathrm{kpc}}{D}\right), \tag{1}$$

for a source placed at the distance of the galactic center.

Here e defines the eccentricity of the ellipsoidally distorted equatorial plane through which the neutrino flux escapes away (Epstein 1978).

In the second approach, basically a numerical computation one (Burrows & Hayes 1995), the enhanced asymmetric density perturbation provokes asymmetric shock propagation and breakout. This induces an asymmetric explosion and likely a ν-kick to the PNS. The GWs burst characteristics are estimated using Newtonian gravity together with the quadrupole approximation to general relativity, which yields

$$h_{ij}^{tt} = \frac{4G}{c^4 D}\int_{-\infty}^{t} \alpha(t')L_\nu(t')dt'. \quad (2)$$

Here $0.2 \leq \alpha(t) \leq 0.8$ is the instantaneous quadrupole anisotropy, and $L_\nu(t)$ the overall luminosity of neutrinos. As we show below, this expression is the right one to compute the GWs amplitude in the neutrino-oscillation process here introduced.

2. NEUTRINO OSCILLATIONS IN DENSE MATTER

We would like to examine the effect of possible neutrino conversions (Mikheyev & Smirnov 1985; Wolfenstein 1979) inside the core of a supernova on the gravitational waves emanating from the star. The amplitude of the GWs produced during the cooling phase (when most of the neutrinos escape by diffusion in the conventional mechanism) is nearly at the limit for detectability of LIGO (Burrows & Hayes 1995). We are interested in the possible contribution to the GWs amplitude by the neutrino flavor conversion within the first few milliseconds after the core collapse.

In order to produce an effect, neutrinos must be able to escape the core without thermalizing with the stellar material. For active neutrino species of energies $\approx$ 10MeV, this is not possible as long as the matter density is $\gtrsim 10^{10}$gcm^{-3}. Since the production rate of neutrinos is a steeply increasing function of matter density (production rate $\propto \rho^n$, where ρ is the matter density and $n > 1$), the overwhelming majority of the neutrinos of all species produced are trapped. So the contribution to the GWs amplitude is negligible, irrespective of the neutrino conversions taking place within the active neutrino flavors.

Sterile neutrinos, on the other hand, would be able to escape the core. Though they are not directly produced inside the star, if any active neutrino species can be copiously converted into sterile neutrinos through oscillations, it may be possible to dramatically increase the number of escaping neutrinos. This effect can be significant only if these active $\leftrightarrow$ sterile transitions take place inside the neutrinospheres of the active neutrinos, i.e. at $\rho \gtrsim 10^{10}$gcm^{-3}.

Let us first consider the case of vacuum oscillations, assuming for simplicity that the oscillations take place between only two neutrino species (ν_a and ν_s, say). For a ν_a state produced at $\vec{x}_0$ which does not undergo any collisions before reaching $\vec{x}$, the probability that it wil be observed at $\vec{x}$ as a ν_s is

$$P_{as}(|\vec{x}-\vec{x}_0|) = \sin^2 2\theta \sin^2\left(\frac{\Delta m^2}{4E_\nu}|\vec{x}-\vec{x}_0|\right), \quad (3)$$

where θ is the $\nu_a \leftrightarrow \nu_s$ vacuum mixing angle, Δm^2 is the mass-squared difference between the two neutrino mass eigenstates in vacuum, and E_ν is the neutrino energy. In order for the estimates, let us consider only the neutrinos travelling radially outwards, so that the problem is reduced to a one dimensional problem, with x and x_0 now representing the distance from the centre of the star. The probability that this neutrino would interact within a distance dx is given by

$$dP(x,x_0) = (1 - P_{as}(x-x_0))\mathrm{d}x/\lambda(x), \quad (4)$$

where $\lambda(x)$ is the mean free path of ν_a, given by $\lambda(x) \equiv [N(x)\sigma(x)]^{-1}$. Here $N(x)$ is the number density of the relevant scatterers at x while $\sigma(x)$ is the cross section for the neutrino scattering at x. From Eqs.(3,4), the probability of a ν_a reaching x from x_0 without interacting is (Mohapatra & Pal 1998)

$$\begin{aligned} P_{surv}(x,x_0) &= \exp\left[-\int_{x_0}^{x}\frac{\mathrm{d}x}{\lambda(x)}\right. \quad (5)\\ &+ \left.\sin^2 2\theta\int_{x_0}^{x}\sin^2\left(\frac{\Delta m^2}{4E_\nu}(x-x_0)\right)\frac{\mathrm{d}x}{\lambda(x)}\right]. \end{aligned}$$

The first integral in the exponential is the answer one gets in the absence of oscillations. It is this term that determines the size of the neutrinosphere for the active neutrino. The second term represents the enhancement of the survival probability due to oscillations.

The effect of the second term is negligible if the vacuum mixing angle is small ($\sin^2 2\theta \approx 0$) or when the oscillation length is large compared to the radius of the neutrinosphere ($4E_\nu/\Delta m^2 \gg R_\nu$). The latter condition is satisfied for $E_\nu \approx$ 10MeV only for $\Delta m^2 \lesssim 10^{-3}$eV^2. For Δm^2 larger than this value, the effect of the second term is an effective reduction in the radius of the neutrinosphere by a factor ζ, which can be estimated by taking $N(x) \propto x^{-\alpha} \longrightarrow \zeta \approx C^{\frac{1}{\alpha-1}}$, where

$$C \equiv 1 - \sin^2 2\theta \int_0^\infty \sin^2\left(\frac{\Delta m^2}{4E_\nu}x\right)\frac{\mathrm{d}x}{\lambda(x+R_\nu)} \quad . \quad (6)$$

The enhancement is maximum when C is smallest, which happens with maximal vacuum mixing angle ($\sin^2 2\theta = 1$), and with the maximum value of the integral in Eq.(6).

Pure vacuum oscillations would take place inside the star only when $\Delta m^2/(2E) \gg \sqrt{2}G_F\rho/m_N$, where G_F is the Fermi constant and m_N is the nucleon mass. Since $\rho \gtrsim 10^{10}$ gcm^{-3} inside the neutrinosphere, this condition is satisfied for $E \sim$ 10MeV neutrinos only for $\Delta m^2 \gg 10^4$eV^2. In this parameter range, the integrand is rapidly oscillating and the integral reduces to 0.5. Even in the most favorable scenario, thus, $C > 0.5$. Then, with $\alpha \approx 3$, we get $\zeta > \sqrt{2}$. Hence, the radius of the effective neutrinosphere cannot change by a large factor, which indicates that the oscillations into sterile neutrinos cannot increase the number of escaping neutrinos dramatically in this parameter range (where matter effects can be neglected).

The matter effects may help in allowing more neutrinos to escape, if resonant neutrino conversions into sterile neutrinos occur inside the neutrinosphere of the active neutrinos. In the case of $\nu_e \leftrightarrow \nu_s$ oscillations, the resonance occurs if

$$\sqrt{2}G_F\left(N_e(x)-\frac{1}{2}N_n(x)\right)\equiv A(x)=\frac{\Delta m^2}{2E_\nu}\cos 2\theta. \quad (7)$$

Here $N_e(x)$ is the electron number density (given by $N_{e^-}-N_{e^+}$) while $N_n(x)$ is the neutron number density. In the case of $\nu_{\mu,\tau} \leftrightarrow \nu_s$ the N_e term is absent, while in the case of antineutrinos, the potential changes by an overall sign. Numerically, for $\nu_e \leftrightarrow \nu_s$ oscillations,

$$A(x)=7.5\times 10^2\left(\frac{\mathrm{eV}^2}{\mathrm{MeV}}\right)\left(\frac{\rho_m(x)}{10^{10}\mathrm{g/cm}^3}\right)\times\left(\frac{3Y_e}{2}-\frac{1}{2}\right) \quad (8)$$

where Y_e is the electron number fraction. For $\nu_{\mu,\tau} \leftrightarrow \nu_s$ oscillations, the last term in parenthesis is changed to $(\frac{Y_e}{2}-\frac{1}{2})$. For all order of magnitude estimates we perform henceforth, we take the last term in the parenthesis to be of order one.

The neutrino conversions in the resonance region can be strong if the adiabaticity condition is fulfilled: the oscillation probability is $P_{as}=\cos^2\theta$, which is close to 1 in the case of small mixing angles. Moreover, after the resonance region, the newly created sterile neutrinos have very a small probability ($P_{sa}^{\mathrm{average}}=(1/2)\sin^2 2\theta$) of oscillating back to active neutrinos, which could be potentially trapped.

In order that the resonance condition is satisfied, we require $10^4\mathrm{eV}^2 \lesssim \Delta m^2\cos 2\theta\left(\frac{10\mathrm{MeV}}{E_\nu}\right) \lesssim 10^8\mathrm{eV}^2$, while the adiabaticity condition is satisfied for

$$\frac{\Delta m^2 \sin^2 2\theta}{2E_\nu\cos 2\theta}\left(\frac{1}{\rho}\frac{\mathrm{d}\rho}{\mathrm{d}x}\right)^{-1}_{x=x_{\mathrm{res}}} \gg 1 \quad (9)$$

where x_{res} is the position of the resonance layer. Inside the core, $\left(\frac{1}{\rho}\frac{\mathrm{d}\rho}{\mathrm{d}x}\right)^{-1}_{x=x_{\mathrm{res}}} \sim 1$ km and therefore the adiabaticity condition is satisfied if

$$\Delta m^2\frac{\sin^2 2\theta}{\cos 2\theta} \gg 10^{-3}\mathrm{eV}^2\left(\frac{E_\nu}{10\mathrm{MeV}}\right), \quad (10)$$

which is easily satisfied by $\Delta m^2 \gtrsim 10^4\mathrm{eV}^2$ as long as $\sin^2\theta \gg 10^{-7}$.

Thus, we find that a substantial fraction of neutrinos may get converted to sterile neutrinos and escape the core of the star, if the mass of the sterile neutrinos is such that $10^4\mathrm{eV}^2 \lesssim \delta m^2_{as} \lesssim 10^8\mathrm{eV}^2$ (Kainulainen, Maalampi & Peltoniemi 1991).

The mass difference of this magnitude cannot solve the observed solar and atmospheric neutrino problem, but the possibility of three active neutrinos explaining these anomalies and a heavy sterile neutrino of mass $m_s \sim$ keV still stays open. At very low values of the mixing angle ($\sin^2 2\theta \lesssim 10^{-2}$), the CHOOZ experiment does not put any constraints on the neutrino mass (Berezinsky, Raffelt & Valle 2000 and references therein). The tritium β decay experiment (Berezinsky, Raffelt & Valle 2000, and references therein) indicates $\sin^2\theta_{es}m_s^2 \lesssim (2.5)^2\mathrm{eV}^2$, which is consistent with $m_s^2 \gtrsim 10^4\mathrm{eV}^2$ and $\sin^2\theta \lesssim 10^{-4}$ (note that the adiabaticity condition Eq.(10) holds at this value of θ). The neutrinoless double beta decay constraint $\sin^2\theta_{es}m_s \lesssim 0.2$eV, is satisfied for $m_s \gtrsim 10^2$eV and $\sin^2\theta_{es} \lesssim 10^{-3}$. The cosmological constraint, which comes from the requirement that the universe should not be overclosed, can be sidestepped if ν_s has a small enough lifetime, which would require physics beyond the standard model. The bounds from the big bang nucleosynthesis still allow one species of sterile neutrinos.

Since the sterile neutrinos escape the core over a timescale of a few milliseconds, the number of neutrinos escaping and their angular distribution is sensitive to the instantaneous distribution of neutrino production sites. Since thermalization cannot occur over such small times, and since the neutrino production rate is sensitive to the local temperature at the production point, the inhomogeneities during the collapse phase get reflected in the inhomogeneities in the escaping neutrino fluxes and their distributions. The asymmetries in these distributions can give rise to dipole moments, which can help in understanding the high peculiar velocities of pulsars (Kusenko & Segrè 1998), or quadrupole moments, which can give rise to gravitational waves, as suggested here.

The *fraction* of neutrinos that can escape in the first few milliseconds is, however, *small*. Firstly, the neutrinos have to be produced roughly within one mean free path of the resonance surface. Secondly, since m_s are the heaviest neutrino species, the sign of the effective potential $V(x)$ and the resonance condition indicates that only $\nu_e, \bar{\nu}_\mu$ and $\bar{\nu}_\tau$ can undergo resonant conversions. Considering that at least six species of νs may oscillate, perhaps this factor is not just of 1/2, but a factor of 1/6. Also, in order to escape, the direction of neutrinos has to be from higher densities to lower densities, which provides a further factor of two reduction. Thence, we are left with $\lesssim 10\%$ of the total ν-flux as sterile neutrinos.

2.1. *GWs from ν-Flavor Conversions*

With the assumption that the neutrino oscillation timescale (the one for the core bounce at the PNS interior) is $T_{\nu_{\bar{\tau},\bar{\mu}}\longleftrightarrow\nu_s} \sim R_{PNS}/c \sim 1.0\times 10^{-3}$s, where R_{PNS} is the PNS radius, we can expect the GWs burst to appear similar to a delta Dirac function centred around 1kHz (the GWs frequency f_{gw}) superimposed onto the overall waveform as computed by Burrows & Hayes (1995). Then, the characteristic, normalized GW amplitude produced during the SNE non-spherical outgoing front of the oscillating s-neutrinosphere is computed using Eq.(2) as

$$h \sim 2.6\times 10^{-21}\left(\frac{55\mathrm{kpc}}{\mathrm{D}}\right)\left(\frac{10\% E_{\nu_s}}{5.10^{53}\mathrm{erg}}\right)\left(\frac{1\mathrm{ms}}{\Delta t_{CB}}\right)\mathrm{Hz}^{-1/2}, \quad (11)$$

provided 10% of the total neutrinos produced undergo oscillations $\nu_{\bar{\tau},\bar{\mu}} \longleftrightarrow \nu_s$, with luminosity $L_\nu = 3\times 10^{55}\mathrm{ergs}^{-1}$, i. e., $\sim 3\times 10^{52}$erg released on a timescale $\Delta t_{CB} \sim 1$ms (Burrows, Hayes & Fryxell 1995; Burrows & Hayes 1995), for a source distance 55kpc (i. e., Large Magellanic Cloud). Therefore, a GWs signal such as this may be detectable by the LIGO, VIRGO and TIGAs with a large signal-to-noise ratio. It is clear then, that for galactic distances this sort of GWs bursts may also be detectable at higher amplitudes (see Figure 1).

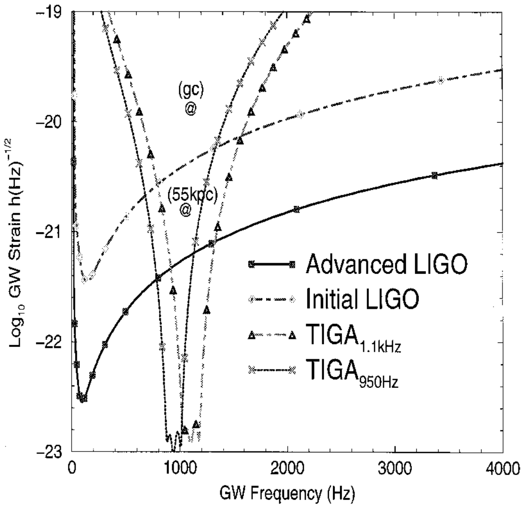

Figure 1: Characteristics of the GWs burst generated by the $\nu_{\mu,\tau} \longrightarrow \nu_s$ oscillation mechanism. For distances upto the Large Magellanic Cloud such a signal would be detectable only by LIGO-II. However, it may be observable even by *Initial* LIGO for sources at the galactic center, as indicated. Truncated Icosahedral GWs Antennas could also observe such events.

3. NATAL RECOIL KICKS AND GWS AT CORE BOUNCE

The evolution of the rapidly rotating PNS during this phase is dominated by vigorous entropy-driven convective motions. Then, contrary to the arguments given by Janka & Raffelt (1999) concerning the existence of a structured PNS (mantle and atmosphere), we suggest that at this post-bounce evolutionary stage (a few ms) there is no room for such a stratified structure to appear. Simply, the convection overturn timescale $\sim ms$ (Müller & Janka 1997) dominates the PNS evolution, resulting in an entangled configuration. Consequently, the analysis by Janka & Raffelt (1999) is not applicable to our model. Instead, we follow the lines of Kusenko & Segrè (1998), and assume the ν-oscillation mechanism drives the PNS kick. Furthermore, since one expects the neutrino oscillation-driven GW burst to be emitted at core bounce ($M_{core} \sim 1.4M_\odot$, $R_{core} \sim 10^7$cm), its timescale ($\tau_{CB} \sim 1$ms) is very short compared to the neutrino diffusion time (the PNS Kelvin-Helmholtz neutrino cooling) ~ 10s, the process itself can be considered as a *transient stage.* As pointed out in Janka & Raffelt (1999) for this phase the neutrino-driven neutron star kick mechanism proposed by Kusenko & Segrè (1997,1998) may be at work.

The relative neutrino recoil momentum is defined as (Kusenko & Segrè 1997,1998)

$$\frac{\Delta\kappa}{\kappa} = \left[\frac{4e\sqrt{2}}{\pi^2}\right]\left(\frac{\mu_e \mu_n^{1/2}}{m_n^{3/2}T^2}\right) B = 0.01\left(\frac{B}{1.2\times 10^{15}\mathrm{G}}\right), \tag{12}$$

where μ_e and μ_n, define the electron and neutron chemical potential, respectively, m_n the neutron mass, and $T \sim 30$MeV the temperature at the s-neutrinosphere. This anisotropic temperature distribution (the matter-induced neutrino potentials) is created by the strong magnetic field ($B \sim 10^{15}$G) assumed to be developed during the supernova core collapse[1]. The *critical* magnetic field strength can be estimated by using the Klúzniak & Ruderman (1998) mechanism in which the PNS core rotates differentially: $B_c^2 = f_{st}\rho_{ns}c_s^2$. As stressed above, a scarcely stratified PNS $\longrightarrow f_{st} \sim 10^{-3}$, which yields $B_c \sim 2\times 10^{15}$G, for PNS sound speed and density $c_s = c/\sqrt{3}$ and $\rho_{ns} \sim 4\times 10^{14}\mathrm{gcm}^{-3}$, respectively.

Now, because the general relativistic description of the gravitational radiation generated by escaping neutrinos during the supernova explosion, as demonstrated by Epstein (1978), requires that the collapsing star is ellipsoidally deformed due to rotation and anisotropic ν-emission [Eq.(1)]. *We can conjecture that such an asymmetry "e" is determined fundamentally by the shape of the neutrino-sphere by the time the oscillation develops.* Therefore, the anisotropic neutrino emission [responsible in the picture by Kusenko & Segrè (1997,1998) for the nascent neutron star kicks] might be considered the appropriated scaling of the eccentricity appearing in Eq.(1). Assuming that is the case, using the result above for $B_c \sim 10^{15}$G and Eq.(1), we can write $h \sim |\theta_{TT}^{ij(\nu)}|$ which implies

$$e^2 \sim \frac{\Delta\kappa}{\kappa} = 0.01, \tag{13}$$

from which we derive the angle θ the final jet opens respect to $\vec{B}$, i. e., the angle under which one, placed at the ν-sphere poles, sees the large part of escaping oscillating νs:

$$\tan\theta \simeq e = \frac{a-b}{a+b} = 0.1 \longrightarrow \theta = 5.7\ deg. \tag{14}$$

This result provides a well-fundamented astrophysical origin for the observed supernovae beamed jets and asymmetries (Nisenson & Papalolios 1999; Gaensler 1999).

4. CONCLUSIONS

We conclude by saying that if the mechanism ν-oscillations-GWs bursts is really realized in nature, then it might afford a natural bridge to link and explain two until now difficult problems in SNe physics, i. e., a) what is the origin of the asymmetries observed in SNe ejecta and remnants? b) what triggers the bi-polar beamed jets demonstrated by Nisenson & Papaliolios (1999) to be in association with SN 1987A? *In our picture, this asymmetry in the power of the jets, which might also kick the PNS; so as to explain observed pulsar velocities and spins, is fundamentally created by the neutrino oscillation mechanism.*

The prospective detection of such bursts in future experiments will be a breakthrough in fundamental physics for, as pointed out above, there exists the possibility that an important parcel of the oscillating neutrinos to arrive at Earth (as during SN1987A) nearly by the time the associated GWs bursts do (the ones produced when changing

[1]See also the process leading to the formation of *magnetars*, i. e., strongly magnetized newly-born NSs (Duncan & Thompson 1992).

flavors $\nu_{\mu,\tau} \leftrightarrow \nu_s$). Such measurements of the arriving times may lead to tight constraints on the neutrino mass spectrum.

We also speculate that despite the small probability for flavor reconversion $\nu_s \longrightarrow \nu_{\mu,\tau}$, a number of sterile neutrinos might convert back to active ones outside the SNE core, stream away freely (matter density in the SN mantle is quite low to trap them), and arrive on Earth in very early stages of the SNE compared to those escaping in the ν-diffusion time. This *anomalous ν-burst* from the s-reconversion may be observed by ν-telescopes.

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