Neutrino propagation in binary neutron star mergers in presence of nonstandard interactions

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We explore the impact of nonstandard interactions on neutrino propagation in accretion disks around binary neutron star mergers remnants. We show flavor evolution can be significantly modified even for values of the nonstandard couplings well below current bounds. We demonstrate the occurrence of I resonances as synchronized MSW phenomena and show that intricate conversion patterns might appear depending on the nonstandard interaction parameters. We discuss the possible implications for nucleosynthesis.

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I. INTRODUCTION

The origin of heavy elements remains one of the key open questions in nuclear astrophysics. Nucleosynthetic abundances produced in the rapid neutron capture process ($r$-process) are formed in dense neutron-rich environments \cite{1} including core-collapse supernovae, accretion disks around black holes or binary compact systems. It was first recognized in \cite{2} that $r$-process nuclei could be formed in neutron star matter. The occurrence of a weak or of a strong $r$-process depends mainly upon the astrophysical conditions and the properties of exotic nuclei. In particular, conditions for a strong $r$-process are met in neutron star mergers, whereas elements with $A > 130$ are not produced in core-collapse supernovae considered for long a favorite $r$-process site (see e.g. \cite{3,4}).

The recent observation of gravitational waves from a binary neutron star (BNS) merger event in coincident with a short gamma-ray bursts and a kilonova constitute the first experimental evidence for $r$-process nucleosynthesis in such sites \cite{5,6}. Weak interactions and neutrinos bring the ejecta to being hot. The role of neutrino flavor evolution in these environments still needs to be fully assessed.

Calculations of nucleosynthetic abundances of heavy elements show that dynamical ejecta can produce a strong $r$-process while a weak $r$-process can take place in neutrino driven winds \cite{7}. In fact, the presence of a significant amount of neutrinos in neutrino-driven winds influences the build-up of heavy elements through the electron neutrino and antineutrino interactions with neutrons and protons respectively. Such interactions tend to be detrimental to the $r$-process since they reduce the number of available neutrons. The occurrence of flavor conversion phenomena can produce swappings of the neutrino spectra and modify the interaction rates that determine the electron fraction, as shown in numerous studies (see e.g. Refs.\cite{8-10}).

Two decades of experiments have contributed to determine the neutrino squared-mass differences and the mixing angles responsible for vacuum oscillations and for some of the flavor conversion mechanisms in matter. Open questions remain concerning the absolute neutrino mass ordering, the neutrino nature (Dirac versus Majorana), the existence of sterile neutrinos and of CP violation in the lepton sector. The combined analysis of available data indicate that the neutrino mass ordering is normal and the Dirac CP violating phase approximately $1.5\pi$, although statistical significance is still low \cite{11}.

Flavor evolution in astrophysical environments reveals a variety of mechanisms due to the nonlinear many-body nature of neutrino propagation in presence of neutrino self-interactions. Regions of flavor instabilities close to the neutrinosphere are found including collective phenomena such as the bipolar instability in core-collapse supernovae \cite{12,13}, or the matter-neutrino resonance (MNR) in neutron star mergers \cite{9}.

The presence of nonstandard interactions can alter flavor conversion. Limits on nonstandard neutrino self-interactions are rather loose \cite{14}, whereas scattering and oscillation experiments furnish tight bounds on nonstandard neutrino-matter interactions (NSI) \cite{15,17}. The first measurement of neutrino-nucleus coherent scattering provides interesting NSI constraints Ref.\cite{18}. The existence of NSI would modify the interpretation of oscillation experiments in particular for the inferred values of the squared-masses and the mixings and could furnish an explanation of observed anomalies.

Within a supernova core, flavor changing neutral current interactions would impact the scattering rates and the electron fraction, altering the infall \cite{19}. Nonstan-
standard four fermion neutrino self-interactions might produce flavor equilibration both in normal and inverted mass ordering \cite{20} or could modify the neutronisation burst signal of a supernova explosion \cite{21}. Novel interactions can also produce resonant conversion near the neutrinosphere and influence the $r$-process in supernovae \cite{22}. In particular, the I resonance – a Mikheev Smirnov Wolfenstein (MSW)-like resonance \cite{23} \cite{24} – can take place due to the cancellation between the matter and the NSI contributions to the neutrino Hamiltonian \cite{25}. Refs. \cite{26} \cite{27} have pointed out that the I location appears to be little affected by neutrino self-interactions. Moreover Ref.\cite{26} has shown that NSI contributions can provide the necessary cancellation for the occurrence of MNR in supernovae.

Neutrino propagation in accretion disks of compact objects presents differences with supernovae. Spherical symmetry breaking makes simulations demanding, even when the geometry of the neutrino emission at the neutrinospheres is the simplest. Detailed simulations of neutrino emission in accretion disks show that the electron antineutrino luminosities are larger than the electron neutrino ones and that the muon and tau neutrino luminosities are small. Such excess of electron antineutrinos over neutrinos can produce the MNR when the matter and the neutrino self-interaction potentials cancel \cite{10}. A MNR is a series of MSW-like resonances \cite{10} \cite{28} \cite{29} that can impact $r$-process nucleosynthesis \cite{10}. A perturbative analysis unravels the conditions for nonlinear feedback that produces the matching of the self-interaction and the matter contributions and maintains the MNR over several tens of kilometers. On the other hand, such a matching is impeded by the geometrical factors radial dependence in the case of helicity coherence \cite{30}. Conditions for the MNRs can be met in detailed simulations of the disk around BNS merger remnants \cite{31} \cite{32}. An analysis with non-stationary evolution shows that conversion can occur on very short timescales \cite{33}. These modes termed as fast have been first pointed out in the supernova context \cite{34}.

Our goal is to explore the role of nonstandard neutrino-matter interactions on the neutrino evolution in accretion disks around binary neutron star merger remnants. First we focus on the NSI impact on flavor evolution and discuss the I resonance. We shed a new light on its mechanism and show that the neutrino self-interactions can produce I resonances as synchronized MSW effects. Moreover we present how NSI can modify both location and adiabaticity of the MNRs. Our calculations are based on the matter density profiles and electron fraction taken from detailed astrophysical simulations of BNS remnants \cite{35}. We discuss the effects of nonstandard interactions on the electron fraction $Y_e$, a key parameter for $r$-process nucleosynthesis in neutrino-driven winds, in the light of the study of Ref.\cite{24}.

The manuscript is structured as follows. Section II presents the model with NSI. Numerical results on the flavor evolution for different sets of NSI parameters are given in Section III. The NSI effects on the I and MNR resonances are discussed. Section IV is a conclusion.

II. THE MODEL

A. Neutrino evolution equations in presence of nonstandard interactions

The evolution of a system of neutrinos and antineutrinos in an astrophysical environment is governed by the Liouville Von-Neumann equations (we will use $\hbar = c = 1$)

\[ i\dot{\rho} = [h, \rho], \quad i\dot{\bar{\rho}} = [\bar{h}, \bar{\rho}], \]

where $\rho$ and $\bar{\rho}$ are single-particle density matrices, $h$ and $\bar{h}$ mean-field Hamiltonians for neutrinos and antineutrinos respectively. For a detailed derivation of Eqs. (1) see e.g. \cite{36}. The mean-field equations (1) correspond to the lowest order truncation of the Born-Bogoliubov-Green-Kirkwood-Yvon hierarchy for the two-point correlators \cite{37}

\[ \rho_{ij}(t, \vec{q}, -\vec{q}', -) = \langle a_j^\dagger(t, \vec{q}, -) a_i(t, \vec{q}', -) \rangle, \]

\[ \bar{\rho}_{ij}(t, \vec{q}, +, \vec{q}'', +) = \langle a_j^\dagger(t, \vec{q}, +) a_i(t, \vec{q}'', +) \rangle, \]

that depend on time, on particle momentum $\vec{q}$ and on positive ($h = +$) or negative ($h = -$) helicity states. The labelling $i, j$ are either mass or flavor indices. The creation and annihilation operators $a^\dagger$ ($b^\dagger$) and $a$ ($b$) for neutrinos (antineutrinos) satisfy the nonzero equal-time anti-commutation relations. The diagonal elements ($i = j$) are the expectation values of the number operator, while the non-diagonal ones ($i \neq j$) include the decoherence due to the neutrino mixing. In the following the explicit dependence on the momentum and helicity variables, as well as the time dependence will not be shown systematically to simplify notations.

Since neutrinos propagate through an astrophysical background, the mean-field Hamiltonians include the neutrino charged- and neutral-current interactions with the particles composing the medium, usually electrons, protons and neutrons, as we will be considering in the present work. Therefore $h$ is given by

\[ h = h_0 + h_{\text{mat}} + h_{\nu\nu}, \]

where the first term corresponds to the vacuum Hamiltonian, the second to the neutrino standard and nonstandard interactions with matter and the last one to neutrino self-interactions. The same expression holds for
neutrino self-interaction Hamiltonian parameters are rather loose, with the exception of (solar–like). One can see that the bounds on the NSI if matter is composed only of protons and electrons the neutrinosphere.

The heavy quark content of the nucleon is neglected.

3 We remind that the standard neutrino-matter neutral current contributions are not included since they are proportional to the identity matrix and therefore do not produce flavor modifications.

4 From now on, only the radial dependence of all quantities is retained and not explicitly shown to simplify notations.

5 In this work, we neglect the bending of the trajectory due to strong gravitational fields.

\[ \hat{h} = \text{diag}(E_i), \quad E_{i=1,N_f} \] being the eigenenergies of the propagation eigenstates with \( N_f \) the number of neutrino flavors. The quantity \( U \) is the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) \( N_f \times N_f \) unitary matrix relating the mass to the flavor basis [33].

As for the matter term, it comprises the standard contribution from neutrino-electron charged currents and a nonstandard term related to neutrino-matter interactions

\[ h_{\text{mat}} = h_{\text{CC}} + h_{\text{NSI}}, \]

where \( h_{\text{CC}} = \text{diag}(V_{\text{CC}}, 0) \) and \( V_{\text{CC}} = \sqrt{2} G_F \rho_e \), with \( G_F \) the Fermi coupling constant and \( \rho_e \) the net electron number density. Note that here anisotropic contributions to the matter Hamiltonian are not included. The nonstandard interaction Hamiltonian is

\[ h_{\text{NSI}} = \sqrt{2} G_F \sum f n_f \epsilon_f, \]

where a sum over the electron, down and up quark number densities is performed \( (f = e, d, u) \). The \( \epsilon \) matrices correspond to the nonstandard interactions couplings, constrained by several observations [15–18].

In the case of three neutrino flavors, these are [15]

\[
\begin{pmatrix}
|\epsilon_{ee}| < 2.5 \\
|\epsilon_{\mu\mu}| < 0.21 \\
|\epsilon_{\mu\tau}| < 1.7 \\
|\epsilon_{\tau\mu}| < 0.21 \\
|\epsilon_{\tau\tau}| < 9.0
\end{pmatrix},
\]

if matter is composed only of protons and electrons (solar–like). One can see that the bounds on the NSI parameters are rather loose, with the exception of \( \epsilon_{\mu\mu} \).

The third contribution in Eq. (3) corresponds to the neutrino self-interaction Hamiltonian

\[ h_{\nu\nu} = \sqrt{2} G_F \sum_\alpha \int (1 - \hat{q} \cdot \hat{p}) \left[ d\nu_\alpha \rho_{\nu\alpha}(\hat{p}) - d\nu_\alpha \bar{\rho}_{\nu\alpha}(\hat{p}) \right], \]

where the quantity \( d\nu_\alpha \) denotes the differential number density of neutrinos (antineutrinos), the underline refers to the neutrinos initially born with \( \alpha \) flavor at the neutrinosphere.

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The heavy quark content of the nucleon is neglected.

\[ \rho = \left( \begin{array}{cc} \rho_{ee} & \rho_{ee} \\ \rho_{e\tau} & \rho_{\tau\tau} \end{array} \right), \]

and similarly for \( \bar{\rho} \). The vacuum Hamiltonian Eq. (4) involves the PMNS matrix that for three flavors depends on three measured mixing angles and three unknown CP violating phases (one Dirac- and two Majorana-type) [33].

In two flavors, these fundamental parameters reduce to one mixing angle \( \theta \) (one phase as well in the case of Majorana neutrinos). Therefore the vacuum contribution becomes

\[ h_0 = \omega \left( \begin{array}{cc} -c_{2\theta} & s_{2\theta} \\ s_{2\theta} & c_{2\theta} \end{array} \right), \]

with \( \omega = \frac{\Delta m^2}{4E} \), \( \Delta m^2 = m_2^2 - m_1^2 \) with \( m_1, m_2 \) the mass values of the mass eigenstates, \( E = q \) the neutrino energy, \( s_{2\theta} = \sin 2\theta \) and \( c_{2\theta} = \cos 2\theta \).

For the standard matter Hamiltonian in Eq. (5) we write

\[ V_{\text{CC}} = \lambda Y_e, \]

where \( \lambda = \sqrt{2} G_F n_B \), with \( n_B \) the baryon number density, and \( Y_e = \rho_e / (n + p) \) the electron fraction, with \( n \) and \( p \) the neutron and proton number densities, respectively. As in Refs. [30–32] our investigation is anchored to the detailed simulations in which the BNS merger remnant is a central object, lasts up to 200 ms and has about 30 km radius. We take information on the baryon number densities and electron fraction from cylindrical averages of detailed three-dimensional Newtonian simulations [34].

In our two-dimensional model neutrino propagate with an azimuthal symmetry axis from point \( (x_0, z_0) \), at the neutrinosphere following a straight line trajectory characterized by a radial \( r \) and an angular \( \theta_q \) variables (Figure 1). Note that we approximate the neutrinospheres as infinitely thin disks of radii \( R_\nu \) that are flavor dependent, as done in Refs. [9, 10, 30–32].

In two-flavors, by retaining only the nonstandard contribution Eq. (5) with loosest constraints, we get for the vacuum contribution Eq. (4)
fraction that, at the MSW resonance in the Sun, with an electron

Finally we follow Ref. \[26\] and impose the requirement

which for up and down quarks can be rewritten as $Y_d = 2 - Y_e$ and $Y_u = 1 + Y_e$. The NSI contribution is then

$$h_{\text{NSI}} = \sqrt{2} G_F n_B \left[ Y_e \epsilon^e + (1 + Y_e) \epsilon^u + (2 - Y_e) \epsilon^d \right].$$

Finally we follow Ref. \[26\] and impose the requirement

that, at the MSW resonance in the Sun, with an electron fraction $Y_\odot \approx 0.7$, the NSI contribution should vanish as

no effect has been observed (see also \[39\]), namely

$$Y_\odot \delta \epsilon^e + (1 + Y_\odot) \delta \epsilon^u + (2 - Y_\odot) \delta \epsilon^d = 0,$$

with $\delta \epsilon^f = \epsilon^e_{\text{iso}} - \epsilon^f_{\text{iso}}$. This equation gives a relation between $\delta \epsilon^e$ as a function of $\delta \epsilon^u, \delta \epsilon^d$. The off-diagonal

couplings $\epsilon^e_{\text{iso}}, \epsilon^u_{\text{iso}}, \epsilon^d_{\text{iso}}$ are fixed at the same value $\epsilon_0$. As a result the NSI Hamiltonian only depends on two NSI

parameters, the diagonal one $\delta \epsilon^u$ and the off-diagonal $\delta \epsilon^e$

$$h_{\text{NSI}} = \lambda \left( \frac{Y_\odot - Y_e}{Y_\odot} \right) \delta \epsilon^e \left( 3 + Y_e \right) \epsilon_0,$$

with the constraints $|\delta \epsilon^e| \lesssim \mathcal{O}(10)$ and $|\epsilon_0| \lesssim \mathcal{O}(1)$. For the neutrino self-interaction Hamiltonian Eq.(8) we

assume, as done in previous works \[9\] \[10\] \[30\] \[32\], that

$$\rho_\nu(r, \hat{p}) = \rho_\nu(r, \hat{p}),$$

namely that the angular dependence of the neutrino density

matrix is not retained. As a consequence, the neutrinos that are coupled by the self-interaction term have the same flavor
evolution. The linearized analysis of Ref. \[33\] has included the angular dependence. We assume in our calculations that neutrinos are emitted as Fermi-Dirac distributions $f_\nu$, with luminosities $L_\nu$, average energies $\langle E_\nu \rangle$ at the neutrinosphere with neutrinosphere radii $R_\nu$ given in Table I. Table VII and Figures 14 – 16 of Ref. \[32\] show the current spread on $L_\nu$ and $\langle E_\nu \rangle$ according to available simulations of neutrino emission in binary neutron star mergers. Concerning the neutrino luminosities and average energies, these are stable for long

times (see Ref. \[35\]). By using Eqs.(8) and (17), the neutrino self-interaction term becomes

\begin{equation}
\rho_\nu(r, \hat{p}) = \rho_\nu(r, \hat{p}),
\end{equation}

\[6\] Note that here we show the full dependence on the variables for clarity.

\[\begin{equation}
\begin{split}
h_{\nu\nu}(r, q, \ell_q) &= \sqrt{2} G_F \sum_{\alpha = e, \mu} \int_0^\infty dp \left[ G_{\nu_\alpha}(r, \ell_q) \rho_{\nu_\alpha}(r, p) \frac{L_{\nu_\alpha} f_{\nu_\alpha}(p)}{\pi^2 R_\nu^4 \langle E_{\nu_\alpha} \rangle} - \bar{\rho}_{\nu_\alpha}(r, p) G_{\bar{\nu}_\alpha}(r, \ell_q) \frac{L_{\bar{\nu}_\alpha} f_{\bar{\nu}_\alpha}(p)}{\pi^2 R_\nu^4 \langle E_{\bar{\nu}_\alpha} \rangle} \right],
\end{split}\end{equation}\n
where the geometrical factor $G_{\nu_\alpha}$ reads

$$G_{\nu_\alpha}(r, \ell_q) = \int_{\Omega_{\nu_\alpha}} d\Omega (1 - \hat{q} \cdot \hat{p}),$$

with $\Omega_{\nu_\alpha}$ the angular variables and similarly for $G_{\bar{\nu}_\alpha}$ for

\begin{equation}
\begin{pmatrix}
|\epsilon_{ee}| < 2.5 \\
|\epsilon_{er}| < 1.7 \\
|\epsilon_{rr}| < 9.0
\end{pmatrix},
\end{equation}

We rewrite the NSI potential Eq.(9) in terms of the fermion fraction $Y_f$. In fact, using the charge neutrality of the medium, we get the relation

$$Y_f \equiv \frac{n_f}{n_B},$$

which for up and down quarks can be rewritten as $Y_d = 2 - Y_e$ and $Y_u = 1 + Y_e$. The NSI contribution is then

\begin{equation}
\begin{split}
h_{\text{NSI}} &= \sqrt{2} G_F n_B \left[ Y_e \epsilon^e + (1 + Y_e) \epsilon^u + (2 - Y_e) \epsilon^d \right].
\end{split}\end{equation}\n
\[\begin{equation}
\begin{split}
h_{\nu\nu}(r, q, \ell_q) &= \sqrt{2} G_F \sum_{\alpha = e, \mu} \int_0^\infty dp \left[ G_{\nu_\alpha}(r, \ell_q) \rho_{\nu_\alpha}(r, p) \frac{L_{\nu_\alpha} f_{\nu_\alpha}(p)}{\pi^2 R_\nu^4 \langle E_{\nu_\alpha} \rangle} - \bar{\rho}_{\nu_\alpha}(r, p) G_{\bar{\nu}_\alpha}(r, \ell_q) \frac{L_{\bar{\nu}_\alpha} f_{\bar{\nu}_\alpha}(p)}{\pi^2 R_\nu^4 \langle E_{\bar{\nu}_\alpha} \rangle} \right],
\end{split}\end{equation}\n
where the geometrical factor $G_{\nu_\alpha}$ reads

$$G_{\nu_\alpha}(r, \ell_q) = \int_{\Omega_{\nu_\alpha}} d\Omega (1 - \hat{q} \cdot \hat{p}),$$

with $\Omega_{\nu_\alpha}$ the angular variables and similarly for $G_{\bar{\nu}_\alpha}$ for
In this section, we have show examples with NSI parameters $\delta\epsilon^n \in \left[-0.9, -0.7\right]$. These are the parameters for which we observed the presence of the I resonance in most of the trajectories explored that were relevant for nucleosynthesis \[7\]. Negative values of $\delta\epsilon^n$ with a greater absolute value lead to the disappearance of the I resonances, as the matter potential $V_M$ Eq.\[21\] would always be negative on the region of space studied, and would also make the MNR further away. Negative values of $\delta\epsilon^n$ with a smaller absolute value would still present I resonances, but in a different region of space, and would also shift the MNRs. It is worth noting that positive values of $\delta\epsilon^n$ have also been considered as they can shift the MNR closer to the neutrinosphere.

As for the value of $\epsilon_0$, we have restricted ourselves to values smaller than $10^{-3}$. Indeed, values larger than that creates oscillation patterns analogous to vacuum oscillations, but driven by the large matter off-diagonal element. These oscillations have a very short wavelength (shorter than a kilometer), and can start as soon as the neutrino propagation begins. Given that, in our calculations, we assume that neutrinos are free streaming, our results are reliable only if flavor conversions happen well outside the neutrinospheres, and therefore using larger values of $\epsilon_0$ would give unphysical results. These oscillations appearing because of a larger $\epsilon_0$ also have a large amplitude, making the behaviors difficult to analysis. For all these reasons, we chose to work with a value of $\epsilon_0$ well below the current experimental constraints.

### A. New conditions for the I resonance

The presence of NSI produces a new MSW-like resonance, called the inner (I) resonance [25]. Refs.\[25\]\[27\] have shown that its occurrence is due to the matter terms only. In the present work we will be discussing two situations in which the I resonance occurs: \(i\) the self-interaction is sub-dominant, in accord with \[25\]\[27\]; \(ii\) the neutrino self-interaction dominates and leads to a I resonance as a synchronized MSW mechanism. We explore this scenario using the SU (2) spin formalism.

#### 1. I resonance with negligible self-interaction

The I resonance occurs when the difference between the diagonal elements of the total Hamiltonian goes to zero, requiring for the total matter potential to meet the condition

$$V_M \equiv \lambda \left[ Y_e \left( Y_e - Y_e \delta \epsilon^n \right) \right] \approx 2 \omega c_{2\theta} - \left( h_{ee} - h_{\nu\nu} \right). \quad (21)$$

Refs.\[25\]\[27\] have pointed out that the presence of $\nu\nu$ self-interactions have negligible effects on the location and adiabaticity of the I resonance, thus making it to occur when the matter potential Eq.\[21\] is very small.
First we consider here a case in which the self
interaction potential is sub-dominant compared to the
matter one. In such cases the location of the I resonance
coincides with the point where the matter potential \( V_{\text{mat}} \)
becomes very small, which is possible in the presence
of NSI because of a cancellation between the standard
matter term and the nonstandard contribution. Figure
2 (left panel) presents the difference of the diagonal ele-
ments of the total neutrino Hamiltonian (solid line),
matter potential \( V_M \) (dashed line) Eq.(21) in presence
of NSI contributions with \( \delta\epsilon^\alpha = -0.7 \) and \( \epsilon_0 = 1 \times 10^{-4} \) and self-interaction oscillated
potential (dotted line), as a function of distance from the emission point. The initial parameters are \( z_0 = -30 \text{ km}, z_0 = 20 \text{ km}, \) and \( \theta_q = 55^\circ. \) Middle and right panels : Survival probabilities for neutrinos (middle), antineutrinos (right). Different energies corresponding to different colors as well as averaged probability (dotted line) are indistinguishable. The results are obtained by using baryon densities and electron fraction from the detailed simulations 15.

The survival probabilities for neutrinos and antineu-
trinos as well as the averaged one are shown in Figure 2
for different neutrino energies (middle and right panels).
Given a specific matter profile, the resonance location
only depends on the value of the diagonal NSI parameter,
\( \delta\epsilon^\alpha; \) whereas the value of \( \epsilon_0 \) impacts the adiabatic-
ity. For the case shown, the I resonance is adiabatic and
induces significant conversion for both neutrinos and an-
tineutrinos. It is worth noting that even in the presence
of a small \( \epsilon_0 \) parameter, the flavor conversion behaviors
stay independent of the energy. This is due to the fact
that the off-diagonal self-interaction contribution to the
Hamiltonian is, at the considered location, much larger
than the vacuum one, therefore suppressing the energy
dependence.

2. I resonance as a synchronized MSW

While exploring the parameter space and different tra-
jectories for the neutrino propagation, we have encoun-
tered situations where, although the self-interaction un-
oscillated potential is several orders of magnitude larger
than the matter potential, a I resonance takes place and
leads to significant flavor conversions. Figure 3 shows a
typical example of this situation with the NSI parameters
\( \delta\epsilon^\alpha = -0.88 \) and \( \epsilon_0 = 1 \times 10^{-4} \). One can see that
although the unoscillated self-interaction potential \( \mu \) dominates the matter one \( \lambda \) (11), flavor conversions oc-
cur at the same location where the I resonance condition is
fulfilled. Note that the difference between the self-
interaction oscillated diagonal elements do cancel at the
same point. We will be unraveling this effect in the light
of synchronized flavor conversions.

a. Spin description
In order to describe this phe-
omenon, we use the SU(2) isospin formalism in flavor
space. The effective isospin vector \( \vec{P}_{\nu_\alpha} (r, q) \) denoting
a neutrino of initial flavor \( \alpha \) is related to the neutrino dens-
ity matrix according to

\[
\rho_{\nu_\alpha} (r, q) = \frac{1}{2} \left( \mathbb{I} + \vec{\sigma} \cdot \vec{P}_{\nu_\alpha} (r, q) \right),
\]

and similarly for antineutrinos, where \( I \) is the \( 2 \times 2 \) identity
matrix and \( \vec{\sigma} = (\sigma_x, \sigma_y, \sigma_z) \) is a vector in flavor space
whose components are the Pauli \( \sigma \) matrices. In this theo-
retical framework, the Liouville Von-Neumann equations
are replaced by precession equations for \( \vec{P}_{\nu_\alpha} (r, q) \) with
an effective magnetic field defined as

\[
h (r, q) = \frac{1}{2} \left( \mathbb{I} + \vec{\sigma} \cdot \vec{B} (r, q) \right),
\]

and receiving three contributions

\[
\vec{B} (r, q) = \vec{B}_{\text{vac}} (q) + \vec{B}_{\text{mat}} (r) + \vec{B}_{\nu\nu} (r),
\]

where
Note that the expressions for $\vec{B}_{\nu_e}$ and $\vec{B}$ are analogous to Eqs.~(22) and (24) respectively. In the antineutrino case, the vacuum contribution in Eq.~(24) has a minus sign. The vacuum term is given by

$$\vec{B}_{\text{vac}} = 2\omega \vec{B}_0 = 2\omega \begin{pmatrix} 0 \\ s_{2\theta} \\ -c_{2\theta} \end{pmatrix},$$

(25)

while the matter term includes the standard and non-standard contributions

$$\vec{B}_{\text{mat}} = \lambda \begin{pmatrix} Y_e(0,0) \\ 2(3 + Y_e) \text{Re}c_0 \\ -2(3 + Y_e) \text{Im}c_0 \end{pmatrix}.$$  

(26)

The third term in Eq.~(24) comes from the self-interaction term of the neutrino Hamiltonian

$$\vec{B}_{\nu\nu} = \sqrt{2} G_F \sum_{\alpha = e, x} \int_{0}^{\infty} dp \left( G_{\nu,\nu_{\alpha}}(p) \vec{P}_{\nu_{\alpha}}(p) \right)$$

$$-G_{\bar{\nu},\bar{\nu}_{\alpha}}(p) \vec{P}_{\bar{\nu}_{\alpha}}(p),$$

(27)

where $j_{\nu_{\alpha}}(p) = \frac{\mu_{\nu_{\alpha}} f_{\nu_{\alpha}}(p)}{\pi R_{\nu_{\alpha}}(E)}$ and similarly for antineutrinos. Note that the explicit $r$-dependences are not shown for readability.

In order to describe the collective neutrino mode associated to the $I$ resonance let us introduce the $\vec{J}$ vector

$$\vec{J} = \sum_{\alpha = e, x} \int_{0}^{\infty} dp \left( G_{\nu,\nu_{\alpha}}(p) \vec{P}_{\nu_{\alpha}}(p) \right) -G_{\bar{\nu},\bar{\nu}_{\alpha}}(p) \vec{P}_{\bar{\nu}_{\alpha}}(p).$$

(28)

We emphasize that, in a BNS merger scenario, one needs to include the geometrical factors in the definition of the collective vector, contrarily to what is usually done in the bulb model for supernovae (single-angle approximation), as e.g. in [41]. The reason is that here the geometrical factors differ for different flavors even when one employs the ansatz given by Eq.~(17). With definition (28) one can write the neutrino self-interaction term proportional to a unique vector $\vec{J}$, namely

$$\vec{B}_{\text{self}} = \sqrt{2} G_F \vec{J}.$$  

(29)

The evolution equation for $\vec{J}$ can be derived from the ones of $\vec{P}_{\nu_{\alpha}}$ (and $\vec{P}_{\bar{\nu}_{\alpha}}$) and using the explicit expressions of $\vec{B}$ ($\vec{B}_{\text{mat}}$). One finds

$$\partial_r \vec{J} = \vec{B}_{\text{mat}} \times \vec{J} + \vec{B}_0 \times \sum_{\alpha = e, x} \int_{0}^{\infty} dp \frac{\Delta m^2}{2 p} \left( G_{\nu,\nu_{\alpha}}(p) \vec{P}_{\nu_{\alpha}}(p) + G_{\bar{\nu},\bar{\nu}_{\alpha}}(p) \vec{P}_{\bar{\nu}_{\alpha}}(p) \right)$$

$$+ \sum_{\alpha = e, x} \int_{0}^{\infty} dp \left( \partial_r G_{\nu,\nu_{\alpha}}(p) \vec{P}_{\nu_{\alpha}}(p) - \partial_r G_{\bar{\nu},\bar{\nu}_{\alpha}}(p) \vec{P}_{\bar{\nu}_{\alpha}}(p) \right).$$

(30)
Let us assume now that, during the evolution, the modes all start along the z-axis, i.e. \( \vec{P}_{\nu_a} (r, p) \approx P_{\nu_a, z} (0, p) \hat{J} \) and stay aligned with the collective mode \( \vec{J} \) (similarly for antineutrinos). If neutrinos and antineutrinos of any momentum stay synchronized in flavor space during the propagation, the evolution equation for \( \vec{J} \) becomes

\[
\partial_r \vec{J} \approx \vec{B}_{\text{mat}} \times \vec{J} + \vec{B}_0 \times \vec{J} \int_0^\infty dp \frac{\Delta m^2}{2p} \left[ G_{\nu_a, \nu_a} (p) + G_{\bar{\nu}_a, \bar{\nu}_a} (p) - 2G_{\nu_a, \bar{\nu}_a} (p) \right] + \vec{J} \frac{\partial_\mu}{\sqrt{2G_F}},
\]

(31)

While the first two terms are ordinary oscillation terms, the last one is a damping term, taking into account that the norm of this collective mode decreases with time. This is due to the fact that the geometry of the problem is included in the definition of \( \vec{J} \). Note that such a decrease should not be interpreted as lepton number conservation violation, but as a neutrino density decrease along a given trajectory, due to the geometry. Let us characterize this decrease by multiplying the evolution equation (31) by \( \vec{J} \)

\[
\vec{J} \cdot \partial_r \vec{J} = \frac{1}{2} \partial_r J^2 \approx \left| \vec{J} \right| \frac{\partial_\mu}{\sqrt{2G_F}},
\]

(32)

which gives \( \left| \vec{J} (r) \right| \approx \frac{\mu(r)}{\sqrt{2G_F}} \). Plugging this expression in Eq. (31), one finds

\[
\omega_{\text{sync}} = \frac{\sqrt{2G_F} \Delta m^2 F_1 (0) F_3 (0)}{2\mu F_2 (0)} \left[ \frac{L_{\nu_a} G_{\nu_a}}{R_{\nu_a}^2 \langle E_{\nu_a} \rangle^2} + \frac{L_{\bar{\nu}_a} G_{\bar{\nu}_a}}{R_{\bar{\nu}_a}^2 \langle E_{\bar{\nu}_a} \rangle^2} - 2 \frac{L_{\nu_a} G_{\nu_a}}{R_{\nu_a}^2 \langle E_{\nu_a} \rangle^2} \right].
\]

(36)

b. Resonance condition In addition to a precession motion, the collective mode \( \vec{J} \) can also meet a MSW-like resonance condition \( B_{J,z} \approx 0 \), which requires

\[
\omega_{\text{sync}} \left( r_1 \right) e^{2\theta} = V_M \left( r_1 \right),
\]

(37)

where \( r_1 \) is the resonance location. From Eq. (36), it can be seen that \( \omega_{\text{sync}} \propto \frac{1}{\mu} \); in situations where the neutrino background dominates, the l.h.s. of Eq. (37) is often several of magnitude smaller than the r.h.s. However, in cases where the total matter potential \( V_M \) goes to zero, this resonance condition can be met. The reversed situation, in which the resonance condition is met because \( \mu \) goes to zero, has been already pointed out in [32].

Figure 4 shows the r.h.s. and the l.h.s. of Eq. (37), corresponding to the case of Figure 3. One can see that the synchronized MSW resonance condition given by Eq. (37) is met almost at location where \( V_M \) goes to zero, i.e. at the location of the I resonance, as can be seen from the conversion probabilities. Another example of synchronized I resonance is shown in Figure 5 with the neutrino self-interaction dominating over the matter potential. Significant conversion can be seen at 29 km, 40 km, 65 km et 78 km.
very small. Therefore, at the resonance, as order of magnitude and of opposite signs, making $\omega$ the fact that for this value of $\epsilon$, the cancellation of the adiabaticity parameter around the behaviours observed for the survival probabilities. Note that the location of the resonance, the adiabaticity parameter in the case of $\epsilon = 1 \times 10^{-5}$ is two order of magnitude smaller than the one for $\epsilon = 1 \times 10^{-4}$, consistently with the behaviors observed for the survival probabilities. Note that the cancellation of the adiabaticity parameter around the resonance in the case of $\epsilon = 1 \times 10^{-5}$ comes from the fact that for this value of $\epsilon$, the matter contribution and the $\omega_{\text{sync}}$ contribution in $B_{J,z}$ (Eq. 37) are of the same order of magnitude and of opposite signs, making $B_{J,z}$ very small. Therefore, at the resonance, as $B_{J,z}$ tends to 0, $\gamma \rightarrow \frac{B_{J,z}^2}{\sigma_c B_{J,z}}$ becomes much smaller at the same time.

d. Effect of the neutrino mass ordering The sign of $\omega_{\text{sync}}$ changes when going from normal to inverted mass ordering. However, due to the fact that the resonance location almost coincides with the location at which $V_M$ changes its sign, the mass ordering will have little impact on it. In our calculations we have found modifications of the resonance location smaller than 1 km between normal and inverted mass ordering. As for the adiabaticity parameter $\gamma$, it also depends on $\omega_{\text{sync}}$ and its derivative. Figure 5 shows the effect of neutrino mass ordering on the adiabaticity of flavor evolution for a case with $\delta \epsilon = -0.90$ and $\epsilon_0 = 1 \times 10^{-4}$ where the I resonance is located very close the neutrinosphere, at 5 km.

### B. NSI, the MNR and the I resonance

The occurrence of the MNR in BNS might impact $r$-process nucleosynthesis in neutrino-driven winds, as discussed in Ref. 10. Two kinds of MNR have been pointed out, either a symmetric one in which both neutrinos and antineutrinos convert [9], or a standard one where only neutrinos undergo flavor conversion [10]. The MNR phenomenon is due to a cancellation between the standard matter term Eq. (11) and the neutrino self-interaction Eq. (18). This occurs because of the excess of the antineutrino over the neutrino near the disk in the BNS context, compared to the supernova case, that gives a negative sign to the neutrino self-interaction potential $\mu$ [20]. However, Ref. [26] has shown the presence of NSI can trigger the MNR also in the supernova context. In our numerical investigations, we have observed various NSI effects on the flavor behaviours in presence of MNR. First the existence of NSI can modify the location of the MNR. Figure 9 shows that the cancellation between the matter and the neutrino self-interaction terms shifts from 10 km to 30 km when NSI are included. Moreover neutrino evolution turns from completely non-adiabatic to adiabatic, as the the survival probabilities show. By looking at the difference of the neutrino Hamiltonian diagonal elements, one can see that they keep being very small from 30 km to 80 km due to the non-linear feedback that matches the nonlinear neutrino self-interaction contribution to the matter potential [30].

Along numerous trajectories and sets of NSI parameters we have observed an intriguing interplay between the I resonance, synchronized or not, and the MNR. Figures 10 [11] and [12] furnish three examples of such behaviours. Figure 10 shows a combination of I resonance and MNR. There are two I resonances, the first at 5 km which is partially adiabatic, and the second at 21 km, which triggers a MNR between 20 km and 100 km, followed by a second one between 160 km and 240 km where the $\nu_e$ are converted while $\bar{\nu}_e$ are not [9]. Note that this is in opposition to what the MNR typically creates in the absence of NSI : indeed, without NSI, the MNR tends to lead to flavor conversions for neutrinos while for antineutrinos the evolution is generally non-

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Note that this corresponds to the same parameters as the ones of Figure 8 with a larger range shown.
FIG. 5. Left panel: Difference of the diagonal elements of the total neutrino Hamiltonian (solid line), matter potential $V_M$ (dashed line) Eq. (21) in presence of NSI contributions with $\delta \epsilon_0 = -0.90$ and $\epsilon_0 = 1 \times 10^{-3}$ and self-interaction oscillated potential (dotted line), as a function of distance from the emission point. The initial parameters are $x_0 = -35$ km, $z_0 = 25$ km, and $\theta_q = 50^\circ$. Middle and right panels: Survival probabilities for neutrinos (middle), antineutrinos (right). The NSI parameters are set to $\delta \epsilon_0 = -0.90$ and $\epsilon_0 = 1 \times 10^{-4}$ (solid lines) and $\epsilon_0 = 1 \times 10^{-5}$ (dotted lines).}

FIG. 6. Left panel: Difference of the diagonal elements of the total neutrino Hamiltonian, as a function of distance from the emission point. The initial parameters are $x_0 = 15$ km, $z_0 = 32$ km, and $\theta_q = 15^\circ$. Middle and right panels: Averaged survival probabilities for neutrinos (middle), antineutrinos (right). The NSI parameters are set to $\delta \epsilon_0 = -0.88$ and $\epsilon_0 = 1 \times 10^{-4}$ (solid lines) and $\epsilon_0 = 1 \times 10^{-5}$ (dotted lines).

FIG. 7. Adiabaticity parameter as the right-hand-side of Eq. (28), corresponding to Figure 6. The solid line corresponds to $\epsilon_0 = 1 \times 10^{-4}$, while the dashed line corresponds to $\epsilon_0 = 1 \times 10^{-5}$. The location of the I resonance is shown as a vertical dotted line.

takes place. Between 100 km and 125 km the difference of the diagonal elements stays very small, creating small conversions. Finally, at 144 km, an I resonance occurs. The third example of a combination of MNR and I resonances is given in Figure 12. This case in point is interesting as it shows four I resonances: the first, located around 2 km, being a standard one, completely adiabatic, and the other three being synchronized resonances. At 12 km, the second resonance is also very adiabatic, then the third, at 26 km creates only partial conversions. A fourth resonance occurs at 58 km and produces a short MNR-like cancellation between 60 km and 66 km, followed by a MNR between 96 km and 126 km. Notice again the peculiar behavior of this MNR, which creates conversions for antineutrinos while the evolution for neutrino is nonadiabatic.

or partially adiabatic. In Figure 11 an I resonance is located at 2 km, followed by a non-adiabatic MNR at 12 km. Then, between 60 km and 70 km MNR conversions
FIG. 8. Left panel: Difference of the diagonal elements of the total neutrino Hamiltonian, as a function of distance from the emission point, for normal (solid lines) and inverted (dotted lines) mass ordering. The initial parameters are $x_0 = -10$ km, $z_0 = 30$ km, and $\theta_q = 25^\circ$. Middle and right panels: Averaged survival probabilities for neutrinos (middle), antineutrinos (right). The NSI parameters are set to $\delta \epsilon^a = -0.90$ and $\epsilon_0 = 1 \times 10^{-4}$.

FIG. 9. Left panel: Difference of the diagonal elements of the total neutrino Hamiltonian, as a function of distance from the emission point, without NSI (dotted line) and with NSI parameters set to $\delta \epsilon^a = -0.70$ and $\epsilon_0 = 1 \times 10^{-4}$ (dotted line). The initial parameters are $x_0 = 12$ km, $z_0 = 27$ km, and $\theta_q = 40^\circ$. Middle and right panels: Averaged survival probabilities for neutrinos (middle), antineutrinos (right).

IV. DISCUSSION AND CONCLUSIONS

In order to assess the role of flavor evolution on nucleosynthesis in neutrino driven winds a self-consistent calculation of the electron fraction modification coupled to the flavor evolution should be performed, as e.g. the one performed in Ref. [12] in core-collapse supernovae. First steps in this direction are presented in Refs. [9, 10]. However the trajectory dependence on the abundances and investigations without the ansatz [17] need to be performed. Such studies go beyond the scope of the present work. Figure [13] shows the I resonance location according to Eq. (21) in the dimensional space. One can see that such resonance can occur close to the neutrinosphere and for a large set of NSI parameters. Obviously, for the cases where only the matter term matters, the resonance location would keep unchanged if the ansatz [17] is relaxed.

9 Note that in these calculations are not fully consistent since the feedback effect of the modified electron fraction on the probabilities is not included.

Using the at-equilibrium $Y_e$ as a reference, one would expect that the $Y_e$ value should be increased by the presence of I resonances since the $\nu_e$ and $\bar{\nu}_e$ conversion to $\nu_x$ and $\bar{\nu}_x$ respectively brings the former to have the average energies of the latter. However, the at-equilibrium $Y_e$ is certainly not a good reference for the conditions encountered very close to the neutrinosphere. Only a consistent calculation of $Y_e$ modification including the feedback on the probabilities and the full angular dependence of the neutrino emission would tell us how much flavor evolution impacts the electron fraction.

In the present work, we have investigated the role of nonstandard matter-neutrino interactions within 2$\nu$ flavor framework. In particular we have included the electron-tau couplings for which current bounds from scattering and oscillation experiments are still rather loose. By solving the mean-field Liouville Von-Neumann equations along a large ensemble of trajectories, we have uncovered aspects of NSI impact on flavor evolution and in particular on the I resonance and the MNR. First, we have shown the conditions for the I resonance are met in this kind of setting, based on detailed BNS simulations,
FIG. 10. Left panel: Difference of the diagonal elements of the total neutrino Hamiltonian (solid line), matter potential $V_M$ (dashed line) Eq. (21) in presence of NSI contributions with $\delta e^\nu = -0.90$ and $\epsilon_0 = 1 \times 10^{-4}$ and self-interaction oscillated potential (dotted line), as a function of distance from the emission point. The initial parameters are $x_0 = -10$ km, $z_0 = 30$ km, and $\theta_\theta = 25^\circ$. Middle and right panels: Survival probabilities for neutrinos (middle), antineutrinos (right). Different energies correspond to different colors, and the averaged probabilities (dotted line) are shown. The slight dependence on the energy is due to the fact that as the MNR occurs further away from the emission point, the difference between the diagonal elements becomes comparable to the vacuum term, which then plays a role.

FIG. 11. Left panel: Difference of the diagonal elements of the total neutrino Hamiltonian (solid line), matter potential $V_M$ (dashed line) Eq. (21) in presence of NSI contributions with $\delta e^\nu = -0.70$ and $\epsilon_0 = 1 \times 10^{-5}$ and self-interaction oscillated potential (dotted line), as a function of distance from the emission point. The initial parameters are $x_0 = -30$ km, $z_0 = 20$ km, and $\theta_\theta = 55^\circ$. Middle and right panels: Survival probabilities for neutrinos (middle), antineutrinos (right). Different energies corresponding to different colors as well as averaged probability (dotted line) are indistinguishable.

FIG. 12. Left panel: Difference of the diagonal elements of the total neutrino Hamiltonian (solid line), matter potential $V_M$ (dashed line) Eq. (21) in presence of NSI contributions with $\delta e^\nu = -0.90$ and $\epsilon_0 = 1 \times 10^{-4}$ and self-interaction oscillated potential (dotted line), as a function of distance from the emission point. The initial parameters are $x_0 = -30$ km, $z_0 = 20$ km, and $\theta_\theta = 55^\circ$. Middle and right panels: Survival probabilities for neutrinos (middle), antineutrinos (right). Different energies correspond to different colors, and the averaged probabilities (dotted line) are shown. The slight dependence on the energy is due to the fact that as the MNR occurs further away from the emission point, the difference between the diagonal elements becomes comparable to the vacuum term, which then plays a role.
when the matter term dominates over the self-interaction contribution to the neutrino Hamiltonian. Then, we have uncovered the role of the neutrino self-interaction term and shown that the I resonance can be a synchronized MSW effect: the self-interaction potential dominates over the matter one. The synchronized precession frequency, depending on the self-interaction potential, matches the matter one. The synchronized precession frequency, when the total matter term becomes very small. This mechanism has been dismissed in previous investigations. Note that in Ref. 32 a synchronized MSW effect is observed when, on the contrary, the self-interactions become very small. Second, for the

MNR we have shown that NSI modify little the resonance location while the adiabaticity can be significantly changed. Third we have shown complex situations where MNR, I and synchronized I combine producing intriguing flavor patterns.

The contribution of heavy elements nucleosynthesis in BNS is an open question. The discovery of gravitational waves 13, the determination of the BNS rate from the LIGO and Virgo collaborations and the kilonova observation 5 bring crucial information to the longstanding puzzle of the origin of r-process nuclei. To answer this question, one needs to assess the BNS rate as well as the amount of elements from each individual event. The kilonova observation constitutes an experimental proof that heavy elements are indeed produced in BNS. In such sites nucleosynthetic abundances can be produced in the dynamical ejecta and neutrino-driven winds (see e.g. 7 14). Which elements are produced in each needs to be assessed. In this respect it is necessary to determine if and under which conditions flavor evolution takes place as well as its influence on nucleosynthetic abundances. The present work provides insights to progress in this direction, also in presence of new physics as nonstandard interactions that might be discovered in the future.

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[1] M. E. Burbidge, G. R. Burbidge, W. A. Fowler and F. Hoyle, Rev. Mod. Phys. 29, 547 (1957).
[2] J. M. Lattimer, D. N. Schramm and L. Grossman, Nature 269, no. 5624, 116 (1977) doi:10.1038/269116a0
[3] Y. Z. Qian, J. Phys. G 41, 044002 (2014) arXiv:1310.4462 [astro-ph.SR].
[4] A. Arcones and F. K. Thielemann, J. Phys. G 40, 013201 (2013) arXiv:1207.2527 [astro-ph.SR].
[5] B. P. Abbott et al. [LIGO Scientific and Virgo Collaborations], Phys. Rev. Lett. 119, no. 16, 161101 (2017) doi:10.1103/PhysRevLett.119.161101 arXiv:1710.05832 [gr-qc].
[6] E. Pian et al., Nature doi:10.1038/nature24298 arXiv:1710.05858 [astro-ph.HE].
[7] D. Martin, A. Perego, A. Arcones, F. K. Thielemann, O. Korobkin and S. Rosswood, Astrophys. J. 813, no. 1, 2 (2015) arXiv:1506.05048.
[8] A. B. Balantekin and H. Yuksel, New J. Phys. 7, 51 (2005) doi:10.1088/1367-2630/7/1/051 astro-ph/0411159.
[9] A. Malkus, J. P. Kneller, G. C. McLaughlin and R. Surman, Phys. Rev. D 86, 085015 (2012) arXiv:1207.6648 [hep-ph].
[10] A. Malkus, G. C. McLaughlin and R. Surman, Phys. Rev. D 93, no. 4, 045021 (2016) arXiv:1507.00946.
[11] F. Capozzi, E. Di Valentino, E. Lisi, A. Marrone, A. Melchiorri and A. Palazzo, Phys. Rev. D 95, no. 9, 096014 (2017) arXiv:1703.04471 [hep-ph].
[12] H. Duan, G. M. Fuller and Y. Z. Qian, Ann. Rev. Nucl. Part. Sci. 60, 569 (2010) arXiv:1001.2799 [hep-ph].
[13] S. Chakraborty, R. S. Hansen, I. Izaguirre and G. Raffelt, JCAP 1601, no. 01, 028 (2016) arXiv:1507.07569 [hep-ph].
[14] M. S. Bilenky and A. Santamaria, hep-ph/9908272.
[15] C. Biggio, M. Blennow and E. Fernandez-Martinez, JHEP 0908, 090 (2009) arXiv:0907.0097 [hep-ph].
[16] T. Ohlsson, Rept. Prog. Phys. 76, 044201 (2013) arXiv:1209.2710 [hep-ph].
[17] S. Davidson, C. Pena-Garay, N. Rius and A. Santamaria, JHEP 0303, 011 (2003) hep-ph/0302093.
[18] D. Akimov et al. [COHERENT Collaboration], Science 357, no. 6356, 1123 (2017) doi:10.1126/science.aao0990 arXiv:1708.01294 [nucl-ex].
[19] F. S. Amanik and G. M. Fuller, Phys. Rev. D 75, 083008
(2007) [astro-ph/0606607].

[20] M. Blennow, A. Mirizzi and P. D. Serpico, Phys. Rev. D 78, 113004 (2008) [arXiv:0810.2297 [hep-ph]].

[21] D. Das, A. Dighe and M. Sen, JCAP 1705, no. 05, 051 (2017) [arXiv:1705.00468 [hep-ph]].

[22] H. Numakawa, Y. Z. Qian, A. Rossi and J. W. F. Valle, Phys. Rev. D 54, 4356 (1996) [hep-ph/9605301].

[23] L. Wolfenstein, Phys. Rev. D 17, 2369 (1978).

[24] A. Das, A. Dighe and M. Sen, JCAP 1705, no. 05, 051 (2017) [arXiv:1705.00468 [hep-ph]].

[25] H. Nunokawa, Y. Z. Qian, A. Rossi and J. W. F. Valle, Phys. Rev. D 54, 4356 (1996) [hep-ph/9605301].

[26] L. Wolfenstein, Phys. Rev. D 17, 2369 (1978).

[27] A. Esteban-Pretel, R. Tomas and J. W. F. Valle, Phys. Rev. D 94, no. 9, 093007 (2016) [arXiv:1605.04903 [hep-ph]].

[28] A. Esteban-Pretel, R. Tomas and J. W. F. Valle, Phys. Rev. D 81, 063003 (2010) [arXiv:0909.2196 [hep-ph]].

[29] M. R. Wu, H. Duan and Y. Z. Qian, Phys. Lett. B 752, 89 (2016) [arXiv:1509.08975].

[30] D. Väänänen and C. McLaughlin, Phys. Rev. D 93, no. 10, 105044 (2016) [arXiv:1510.00751 [hep-ph]].

[31] A. Chatelain and C. Volpe, Phys. Rev. D 95, no. 4, 043005 (2017) [arXiv:1611.01862 [hep-ph]].

[32] Y. L. Zhu, A. Perego and C. McLaughlin, Phys. Rev. D 94, no. 10, 105006 (2016) [arXiv:1607.04671 [hep-ph]].

[33] M. Frensel, M. R. Wu, C. Volpe and A. Perego, Phys. Rev. D 95, no. 2, 023011 (2017) [arXiv:1607.05938 [astro-ph.HE]].

[34] M. R. Wu and I. Tamborra, Phys. Rev. D 95, no. 10, 103007 (2017) [arXiv:1701.06580 [astro-ph.HE]].

[35] R. F. Sawyer, Phys. Rev. Lett. 116, no. 8, 081101 (2016) doi:10.1103/PhysRevLett.116.081101 [arXiv:1509.03325 [astro-ph.HE]].

[36] A. Perego, S. Rosswog, R. M. Cabez, O. Korobkin, R. Kppeli, A. Arcones and M. Liebendörfer, Mon. Not. Roy. Astron. Soc. 443, no. 4, 3134 (2014) [arXiv:1405.6730].

[37] J. Serreau and C. Volpe, Phys. Rev. D 90, no. 12, 125040 (2014) [arXiv:1409.3594].

[38] C. Volpe, D. Väänänen and C. Espinoza, Phys. Rev. D 87, no. 11, 113010 (2013) [arXiv:1302.2374 [hep-ph]].

[39] Z. Maki, M. Nakagawa and S. Sakata, Prog. Theor. Phys. 28, 870 (1962).

[40] M. Maltoni and A. Y. Smirnov, Eur. Phys. J. A 52, no. 4, 87 (2016) [arXiv:1507.05287 [hep-ph]].

[41] C. Patrignani et al. [Particle Data Group], Chin. Phys. C 40, no. 10, 100001 (2016). doi:10.1088/1674-1137/40/10/100001

[42] S. Pastor, G. G. Raffelt and D. V. Semikoz, Phys. Rev. D 65, 053011 (2002) doi:10.1103/PhysRevD.65.053011 [hep-ph/0109035].

[43] J. Tamborra, G. G. Raffelt, L. Hudepohl and H. T. Janka, JCAP 1201, 013 (2012) doi:10.1088/1475-7516/2012/01/013 [arXiv:1110.2104 [astro-ph.SR]].

[44] B. P. Abbott et al. [LIGO Scientific and Virgo Collaborations], Phys. Rev. Lett. 116, no. 6, 061102 (2016) doi:10.1103/PhysRevLett.116.061102 [arXiv:1602.03837 [gr-qc]].

[45] O. Just, A. Bauswein, R. A. Pulpillo, S. Goriely and H.-T. Janka, Mon. Not. Roy. Astron. Soc. 448, no. 1, 541 (2015) doi:10.1093/mnras/stv009 [arXiv:1406.2687 [astro-ph.SR]].