Imprints of neutrino-pair flavor conversions on nucleosynthesis in ejecta from neutron-star merger remnants

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The remnant of neutron star mergers is dense in neutrinos. By employing inputs from one hydrodynamical simulation of a binary neutron star merger remnant with a black hole of 3 $M_\odot$ in the center, dimensionless spin parameter 0.8 and an accretion torus of 0.3 $M_\odot$, the neutrino emission properties are investigated as the merger remnant evolves. Initially, the local number density of $\bar{\nu}_e$ is larger than that of $\nu_\ell$ everywhere above the remnant. Then, as the torus approaches self-regulated equilibrium, the local abundance of neutrinos overcomes that of antineutrinos in a funnel around the polar region. The region where the fast pairwise flavor conversions can occur shrinks accordingly as time evolves. Still, we find that fast flavor conversions do affect most of the neutrino-driven ejecta. Assuming that fast flavor conversions lead to flavor equilibration, a significant enhancement of nuclei with mass numbers $A > 130$ is found as well as a change of the lanthanide mass fraction by more than a factor of a thousand. Our findings hint towards a potentially relevant role of neutrino flavor oscillations for the prediction of the kilonova (macronova) light curves and motivate further work in this direction.

I. INTRODUCTION

Compact binary mergers originate from the coalescence of a neutron star (NS) with another NS or a black hole (BH). They have long been considered to be precursors of short gamma-ray bursts (sGRBs), main hosts of the nucleosynthesis of heavy elements [1, 2], and sources of gravitational waves (see [3] for a recent review and references therein). The radioactive decay of the synthesized neutron-rich nuclei has also been assumed to power electromagnetic transients, called kilonovae or macronovae [4–6] (see also, e.g., Refs. [7, 8] for recent reviews). These conjectures have recently been confirmed by the detection of the GW170817 event and its related electromagnetic counterparts [9–11] (see also Ref. [8] and references therein for the kilonova interpretation). Beyond the kilonova associated to the GW170817 event, a few more potential kilonova candidates have been identified through the infrared excess linked to the sGRB afterglows [12–15]. The upcoming increasing statistics in the detection of such events, both for what concerns gravitational waves and electromagnetic counterparts, will greatly improve our understanding of compact binary mergers and provide rich implications on their physics [16, 17].

In binary NS mergers or NS-BH mergers, neutrinos can be copiously generated due to the violent collision of the two NSs and the presence of the hot and dense post-merger massive NS or BH accretion disk. Similarly to core-collapse supernovae (CCSNe) [18], neutrinos play an important role in mergers as they dominate the cooling of the merger remnants, change the composition of the ejecta, and affect the nucleosynthesis outcome in the ejecta and eventually the electromagnetic light curves [19–24]. Neutrinos can also contribute to energizing sGRBs via pair annihilation above the BH accretion disk [1, 25–28].

The exploration of the role of neutrinos in the merger remnants is still preliminary due to the highly demanding computational requirements for fully three-dimensional, general-relativistic magnetohydrodynamical modeling with detailed neutrino transport. On the other hand, various simulations show a generic feature which is the protonization of the merger remnant (i.e., more $\bar{\nu}_e$’s than $\nu_\ell$’s are emitted). The protonization of the merger remnant has peculiar implications for the flavor conversion of neutrinos. For example, the so-called “matter-neutrino resonance” (MNR) [29–34] is expected to occur. The MNR is due to the cancellation of the matter potential describing the interactions of neutrinos with electrons ($\nu-e$) and the neutrino-neutrino potential ($\nu-\nu$). The MNR is typical of merger remnants; it does not occur in, e.g., CCSNe, unless physics beyond the Standard Model is involved [35, 36].

Given the nature of the $\nu-\nu$ interaction, the neutrino angular distribution may affect the overall flavor conversion. In particular, “fast” pairwise flavor conversions [37–39] strongly depend on the angular distribution of the electron neutrino lepton number (ELN) [40–42] and can...
develop on time scales of $(G_F |n_{\nu_e} - n_{\bar{\nu}_e}|)^{-1} \simeq O(10)$ cm at the vicinity of the neutrinosphere, with $G_F$ being the Fermi constant and $n_{\nu_e}$ ($n_{\bar{\nu}_e}$) the local $\nu_e$ ($\bar{\nu}_e$) number density. For any location above the neutrino emitting surface at a given time, fast conversions may occur if the ELN angular distribution has crossings, i.e., the ELN angular distribution is positive in a certain solid-angle range, but negative in another one. As the temporal and spatial scales during which fast conversions develop are small compared to the size of the astrophysical object, they may lead to flavor equilibration.

The role of the angular distribution of neutrinos in the context of flavor conversions in merger remnants has been investigated in Ref. [42] for the first time. A simple two-disk neutrino emission model, with the decoupling region of $\nu_e$ sitting within that of $\nu_x$, has been used. This model, however, was stationary and assumed uniform properties for neutrinos everywhere on the neutrino emitting surfaces. Moreover, the geometrical shape of the neutrino emitting surfaces was approximated with flat disks. Due to the remnant protonization and the emission geometry, the authors of Ref. [42] found that favorable conditions for flavor instabilities exist for any point above the $\nu_e$ emission surfaces.

In this work, we intend to further explore the role of fast conversions in merger remnants and investigate their potential impact on the synthesis of elements by adopting a realistic remnant configuration based on numerical simulations from Ref. [24]. Figure 1 sketches the typical geometry of the merger remnant and the different components of the matter ejected during and after the merger. The central compact object can be a BH or, in the case of a binary NS merger, a NS that may collapse to a BH at some later time. The outermost layer consists of the earliest dynamical ejecta which become unbound via tidal torques during the merger or through the violent collision of two NSs from the contact interface within $\lesssim 10$ ms after the merging. The total amount of these dynamical ejecta can be up to $10^{-3} - 10^{-1} M_\odot$ [43–48]. Neutrinos have a negligible role for the ejecta in the equatorial plane but may greatly influence the polar ejecta [45, 49].

Following the merger, a remnant accretion disk of up to $\sim 10^{-1} M_\odot$ surrounding the central massive NS or BH can form. Recent hydrodynamical simulations based on idealized initial conditions show that $\gtrsim 20\%$ of the initial disk mass can be further ejected via various mechanisms. In the first few hundred milliseconds, neutrinos coming from the hot and dense region of the inner disk and/or the central massive NS (prior to its collapse to a BH) can cause a neutrino-driven wind, dominantly around the polar region with a total mass of $\sim 10^{-3} M_\odot$ [21, 24]. On a longer time scale of a few seconds, $\sim 10^{-2} M_\odot$ can be further ejected by viscous heating and nuclear recombination [24, 50–54]. Both components are shown schematically in Fig. 1 in green and orange respectively.

One can see that the neutrino-driven ejecta may become the dominant component in the polar direction. Consequently, the potential change of the relative abundance of neutrinos of different flavors due to fast flavor conversions may eventually affect the nucleosynthesis outcome in that region.

Throughout the rest of the paper, we focus on the neutrino-driven ejecta in the post-merger phase. For the first time, we study the evolution of the neutrino emission properties by adopting inputs from one hydrodynamical simulation of a post-merger BH accretion disk. Based on model M3A8m3a5 of Ref. [24], we perform a flavor stability analysis for several time snapshots to pinpoint the eventual occurrence of fast pairwise conversions. In order to gauge the importance of neutrino flavor conversions in the remnant, we then quantify whether the flavor equipartition induced by fast pairwise conversions may be responsible for a non-negligible effect on the synthesis of heavy elements in the neutrino-driven ejecta.

The manuscript is organized as follows. In Sec. II, we first study and characterize the neutrino emission properties obtained in the hydrodynamical simulation of model M3A8m3a5 presented in Ref. [24]. In Sec. III, we perform a time-dependent stability analysis. In Sec. IV we discuss how the possible occurrence of flavor equipartition may affect the nucleosynthesis in the neutrino-driven ejecta from the merger remnants. Our conclusions and an outlook are presented in Sec V. Further details on the
flavor stability analysis are reported in Appendix A.

II. NEUTRINO EMISSION PROPERTIES

In this section, we describe the main features of the BH-torus evolution of model M3A8m3a5 with a central BH of 3 $M_\odot$, dimensionless BH spin parameter 0.8 and torus of 0.3 $M_\odot$. Note that we have chosen model M3A8m3a5 from the ones presented in Ref. [24], as this is the one with the largest fraction of neutrino-driven ejecta due to its large torus mass. Therefore, this case represents an optimistic scenario to explore the role of neutrino flavor conversions in the merger remnant. Special attention is dedicated to the evolution of the ELN and its angular distribution, the crucial quantity in the context of fast flavor conversions.

A. Evolution of binary neutron star merger remnants

The authors of Ref. [24] adopted a pseudo-Newtonian Artemova-Novikov gravitational potential [55] and energy-dependent neutrino transport scheme coupled to the Navier-Stokes equations with a Shakura-Sunyaev viscosity prescription to model the post-merger long-time evolution of the BH torus system. We refer the interested reader to Ref. [24] for more details on the simulation setup.

Depending on the initial condition, an accretion torus evolves during time scales ranging from tens of milliseconds to seconds. It loses mass by accreting onto the BH and by thermally and viscously driven outflows. The evolution of such a massive and dense torus can be divided mainly into three different stages, as shown in Fig. 2. Initially, the environment is dense enough to produce optically thick conditions for neutrinos, for which the latter are partially trapped and advected with the flow and neutrino cooling is less efficient. This first phase lasts for about $O(10)$ ms.

As the mass of the torus decreases and the density drops, the phase of a “neutrino-dominated accretion flow” begins. During this phase, neutrinos radiate away most of the gravitational energy that is converted into internal energy via viscous heating. As the mass, density and temperature of the torus continue to decrease, the neutrino production rate decreases until neutrino cooling becomes inefficient again at $t \approx 0.2 - 0.3$ s. This can be seen in the top panel of Fig. 2 which shows the evolution of the energy luminosities of both $\nu_e$ and $\bar{\nu}_e$ in the laboratory frame at a spherical radius of 500 km. The bottom panel of Fig. 2, displaying the ratio of the number luminosities of $\nu_e$ and $\bar{\nu}_e$, shows that during the entire evolution phase of the torus when neutrinos are efficiently produced, the torus on average continues to protonize (apart from the first $\sim 2$ ms during which the electron fraction in the densest parts of the torus settles from its initial value of 0.1 to a new, slightly lower weak equilibrium value).

The behavior of the average electron fraction of the torus and its temporal change (the latter being given by the ratio of number luminosities, $L_{N,\nu_e}/L_{N,\bar{\nu}_e}$, see bottom panel of Fig. 2) can be understood as follows: Accreting torus material at all times tends to achieve $\beta$-equilibrium (e.g. Ref. [56]). The electron fraction corresponding to this equilibrium (and therefore the actual $Y_e$ of the torus) remains rather low ($< 0.5$) during the first two evolutionary phases as a result of self-regulation between viscous heating and neutrino cooling into a state with semi-degenerate electrons [24, 54, 57]. However, as the torus becomes more and more diluted due to accretion onto the BH, the electron degeneracy is (on average) lifted, causing the gas to favor a more symmetric $\beta$-equilibrium regarding the abundance of neutrons and protons, i.e. higher electron fractions.

The particularly low values of $L_{N,\nu_e}/L_{N,\bar{\nu}_e}$ during the early, optically thick phase (indicating strong protoniza-
FIG. 3: BH-torus remnant properties for model M3A8m3a5 around the $z$-axis at 20, 35 and 50 ms (from left to right) as functions of $x$ and $z$ (assuming cylindrical symmetry around the $z$-axis). First row: Baryon mass density $\rho$. Second row: Temperature $T$. Third row: Degeneracy parameter $\mu_e/T$ of electrons with $\mu_e$ being the electron chemical potential. Fourth row: Relative electron neutrino lepton number $(n_{\nu_e} - n_{\bar{\nu}_e})/n_{\nu_e}$. The degeneracy parameter in the innermost torus region slightly increases over time as the torus evolves (see text for details). The BH-torus evolves from a configuration where $n_{\bar{\nu}_e} > n_{\nu_e}$ to a configuration where $n_{\bar{\nu}_e} < n_{\nu_e}$ around the polar axis. Also shown are the emitting surfaces of $n_{\nu_e}$ and $n_{\bar{\nu}_e}$ which are computed as described in Sec. II C and are marked in red and blue, respectively.

In the last stage, once neutrino cooling becomes inefficient, viscous angular momentum transport, together with viscous heating and nuclear recombination drive an inflation of the torus. This also leads to mass outflows which, however, are more massive near the equatorial plane. Such a viscously driven wind proceeds gradually and the matter becomes gravitationally unbound at large radii, with velocities lower than that of the supersonic neutrino-driven wind. More matter can be ejected during this phase.

For model M3A8m3a5 studied in this paper, the total mass of the neutrino-driven ejecta is approximately $1.47\times10^{-3} \, M_\odot$ compared to the total outflow mass which is about $66.2 \times 10^{-3} \, M_\odot$ [24]. Since we are interested in the effect of neutrino flavor conversions on the ejecta properties, throughout the paper, we will focus on the early outflow driven by neutrinos (i.e., $t \lesssim 60$ ms).
B. Electron neutrino lepton number and other remnant emission properties

The bottom panel of Fig. 2 shows that the relative rate of protonization changes as a function of time for \( t \lesssim 60 \text{ ms} \) (i.e., the time window relevant for the neutrino-driven ejecta). We, therefore, show in Fig. 3 the matter density \( \rho \), the temperature \( T \), the degeneracy parameter \( \mu_e/T \) with \( \mu_e \) being the electron chemical potential, and the ELN (\( \equiv n_{e^-} - n_{e^+} \)), as functions of \( x \) and \( z \) (assuming cylindrical symmetry) from top to bottom. Each quantity is shown for three selected snapshots at \( t = 20, 35 \) and \( 50 \text{ ms} \) (from left to right) to illustrate the evolution of the torus conditions. The surfaces where \( \nu_e \) and \( \bar{\nu}_e \) decouple are also shown in red and blue, respectively (see Sec. II C for more details).

As the torus continuously accretes onto the BH, both \( T \) and \( \rho \) decrease. Consequently, the size of both neutrino surfaces shrinks. However, the electron degeneracy \( \mu_e/T \) in the innermost part of the torus increases from \( \mu_e/T < 1 \) to \( \mu_e/T \sim 1 \) as the torus evolves and neutrino cooling becomes more efficient with decreasing optical depth. This increase of the electron degeneracy \(^1\) leads to a relatively larger ratio of the electron capture rate to the positron capture rate. Since most of the neutrinos ending up in the polar region are emitted from this inner region of the torus, this has consequences on the ELN above the neutrino surfaces. The bottom panels of Fig. 3 show that, at 20 ms, the whole region above the torus is characterized by \( n_{\nu_e} > n_{\bar{\nu}_e} \). The torus gradually evolves towards a configuration where \( n_{\nu_e} > n_{\bar{\nu}_e} \) in the polar region at later times. The main reason for having a \( \nu_e \) excess in the polar region is due to a geometrical effect. As the \( \nu_e \) surface with a conical shape is more extended than the \( \bar{\nu}_e \) surface with nearly the same half-opening angle (Fig. 3), more \( \nu_e \)'s are emitted towards the polar region than \( \bar{\nu}_e \)'s from their respective surfaces. This results in a \( \nu_e \) excess when the torus is only slightly protonizing at later times. Figure 4 provides more insight into this evolutionary effect as a consequence of the neutrino transport conditions around the torus. As matter flows towards the BH, it protonizes \((d\nu_e/dt > 0)\) in all of the near-surface regions of the torus at all times, while the high-density inner parts have achieved a steady state condition \((d\nu_e/dt \approx 0)\) or neutronize with a very low rate.

Nevertheless, it is only at early times that all the volume above the neutrino surfaces is dominated by the number densities and number fluxes of \( \bar{\nu}_e \) (Fig. 4, left panels). In contrast, at later times \((t \gtrsim 25 \text{ ms})\) a growing conical volume around the rotation axis develops an excess of \( \nu_e \) in number density and number flux. The reason is twofold: First, the decreasing rate of protonization with progressing evolution (compare left and right columns of Fig. 4) near the torus surface reduces the difference between the overall higher \( \bar{\nu}_e \) number flux compared to the \( \nu_e \) number flux, as well as locally at the neutrino surfaces. Second, the different emission geometry of the \( \nu_e \) and \( \bar{\nu}_e \) surfaces plays an increasingly more important role: because the neutrino surface of \( \nu_e \) is more extended, it irradiates the region around the rotation axis from a wider angle than the \( \bar{\nu}_e \) surface does. Both effects combined lead to the growing excess of \( \nu_e \) compared to \( \bar{\nu}_e \) around the \( z \)-axis.

In the BH-torus model M3A8m3a5, the transition between the two regimes from polar \( \bar{\nu}_e \) excess to polar \( \nu_e \) excess happens at about 25 ms. As we will see in Sec. III B, this has important consequences on the flavor conversions of neutrinos. We note here that such a transition is a generic feature seen in all BH-torus models in Ref. [24], while only the transition time depends on the model parameters.

C. Neutrino emission surfaces

Since the inner torus is dense enough to trap neutrinos, we can define the \( \nu_e \) and \( \bar{\nu}_e \) emitting surfaces to approximate the boundaries above which neutrinos can be considered as free-streaming particles, similarly to the neutrinosphere usually defined in CCSNe. Noticeably, the concept of emitting surfaces is nothing more than a formal definition; it is, however, useful for the flavor instability study done in this work. In CCSNe, the “neutrinosphere” is defined as the surface where the optical depth along the radial direction is about 2/3. However, in the torus case the optical depth becomes direction-dependent since the geometry is highly non-spherical. It is therefore difficult to unambiguously define a direction for calculating the neutrino emission surface.

We therefore adopt a simpler approach guided by CCSN simulations. At any location, the neutrino number flux \( \mathbf{F}_{\nu_\alpha} \) and number density \( n_{\nu_\alpha} \) of the flavor \( \nu_\alpha \) are given by (\( \hbar = c = 1 \))

\[
\mathbf{F}_{\nu_\alpha} = \frac{1}{(2\pi)^3} \int d\Omega dE \hat{\mathbf{p}} E^2 f_{\nu_\alpha}(\mathbf{p}),
\]

and

\[
n_{\nu_\alpha} = \frac{1}{(2\pi)^3} \int d\Omega dE E^2 f_{\nu_\alpha}(\mathbf{p}),
\]

where \( d\Omega \) is the differential solid angle, \( E \) and \( \mathbf{p} \) are the neutrino energy and momentum and \( f_{\nu_\alpha}(\mathbf{p}) \) is the neutrino phase-space distribution function.

By examining the CCSN simulations from Ref. [60], we found that the location of neutrinospheres obtained by using the definition based on the optical depth agrees well with the location where the flux factor \( F_{\nu_\alpha} \equiv |\mathbf{F}_{\nu_\alpha}|/n_{\nu_\alpha} \approx 1/3 \). In fact, when the neutrino distribution is isotropic

\(^1\) We note that a local increase of the electron degeneracy is not inconsistent with the previous statement that this quantity globally (i.e., averaged over the entire torus) decreases.
and neutrinos are trapped, the flux factor is zero, whereas when neutrinos start to free stream, the flux factor increases until it approaches unity at large distance from the emitting surface.

Figure 5 shows the neutrino number density of $\nu_e$ and $\bar{\nu}_e$ on their respective neutrino surfaces at 20, 35, and 50 ms as functions of $x$. The effect of the reduced protonization rate as the torus evolves and the different emission geometry of the $n_{\nu_e}$ and $n_{\bar{\nu}_e}$ surfaces (red and blue curves) lead to the growing excess of $\nu_e$ compared to $\bar{\nu}_e$ around the $z$-axis.

In this rest of the paper, we will take these $\nu_e$ and $\bar{\nu}_e$ surfaces as the inner boundaries where neutrinos are emitted and propagate freely afterwards. The impact of the above transition on the flavor instability will be discussed in the next section.

III. NEUTRINO FLAVOR CONVERSIONS IN COMPACT BINARY MERGERS

In this section, we introduce the dispersion relation (DR) in the neutrino flavor space. Our results on the flavor instabilities regarding fast pairwise conversions are also presented.

A. Dispersion relation in the flavor space

In the free-streaming regime, the equations of motion (EoMs) governing the flavor evolution of neutrinos are usually expressed in terms of the density matrix $\rho$, which encodes the flavor occupation numbers in the diagonal terms and the flavor correlations in the off-diagonal terms. The EoMs are

$$\{\partial_t + \mathbf{v} \cdot \nabla\} \rho = -i[H, \rho],$$

where $\mathbf{v} = (\sin \theta \cos \phi, \sin \theta \sin \phi, \cos \theta)$ is the velocity of an ultra-relativistic neutrino and its 4-vector is $v^\mu = (1, \mathbf{v})$. The Hamiltonian, $H$, consists of the vacuum term which takes into account the flavor mixing of neutrinos in vacuum [61], the matter term describing the coherent forward scattering of neutrinos with the matter background [62, 63] and the $\nu - \bar{\nu}$ term taking into account the interactions of neutrinos with their own background [64–66].

In this work, we are interested in investigating the role of fast flavor conversions [37–39]. Therefore, we will neglect the vacuum term in the Hamiltonian as well as the dependence of the neutrino energy and rely on a two-flavor framework $(\nu_e, \nu_x)$ where $\nu_x$ is a linear combination of the non-electron flavors [41, 42].

Expressing the neutrino density matrix as a function of the “flavor isospin” $\xi$ and the occupation number $f_{\nu_e}$ for the neutrino flavor $\nu$, $\rho = [(f_{\nu_e} + f_{\bar{\nu}_e}) + (f_{\nu_e} - f_{\bar{\nu}_e})\xi]/2$ ($\bar{\rho} = -(f_{\nu_e} + f_{\bar{\nu}_e}) + (f_{\nu_e} - f_{\bar{\nu}_e})\xi^*/2$) for neutrinos (antineutrinos) and introducing the metric $\eta^{\mu\nu} = \text{diag}(1, -1, -1, -1)$, the Hamiltonian for $\xi(\mathbf{v})$ can be written as

$$H = v^\mu \lambda_\mu \frac{\sigma_3}{2} + \int d\Omega' v'^\mu \rho' \xi(\mathbf{v}') g(\mathbf{v}') .$$

The term $v^\mu \lambda_\mu = \lambda_0 - \mathbf{v} \cdot \mathbf{\lambda}$, $\lambda = \lambda_0 \mathbf{v}_F$, where $\mathbf{v}_F$ is the
local fluid velocity, and $\lambda_0 = \sqrt{2}G_F n_e$ where $n_e$ is the net electron number density. In the following, we will work in the frame corotating with $\lambda$ and therefore take $v^\mu \lambda_\mu = \lambda_0$.

The neutrino angular distribution potential $g(\mathbf{v})$ per unit length per unit solid angle is proportional to the ELN angular distribution
\begin{equation}
 g(\mathbf{v}) = \sqrt{2}G_F (\Phi_{\nu_e} - \Phi_{\bar{\nu}_e}) ,
\end{equation}
where $\Phi_{\nu_e} = \frac{dn_{\nu_e}(\mathbf{v})}{d\Omega}$.

We are now interested in looking for non-null off-diagonal terms in the density matrix which could eventually arise, giving origin to fast conversions. Hence, we linearize the EoM [67, 68] and follow the evolution of the off-diagonal term $S$ in $\xi$, neglecting terms larger than $O(|S|)$:
\begin{equation}
\xi = \begin{pmatrix}
 1 & S \\
 S^* & -1
\end{pmatrix} .
\end{equation}

We make the ansatz that $S$ evolves as $S(\mathbf{v}, t, x) = Q(\mathbf{v}, \omega, \mathbf{k}) e^{-i(\omega t - \mathbf{k} \cdot \mathbf{x})}$. The EoM becomes [41]
\begin{equation}
 v_\mu s^\mu Q(\mathbf{v}, \omega, \mathbf{k}) + \int d\Omega' v_\mu' s^\mu g(\mathbf{v}') Q(\mathbf{v}', \omega, \mathbf{k}) = 0 .
\end{equation}

Here, we have introduced the 4-vector $s^\mu = (\omega - \lambda_0 - \epsilon_0, \mathbf{k} - \epsilon)$, $\epsilon_0 = \int d\Omega g(\mathbf{v})$ and $\epsilon = \int d\Omega g(\mathbf{v}) v_\mu$. From the structure of Eq. (7), one sees that the eigenfunction satisfying this equation is given by $Q(\mathbf{v}, \omega, \mathbf{k}) \propto (v_\mu a^\mu)/(v_\mu s^\mu)$, with $a^\mu$ being the coefficient of the eigenfunction solution. A non-trivial solution for the $Q_\mu(\omega, \mathbf{k})$ mode exists for non-zero $a^\mu$, if
\begin{equation}
 \det[\Pi_{\mu\nu}(\omega, \mathbf{k})] = 0 ,
\end{equation}
with $\Pi_{\mu\nu}(\omega, \mathbf{k}) = \eta_{\mu\nu} + \int d\Omega' v_\mu v_\nu g(\mathbf{v})/(v_\mu s^\mu)$. The above equation corresponds to defining a DR in the flavor space.

If the solutions of Eq. (8) are real, any initial perturbations in the flavor space do not grow. However, if any complex solution in $(\omega, \mathbf{k})$ satisfies the DR equation, then an exponentially growing instability may occur and lead to flavor conversions.

In what follows, we will look for temporal instabilities, mainly originating from crossings in the ELN angular distribution, i.e. look for complex solutions of $\omega$ with a given $\mathbf{k}$ satisfying Eq. (8) [41]. The growth rate of the flavor instability will be given by $|\text{Im}(\omega)|$. We neglect the occurrence of spatial instabilities (occurring for $|\text{Im}(k_i)| \neq 0$). In fact, the authors of Ref. [42] found that the latter should cover a much smaller spatial region than the temporal instabilities. Also, the authors of Ref. [69] recently concluded that non-zero $|\text{Im}(k_i)|$ alone might not lead to an exponentially growing instability.

B. Instabilities in the flavor space

The authors of Ref. [42], by approximating the merger remnant with a two-disk model, found that crossings in the ELN are present everywhere above the neutrino emitting surface. We now intend to verify whether such conclusions still hold within a realistic torus configuration evolving in time. To this purpose, we investigate the DR [Eq. (8)] assuming neutrinos are emitted from their respective decoupling surfaces defined in Sec. II C.

Due to the approximate neutrino transport still adopted in merger simulations, detailed neutrino angular distributions at decoupling cannot be extracted. Hence, an assumption needs to be made in order to obtain $g(\mathbf{v})$ above the neutrino surfaces.

Our neutrino surfaces have been defined as the surfaces where the flux factor $j_{\nu_e} = 1/3$ (see Sec. II C). One simple way to parametrize the local neutrino angular distribution on the surface, that is consistent with this definition, is to assume that the angular distribution grows linearly in $\cos \theta'$, where $\theta'$ is the angle with respect to the normal direction of the emission surface:
\[ \Phi_{\nu_e,\bar{\nu}_e}(\cos \theta') = \frac{n_{\nu_e,\bar{\nu}_e}}{4\pi} \left( 1 + \cos \theta' \right). \]  

One can easily verify that this angular distribution directly leads to \( j_{\nu_e,\bar{\nu}_e} = 1/3 \) while respecting the torus emission geometry. At any location above the \( \nu_e \) surface, we can then calculate the corresponding neutrino angular distribution potential \( g(\nu) \), by applying the ray-tracing method (see e.g., Appendix A in Ref. [42]) \(^2\).

Figure 6 shows the resulting ELN distribution as a function of \( \cos \theta \) and \( \phi \) (\( \theta \) is the angle relative to the \( z \)-direction, and \( \phi \) is the angle relative to the \( x \)-direction) in representative locations close to the inner [panels (a), (d), (g)], middle [panels (b), (e), (h)], and outer [panels (c), (f), (i)] regions above the \( \nu_e \) surface, using the procedure described above for the model M3A8m3a5 at 20, 35 and 50 ms. The red and blue shaded areas distinguish between regions where the ELN potential is positive and negative, respectively. The angular space where no neutrinos arrive from the emitting surfaces are left in white.

One sees from Fig. 6 that, as the torus protonizes less, the stronger \( \nu_e \) emission from the inner torus leads to a smaller solid angle where \( g(\nu) < 0 \) for the locations at the inner region. In particular, at later times, e.g. \( t = 50 \) ms, the ELN crossing in the inner region vanishes entirely [see panel (g)]. On the other hand, due to the persistently larger \( \bar{\nu}_e \) emission in the outer torus, the ELN crossing still occurs for locations in the middle and outer parts above the \( \nu_e \) surface.

We note here that, different from Ref. [42] where it was assumed that neutrino number densities were constant across their neutrino emitting surfaces, \( n_{\nu_e} \) and \( n_{\bar{\nu}_e} \) in this work are location dependent (see Fig. 5). Therefore, \( g(\nu) \) is in general not uniform. The color shading in

\( ^2 \) We here have neglected the neutrino ray bending effect due to general relativity. However, this effect should be minor in most of the regions, except those immediately next to the BH.
Fig. 6 is meant to only illustrate the structure of the ELN crossing.

One should also expect neutrinos to stream in the negative \( \cos \theta \) direction for locations above the \( \nu_e \) emitting surface because of projection effects due to the toroidal shape of the remnant. Moreover, a non-zero neutrino distribution in the negative \( \cos \theta \) direction should also be expected due to neutrino scattering that results in the neutrino halo effect discussed in Refs. [70, 71].

Similarly to the conclusion in Refs. [41, 42], complex solutions of the DR for a given \( k \) (i.e. the temporal flavor instabilities) exist whenever there is an ELN crossing. For the purpose of illustration, we show in Fig. 7 the complex part of the DR solution \( |\text{Im}(\omega)|/\mu \) for the spatially homogeneous mode \( k = 0 \) for locations above the \( \nu_e \) surface at \( t = 20, 35, 50 \) ms, where \( \mu = \sqrt{2} G_F n_{\nu_e} \).

The full solution of the DR for \( k = (0, 0, k_z) \) at locations corresponding to the ELN distribution shown in Fig. 6 is provided in the Appendix A.

Figure 7 shows that fast conversions occur everywhere above the \( \nu_e \) surface at \( t = 20 \) ms. This can easily be understood by looking at Fig. 6 where for \( t = 20 \) ms crossings in the ELN distribution appear for any point above the torus when the torus is strongly protonizing.

At later times (\( t > 30 \) ms), as the torus protonizes less and the local \( \nu_e \) abundance starts to become larger and even dominates the \( \bar{\nu}_e \) one around the polar region, we see that the unstable region of \( k = 0 \) shrinks, particularly in the region close to the pole and immediately above the middle part of the \( \nu_e \) surface. At 50 ms, when the ELN crossing completely disappears in the funnel region near the polar axis, the temporal instability is suppressed entirely. However, the local excess of \( \bar{\nu}_e \) with respect to \( \nu_e \) in the outer part of the disk still allows a large region for the flavor instability to exist for \( t = 50 \) ms. We also note that \( |\text{Im}(\omega)|/\mu \) becomes smaller at later times. We note that the growth rates of flavor instability shown in Fig. 7 range from 10 ns\(^{-1} \) \( \lesssim |\text{Im}(\omega)| \lesssim 1 \) \( \mu s^{-1} \).

Our results confirm the findings of Ref. [42] and conclude that favorable conditions for fast flavor conversions exist for the M3A8m3a5 torus. As discussed in Refs. [37–39], the fact that fast pairwise conversions could develop on time scales of ns to \( \mu s \), much smaller than the typical dynamical time scale of the system of \( \gtrsim \) ms, means that neutrinos of different flavors could potentially reach flavor equilibration and share the same properties.

As we will discuss in the next section, nearly all the neutrino-driven trajectories are affected by neutrinos that cross the instability regions. Hence, we will work under the assumption that flavor equilibration occurs because of pairwise conversions at \( t \leq 50 \) ms to investigate the potential role of neutrino flavor conversions in nucleosynthesis, instead of solving the full neutrino quantum kinetic equations in the non-linear regime. In fact new numerical tools need to be developed in order to incorporate fast pairwise conversions in the neutrino transport self-consistently. Our preliminary analysis here should serve as a test study to see whether more in-depth work on the modeling of pairwise neutrino conversions in binary neutron star merger remnants is needed.

The stability analysis set forth in this section has been developed within a two neutrino flavor \((\nu_e, \nu_x)\) framework, as often adopted in the investigation of \( \nu-\nu \) interactions; see e.g., references in Ref. [18]. In the following, we will generalize our conclusions to a full three flavor framework. As discussed above, new numerical tools are needed to exactly estimate the expected flavor outcome; however it is conceivable that if a flavor instability develops within extremely small time scales this may lead to a full mixing of all three flavors. As a consequence, in the following we will assume that flavor equilibration is reached for all three flavors

\[
 f_{\nu_e} = f_{\nu_x} = f_{\nu_x} = \frac{f_{\nu_0}}{3},
\]
for neutrinos that cross the unstable region shown in Fig. 7. In the above equation (10), \( f_{\nu_e} \) is the \( \nu_e \) phase-space distribution function before flavor conversions occur. An analogous relation is applied for antineutrinos.

IV. IMPACT OF FLAVOR EQUILIBRATION ON THE NUCLEOSYNTHESIS IN THE NEUTRINO-DRIVEN WIND EJECTA

After briefly introducing the physics of heavy element nucleosynthesis in the neutrino-driven wind in merger remnants, in this section we explore the role of flavor equilibration on the nucleosynthesis outcome of the neutrino-driven ejecta.

A. Neutrino driven wind in neutron star merger remnants

Heavy elements in the neutrino-driven wind of the post-merger BH-torus remnant are produced in a way similar to the CCSN neutrino-driven wind: matter recombines from free nucleons to form heavy nuclei as the ejecta expand and cool. The detailed calculation of the nucleosynthesis process requires solving a large set of equations describing the nuclear reaction network connecting different nuclei. In this work, we use the established nuclear reaction network suitable for the r-process nucleosynthesis calculation based on the nuclear physics inputs of Ref. [72]. It contains 7360 nuclei and all relevant nuclear reactions.

Despite the complicated nuclear physics needed to model the quantitative nucleosynthesis results, there are a few key quantities that determine the qualitative outcome (see e.g., Refs. [73, 74]), namely the entropy, the ejecta expansion time scale and, most importantly, the electron abundance fraction per nucleon

\[
Y_e = \frac{N_e}{N_p + N_n} = X_p + \sum_{A > 2} \frac{Z_A}{A} X_A ,
\]

where \( N_e \) (\( N_p, N_n \)) is the net electron (proton, neutron respectively) number density. \( X_p \) and \( X_A \) are the mass fractions of free protons and nucleons with charge numbers \( Z_A \geq 2 \).

In the early phases of the ejecta expansion when the temperature \( T \gg 1 \) MeV, matter mostly consists of free protons and neutrons, and the evolution of \( Y_e \) is then set by the \( \beta \)-interactions of neutrinos with free neutrons and protons:

\[
\nu_e + n \leftrightarrow p + e^- \quad \text{and} \quad \bar{\nu}_e + p \leftrightarrow n + e^+ .
\]

Therefore, the evolution of \( Y_e \) during this phase can be approximated as

\[
\frac{dY_e}{dt} \approx (\lambda_{\nu_e} + \lambda_{e^+}) Y_{n,f} - (\lambda_{\bar{\nu}_e} + \lambda_{e^-}) Y_{p,f} ,
\]

with \( \lambda_i \) being the reaction rates and \( Y_{n/p,f} \approx X_{n/p} \) the abundances of free nucleons.

When the temperature drops to \( T \gtrsim 1 \) MeV before nucleons recombine to \( ^4 \)He, the \( e^\pm \) capture rates \( (\lambda_{\nu_e - e^+} \propto T^5) \) become much smaller than the neutrino absorption rates \( (\lambda_{\nu,e}) \) and can later on be ignored. Moreover, when both \( \lambda_{\nu_e} \) and \( \lambda_{\bar{\nu}_e} \) become smaller than the inverse of the radial expansion dynamical time scale of the ejecta, \( \tau_{\text{dyn}} \approx r_{\text{ej},r} / v_{\text{ej},r} \), where \( v_{\text{ej},r} \) and \( r_{\text{ej},r} \) are the radial velocity and the radius for each given ejecta trajectory, one can define this time as the weak-interaction freeze-out time, \( t_{\text{FO}} \). At \( t_{\text{FO}} \), \( Y_e \) of the trajectory roughly approaches an asymptotic value \( Y_{\text{e asym}} \), until much later when the beta-decay of r-process nuclei sets in to further raise \( Y_e \) (see Fig. 9 for a few examples). Note that at \( t_{\text{FO}} \), all the neutrino driven ejecta roughly expand along the radial directions.

Another quantity which is relevant for the subsequent discussion is the so-called equilibrium electron fraction, \( Y_{e}^{\text{eq}} \), defined by

\[
Y_{e}^{\text{eq}} = \frac{\lambda_{\nu_e}}{\lambda_{\nu_e} + \lambda_{\bar{\nu}_e}} .
\]

When the neutrino irradiation is very strong, \( Y_e \) may reach \( Y_{e}^{\text{eq}} \) before the freeze-out. Equation (14) can be easily derived from Eq. (13) by assuming \( dY_e/dt = 0 \), neglecting \( \lambda_{\nu_e - e^+, \nu e} \), and taking \( Y_{n,f} = 1-Y_{p,f} = 1-Y_e \) [73]. In the typical CCSN neutrino-driven wind, this condition generally holds as the ejecta overcome the gravitational potential of the proto-neutron star by neutrino energy deposition. However, in the next section, we will see that it is not generally true for the neutrino-driven wind from post-merger BH-torus remnants, as matter is more loosely bound in this case.

The amount of the neutrino-driven ejecta for the M3A8s3a5 model is shown as a function of \( t_{\text{FO}} \) in the left panel of Fig. 8.\(^3\) Note that we have used for \( t_{\text{FO}} \) the same time coordinate as from the hydrodynamical simulation of model M3A8s3a5. The trajectories for all the neutrino-driven ejecta in the \( (x, z) \) plane are plotted in the right panel of Fig. 8. First, one sees that most of the neutrino-driven ejecta have the freeze-out time \( t_{\text{FO}} \approx 15-40 \) ms. Second, from the color coding, one can see that the trajectories ejected at early times originate mainly from regions next to the polar axis while later ejecta come from the outer edges of the torus next to the equatorial plane.

As the flavor instability exists at any point above the \( \nu_e \) surface at early times, and the outer part above the torus at later times, most of the neutrino-driven ejecta with \( t_{\text{FO}} \lesssim 50 \) ms will be influenced by neutrinos that stream through the unstable regions (see Fig. 7 and Sec. IIIIB for

\(^3\) For a comparison of the mass distribution histograms between the neutrino-driven ejecta and the viscously-driven ones, we refer the reader to Fig. 9 of Ref. [24].
comparison). As a consequence, we will assume flavor equilibration [see Eq. (10)] happens for neutrino fluxes on ejecta trajectories at $t \leq 50$ ms in the following.

**B. Impact of flavor equilibration on the element production**

We now explore the impact of flavor equilibration on the nucleosynthesis outcome of the neutrino-driven ejecta. Since the muon and tau (anti)neutrinos are produced only in the very innermost and dense regions of the torus, their luminosities are about 10 times lower than the ones of $\nu_e$ and $\bar{\nu}_e$ [28, 75, 76]. We here neglect the non-electron flavors and perform nucleosynthesis calculations by assuming that, when flavor equilibration occurs [see Eq. 10], both the $\nu_e$ and $\bar{\nu}_e$ capture rates on nucleons are reduced to 1/3 of their original values without oscillations.

To discuss the impact of flavor equilibration on the electron fraction $Y_e$, we first examine three representative trajectories with $t_{FO} = 16$, 25, and 31 ms. The top panel of Fig. 9 shows the selected trajectories in the ($x, z$) plane. The bottom panel of Fig. 9 shows the evolution of $Y_e$ with and without flavor equilibration (thick and thin lines respectively). The earlier ejecta with $t_{FO} = 16$ ms originate from the region close to the pole and, therefore, are exposed to stronger neutrino fluence. As a result, despite the strong reduction of the neutrino absorption rates, $Y_e(t_{FO}) \approx Y_{e}^{eq}$. This explains why a reduction of the neutrino rates due to flavor conversions has only little effect on the $Y_e$ evolution. For the later ejecta, such as the ones with $t_{FO} = 25$ and 31 ms, the asymptotic value of $Y_e$, $Y_{e}^{asym}$, never reaches $Y_{e}^{eq}$ even in the case without oscillations. Thus, the reduction of the neutrino capture rates due to flavor equilibration dramatically lowers the asymptotic value of $Y_e$ ($Y_{e}^{asym}$) from $\sim 0.41$ and $0.34$ to $\sim 0.3$ and $0.23$, respectively. Note that $Y_e$ for the $t_{FO} = 25$ ms trajectory shows a slight rise at $t \sim 80$ ms; this is due to the $\beta$-decays of neutron-rich nuclei during and after the r-process.

Figure 10 shows the final nucleosynthesis outcome of these three trajectories. The mass fraction $X(A)$ is plotted as a function of the nuclear mass number $A$. As a consequence of the $Y_e$ evolution shown in Fig. 9, there is only a small change in the nucleosynthesis pattern of the earliest trajectory with $t_{FO} = 16$ ms, while the produced heavy nuclei in the later ejecta are shifted from peaking around $A \sim 80$ to $A \gtrsim 130$, and even reaching the third peak $A \sim 195$ for the case with $t_{FO} = 31$ ms.

Figure 11 shows the ejecta masses as a function of the asymptotic $Y_{e}^{asym}$ as well as the mass fraction for all the neutrino-driven trajectories shown in the right panel of Fig. 8 for the cases with and without flavor equilibration. As evident from the top panel of Fig. 11, the eventual occurrence of flavor equipartition greatly changes the $Y_{e}^{asym}$ distribution of the ejecta, from uniformly distributed in the range $Y_e \in [0.35, 0.5]$ to being peaked around $Y_e \sim 0.25$ with a tail distribution reaching $\sim 0.5$.

The overall production of heavy elements is therefore shifted from abundance peaks around $A \sim 80$ to $A \sim 130$, as shown in the bottom panel of Fig. 11. In addition, the production of nuclei above $A \sim 130$ is enhanced by more than a factor of 1000.\(^4\)

\(^4\) Note that for the no-oscillation case, the production of nuclei is
Our explorative study suggests that fast pairwise conversions may indeed greatly affect heavy element production in the neutrino-driven wind of the merger remnant and strongly justifies further work in this direction. In particular, the enhancement of the production of lanthanides and the third-peak nuclei can be substantial. This can potentially lead to interesting observational consequences on the kilonova (macronova) lightcurve, if the neutrino-driven wind dominates the polar ejecta. For example, observations of the kilonova associated to the GW170817 event suggest blue (high $Y_e$) ejecta in the polar direction. Our results may support the interpretation that this observation points to a massive NS remnant that was stable for some time before collapsing to BH with some delay [77]. In fact, the specific spectrum of the electromagnetic signal may sensitively depend on the fraction of lanthanides [78, 79]. If this should be the case, an increasing number of face-on observations of the kilonova lightcurves along with theoretical improvements in the modeling of binary mergers may also be able to put indirect constraints on fast flavor conversions and neutrinos.

We here assumed that flavor equilibration occurs for any time $t \leq 50$ ms. Given the change of sign of the ELN, our assumption may seem extreme as the regions above the torus where the system is unstable shrink for $t > 30$ ms (see Fig. 7). However, we stress again that the neutrino-dominated trajectories although ejected at different times are always influenced by the neutrinos crossing the instability regions. On the other hand, flavor equilibration will also reduce the heating and the amount of neutrino-driven ejecta. Matter can take longer to be ejected and the real asymptotic $Y_{e\text{ asym}}$ would likely be sitting in between our results with and without oscillations. Given the potential major implications of neutrino conversions on heavy element synthesis, further exploration beyond the scope of this work is definitely needed to fully understand the role of neutrino flavor conversions in post-merger nucleosynthesis.

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slightly different with respect to Fig. 13 of Ref. [24]. This is due to the fact that we ignore ejecta with $t_{FO} < 10$ ms from the torus in this work, as the torus is still going through an artificially high $\nu_e$ emission for $t_{FO} < 10$ ms [24]. By including the first 10 ms of the neutrino-driven ejecta, we can indeed reproduce the results in Ref. [24], except for small differences due to different nuclear physics inputs.
from being negative ($n_{\nu_e} < n_{\bar{\nu}_e}$) everywhere to exhibiting a $\nu_e$ excess in the funnel region around the polar axis after $\sim 30$ ms. The fact that $\nu_e$ dominate with respect to $\bar{\nu}_e$ for $t > 30$ ms, implies the need for a time-dependent analysis of the flavor conversion phenomenology, which has never been addressed before.

By performing a flavor stability analysis for different time snapshots of the remnant, we found that favorable conditions for fast conversions exist for every point above the $\nu_e$-emitting surface for $t < 30$ ms. As $n_{\nu_e}$ starts to become larger than $n_{\bar{\nu}_e}$ around the polar axis, the region where the flavor instability can develop shrinks, but it persists in the outer part above the torus where crossings in the electron neutrino lepton number distributions occur.

Under the assumption that fast pairwise conversions lead to full flavor equilibration, we further investigated the impact of the reduced neutrino absorption rates on nucleosynthesis in the neutrino-driven ejecta. The $Y_e$ of the neutrino-driven outflow can be largely reduced. Consequently, the production of nuclei with mass numbers larger than 130 can be enhanced by more than a factor of 1000 with respect to the case where flavor conversions are neglected. The enhanced production of lanthanides may also greatly change the opacity of the neutrino-driven ejecta and thereby affect the resulting kilonova lightcurves.

In conclusion, our findings hint towards a relevant role of neutrino flavor conversions in binary neutron star merger remnants. The details of our findings should be taken with caution as our work was only meant to be exploratory and future work needs to be done to address the caveats adopted in this study.

One of the largest caveats is that we did not compute the exact flavor distribution due to fast pairwise conversions, but we assumed that flavor equilibration is reached given the temporal and spatial scale on which flavor instabilities are supposed to develop. Further numerical tools tackling fast pairwise conversions in a highly asymmetric environment need to be developed. Furthermore, if fast conversions happen so close to the neutrino decoupling surfaces, a self-consistent feedback effect of flavor conversions on the ejecta composition and the merger dynamics needs to be carefully implemented. The role of residual non-forward scatterings between neutrinos and matter above the torus also needs to be examined.

We here only analyzed in detail the case of a remnant with a central black hole. However, we note that different BH-torus configurations in Ref. [24] all share the same qualitative behavior of protonization and neutrino emission. Therefore, fast flavor conversions should occur in any BH-torus remnants and have similar impact on the nucleosynthesis of heavy elements in the associated neutrino-driven wind. Remnants with a massive neutron star in the center should also be studied in this respect as the amount of neutrino-driven ejecta can be much larger in this case [51, 77]. We expect that the

V. CONCLUSIONS AND OUTLOOK

Binary neutron star mergers are neutrino-dense environments and likely sites for the rapid-neutron capture process. By adopting inputs from the hydrodynamical simulation of a binary neutron star merger with a black hole of $3M_\odot$, dimensionless spin parameter $0.8$ and an accretion torus with mass of $0.3M_\odot$ (M3A8m3a5) [24], for the first time, we have studied the neutrino emission properties as a function of time, investigated the conditions under which fast pairwise conversions should develop in this environment, and examined the impact of flavor equilibration on the nucleosynthesis of heavy elements in the neutrino-driven ejecta.

During the first 50 ms of the neutrino-dominated accretion, when neutrinos are efficiently emitted, the torus strongly protonizes initially. Then, it gradually approaches a self-regulated semi-degenerate state with a lower protonization rate. Together with the geometry of the torus, this changes the electron neutrino lepton number distribution above the neutrino emitting surfaces
condition for fast flavor conversions to occur above the remnant NS-disk system should also exist, based on the toy models studied in Ref. [42]. The impact of flavor equilibration in such systems needs to be carefully examined because substantial amount of non-electron flavor neutrinos can be emitted from the central massive neutron star. Moreover, the condition of neutrino flavor conversions during the dynamical merger phase should also be considered in the future as several recent studies have shown that neutrinos can play an important role in driving the polar ejecta in the case of binary neutron star mergers [22, 49]. Together with the theoretical improvements, future gravitational-wave follow-up kilonova observations like the recent detection of GW170817 will offer unique opportunities to shed light on the role of neutrinos and their flavor conversions in compact binary mergers.

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Appendix A: Fast Pairwise Conversions: Instability regions

In this appendix, we provide more details on the solution of the dispersion relation (Eq. (8)) at various locations above the merger remnant at different times. Figure 12 shows the solutions for \( k = (0, 0, k_z) \) for the model M3A8n3a5 at the same locations and times for given ELN distributions in Fig. 6. The dash-dotted lines show the real part of the solutions while the continuous lines represent the imaginary part. One can see that at 20 ms, unstable solutions always exist for a wide range of \( k_z/\mu \) and shift from mainly at \( k_z/\mu \lesssim 0 \) to \( k_z/\mu \gtrsim 0 \) as one moves from the inner part above the remnant to the outer part [panels (a)–(c)]. This is similar to what was found in Ref. [42]. At later times, \( t = 35 \) and 50 ms, instabilities in general still exist [see panels (d)–(i)], however, the system becomes less unstable, particularly in regions close to the polar axis as the ELN angular distribution becomes dominated by neutrinos (see Fig. 6); see e.g., panels (d), (g), and (h). This is related to the fact that the neutrino local density starts to exceed the antineutrino one and the BH-torus remnant approaches the self-regulated equilibrium discussed in Sec. II.
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FIG. 12: Dispersion relation of $k = (0, 0, k_z)$ for the ELN distribution of $t = 20, 50$ ms of the model M3A8m3a5. For the complex $\omega$ solutions that lead to flavor instability, $\Re(\omega)$ are plotted in red dash-dotted curves and $\Re(\omega) \pm \Im(\omega)$ are plotted with red solid curves. The gray region indicates the zone of avoidance for real $(\omega, k_z)$. The system is unstable to fast conversions for a large range of $k_z/\mu_0$ and it is therefore unavoidable that fast pairwise conversions would occur.
