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On the importance of viscous dissipation and heat conduction in binary neutron-star mergers

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Inferring the properties of dense matter is one of the most exciting prospects from the measurement of gravitational waves from neutron star mergers. However, it requires reliable numerical simulations that incorporate viscous dissipation and energy transport as these can play a significant role in the survival time of the post-merger object. We revisit this assumption by exploring the impact of viscosity and thermal transport after merger, exploiting results of simulations. These have seen enormous progress [9–13] and found that for not too massive or too asymmetric systems, the post-merger object is metastable to gravitational collapse over tens of milliseconds. The inner region of this object, ~10 km across, can reach several times nuclear-matter saturation (number) density $n_0 \approx 0.16\text{fm}^{-3}$ and temperatures of tens of MeV.

The details of the complicated post-merger phase depend on the mass of the system, the equation of state (EOS), and the magnetic fields that develop after the merger [4, 5]. Quite generically, unless it collapses promptly to a black hole [10], the binary-merger product will oscillate in modes that leave the magnetic fields that develop after the merger [4, 5]. Quite generally, unless it collapses promptly to a black hole [10], the binary-merger product will oscillate in modes that leave the magnetic fields that develop after the merger [4, 5]. Quite generally, unless it collapses promptly to a black hole [10], the binary-merger product will oscillate in modes that leave the magnetic fields that develop after the merger [4, 5].

The recent discovery of a binary neutron star merger both across nearly the entire electromagnetic spectrum [1] and in gravitational waves [2]—not even two years after their first detection by LIGO in black-hole mergers [3]—as well as the striking confirmation of such mergers as the central engine of short gamma-ray bursts (see [4, 5] for reviews) heralds the era of gravitational wave astronomy. Detailed observations of such events could provide valuable information about the properties of matter at extreme density and temperature. With a few exceptions [6, 7] current simulations of neutron-star mergers neglect the transport properties of the material, assuming that they are too small to operate on dynamical timescales [8]. We revisit this assumption by exploring the impact of viscosity and thermal transport after merger, exploiting results of simulations. These have seen enormous progress [9–13] and found that for not too massive or too asymmetric systems, the post-merger object is metastable to gravitational collapse over tens of milliseconds. The inner region of this object, ~10 km across, can reach several times nuclear-matter saturation (number) density $n_0 \approx 0.16\text{fm}^{-3}$ and temperatures of tens of MeV.

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**Thermal equilibration.** To establish whether heat diffusion is significant, consider a region of size $z_{\text{typ}}$ that is hotter than its surroundings by a temperature difference $\Delta T$. For a material with specific heat per unit volume $c_V$ and thermal conductivity $\kappa$, this region has an additional thermal energy $E_{\text{th}} \approx (\pi/6)c_V z_{\text{typ}}^3 \Delta T$ and (assuming a smooth temperature distribution so that the thermal gradient is $\Delta T/z_{\text{typ}}$) heat is conducted out of the region at a rate $W_{\text{th}} \approx \pi \kappa \Delta T z_{\text{typ}}^2$. The thermal equilibration time, needed to conduct away a significant fraction of the extra thermal energy, is $\tau_\kappa \equiv E_{\text{th}}/W_{\text{th}} = c_V z_{\text{typ}}^2/(6\kappa)$. The specific heat is dominated by neutrons, which have the largest phase space of low-energy excitations, giving $c_V \approx 1.0 m_n^2 \pi^3 / T$, assuming a Fermi liquid of neutron density $n_n$ with Landau effective mass $m_n^* [19]$. Particles of number density $n_i$, typical speed $v_i$, and mean free path (mfp) $\lambda_i$, contribute to the thermal conductivity as $\kappa \propto \sum_i \kappa_i \propto \sum_i n_i v_i \lambda_i$, so $\kappa$ is effectively dominated by particles with the optimal combination of high density and long mfp. Neutrons, though numerous, are strongly interacting and have a very short mfp, thus thermal conductivity is dominated by electrons or neutrinos.

Below a few MeV, the neutrino mfp becomes longer than the merger region [4, 20], so neutrinos escape and thermal conductivity is dominated by electrons which scatter mainly via exchange of Landau-damped transverse photons. In this approximation, the thermal conductivity is temperature-independent $\kappa_e \approx 1.5 n_e^{2/3} / \alpha [\text{Eq. (40) of [21]}]$, where $n_e$ is the electron number density and $\alpha \approx 1/137$. This yields a lower bound for the thermal equilibration time in the electron
dominated regime

\[ \tau_{\kappa}^{(r)} = 5 \times 10^8 \text{s} \]

\[ \times \left( \frac{0.1}{x_p} \right)^{\frac{2}{3}} \left( \frac{m_n^*}{0.8 m_n} \right)^{\frac{1}{3}} \left( \frac{n_0}{n_B} \right)^{\frac{1}{3}} \left( \frac{z_{\text{typ}}}{1 \text{ km}} \right)^{2} \left( \frac{T}{1 \text{ MeV}} \right), \]

where \( n_B \) is the baryon number density, \( n_0 \) nuclear saturation density and \( x_p \equiv n_e/n_B \) is the proton fraction. Clearly, this timescale is far too large to have an impact on the \( \sim 10 \text{ ms} \) timescale of post-merger processes [4].

At temperatures \( T \gtrsim 10 \text{ MeV} \), neutrinos become trapped for nucleon density \( n \gtrsim n_0 \), since the neutrino mfp, which at high density depends strongly on in-medium corrections [20, 22], becomes smaller than the star. Electron neutrinos form a degenerate Fermi gas with a Fermi momentum \( p_{F,\nu} \) of about half that of the electrons. Their mfp is longer than that of the electrons, so they dominate the thermal conductivity [23], which is given by \( \kappa_{\nu} \approx 0.33 n_{\nu}^{2/3}/(G_F^2 (m_n^*)^2 n_e^{1/3} T) \), where \( G_F \equiv 1/(293 \text{ MeV})^2 \) is the Fermi coupling. This yields the timescale for thermal transport via neutrinos

\[ \tau_{\kappa}^{(\nu)} \approx 0.7 \text{ s} \]

\[ \times \left( \frac{0.1}{x_p} \right)^{\frac{2}{3}} \left( \frac{m_n^*}{0.8 m_n} \right)^{\frac{3}{2}} \left( \frac{\mu_e}{2 \mu_\nu} \right)^{\frac{2}{3}} \left( \frac{z_{\text{typ}}}{1 \text{ km}} \right)^{2} \left( \frac{T}{10 \text{ MeV}} \right)^{\frac{2}{3}}. \]

In summary, for neutrino-driven thermal transport to be important on a timescale of tens of milliseconds, there would e.g. have to be thermal gradients (e.g., from turbulence) on length scales of order 0.1 km. Moreover, heat transport into cooler regions should manifest itself even more quickly.

**Shear dissipation.** To estimate the shear-viscosity timescale, consider a fluid of rest-mass density \( \rho \) flowing in the \( x \) direction at velocity \( v_x \), having kinetic energy per unit volume \( E_{\text{kin}} = \frac{1}{2} \rho v_x^2 \). If the fluid has shear viscosity \( \eta \), then the energy dissipated per unit time and unit volume is \( W_{\text{shear}} \approx \eta (d v_x/dz)^2 \), and the time needed for shear viscosity to dissipate a significant fraction of the kinetic energy is \( \tau_\eta \equiv E_{\text{kin}}/W_{\text{shear}} \). We assume that the flow is fairly uniform, with the velocity varying by a factor of order unity over a distance \( z_{\text{typ}} \), in the \( z \) direction, so \( d v_x/dz \approx v_x/z_{\text{typ}} \) which gives \( \tau_\eta \approx \rho z_{\text{typ}}^2/(2\eta) \).

In the low-temperature, electron-dominated regime \( (T \lesssim 10 \text{ MeV}) \), using the dominant transverse contribution from [24] [Eq. (2.4) in [25]] with the damping scale \( q_l^2 \equiv 4\alpha p_{F,e}/\pi \), we find \( \eta^{(e)} \approx 0.2 n_e^{14/9}/(\alpha^{5/3} T^{5/3}), \) so

\[ \tau_{\eta}^{(r)} \approx 1.6 \times 10^8 \text{s} \left( \frac{z_{\text{typ}}}{1 \text{ km}} \right)^2 \left( \frac{1}{1 \text{ MeV}} \right)^5 \left( \frac{n_0}{n_B} \right)^{\frac{1}{9}} \left( \frac{0.1}{x_p} \right)^{\frac{14}{9}} \left( \frac{T}{10 \text{ MeV}} \right), \]

again being far too large to be relevant.

However, in the high-temperature, neutrino-dominated regime \( (T \gtrsim 10 \text{ MeV}) \) neutrinos produce a much larger shear viscosity \( \eta^{(\nu)} \approx 0.46 n_e^{4/3}/(G_F^2 (m_n^*)^2 n_e^{1/3} T^2) \) [23], which yields

\[ \tau_{\eta}^{(\nu)} \approx 54 \text{s} \left( \frac{0.1}{x_p} \right)^{\frac{2}{3}} \left( \frac{m_n^*}{0.8 m_n} \right)^{\frac{2}{3}} \left( \frac{\mu_e}{2 \mu_\nu} \right)^{\frac{4}{3}} \left( \frac{z_{\text{typ}}}{1 \text{ km}} \right)^{2} \left( \frac{T}{10 \text{ MeV}} \right)^{\frac{2}{3}}. \]

Interestingly, like eq. (2), this result depends only weakly on the density, via the proton fraction \( x_p \), the effective mass \( m_n^* \), and the ratio \( \mu_e/\mu_\nu \). In summary, neutrino shear viscosity could play an important role, i.e., \( \tau_{\eta}^{(\nu)} \) could be in the millisecond range, if the neutrino density is anomalously high or if there are flows that experience shear over short distances, \( z_{\text{typ}} \approx 0.01 \text{ km} \), for example, due to turbulence or high-order non-axisymmetric instabilities [26–29].

**Bulk viscosity.** To study the impact of bulk viscosity, we consider an “averaged” bulk viscosity \( \zeta \) in response to a periodic compression-rarefaction cycle. In nuclear matter, dissipation arises because the rate of beta equilibration of the proton fraction via Urca processes occurs on the same timescale, so that the proton fraction lags behind the applied pressure. If the oscillations after the merger are roughly periodic, we expect that the dissipation induced by pressure variations occurring on a timescale \( t_{\text{dens}} \) can be estimated by using the bulk viscosity evaluated at frequency \( f = 1/t_{\text{dens}} \). The bulk viscosity is largest when the internal equilibration rate matches the frequency of the oscillation. Furthermore, because the equilibration rate is sensitive to the temperature, the bulk viscosity shows a resonant maximum as a function of temperature (e.g., Fig. 7 in [30]). For oscillations with a timescale \( t_{\text{dens}} \), the resonant maximum value is [30]

\[ \zeta_{\text{max}} \equiv Y_\zeta \bar{n} t_{\text{dens}}, \quad Y_\zeta \equiv C^2/(4\pi B \bar{n}), \]

where \( B \equiv -(1/\bar{n}) \left( \partial \delta \mu / \partial x_p \right)_\bar{n} \) and \( C \equiv \bar{n} \left( \partial \delta \mu / \partial n \right)_{x_p} \) are the nuclear susceptibilities with respect to baryon density and proton fraction, where the chemical potential \( \delta \mu \equiv \mu_n - \mu_p - \mu_e \) characterises, in the absence of neutrino trapping, the degree to which the system is out of beta equilibrium. This maximum value \( \zeta_{\text{max}} \) depends only on properties of the EOS and is independent of the flavor re-equilibration rate. Changing the re-equilibration rate moves the curve in Fig. 7 in [30] “horizontally”, changing the temperature at which the maximum value is attained.

We note that the maximum bulk viscosity is a monotonically
cally increasing function of number density and Fig. 1 shows the prefactor $Y_\xi$ for nuclear matter obeying various EOSs, all of which can sustain a $2 M_\odot$ neutron star [31, 32]. Whereas APR [33] is a cold EOS and is included here for comparison, for all the others we use “hot” EOSs calculated using a model of nuclei and interacting nucleons in statistical equilibrium [34]. In addition to the LS220 [35], used for the simulations below, these EOSs range from the moderately soft SFHo [36] through the increasingly stiff DD2 [36, 37] and TMA [34], to the extremely stiff NL3.

Now consider the temperature $T_{\text{max}}$ at which bulk viscosity reaches its resonant maximum. For small-amplitude oscillations $T_{\text{max}} = (2\pi f/(\bar{\Gamma}))^{1/\delta}$ [30], where $\Gamma$ is the prefactor in the equilibration rate, $\Gamma = \Gamma^0 n^\delta \rho$. For modified-Urca processes, $\delta = 6$, so $1/\delta$ is small, making $T_{\text{max}}$ insensitive to details of the EOS. As a result, over the entire relevant frequency range, i.e., from a few tenths to several kHz, we find for flavor equilibration via nuclear modified-Urca "nmU" processes

$$T_{\text{max}}^{\text{nmU}} \approx 4 - 7 \text{ MeV} \approx 5 - 8 \times 10^{10} \text{ K},$$

which is well within the range of temperatures expected for dense matter in the post-merger [4, 38, 39].

It should be noted that flavor re-equilibration might instead occur via direct-Urca reactions, which are orders of magnitude faster than modified-Urca processes, giving much lower bulk viscosities at $T \sim 5 \text{ MeV}$, since the resonant maximum of bulk viscosity would have moved to lower temperatures (Fig. 7 in [30]). In neutrino-transparent matter at $T = 0$, direct-Urca processes are allowed when $\Delta_F = p_{F,n} - p_{F,p} - p_{F,e} < 0$. In Fig. 2 we plot this kinematic constraint as a function of density for the same EOSs in Fig. 1. For softer EOSs (e.g., SFHo, DD2) direct-Urca processes are never possible at $T = 0$; however, for APR the direct-Urca channel opens at $n > 5 n_0$. For even stiffer EOSs (LS220, NL3, TMA) it already opens around twice saturation density, yet these EOSs have been challenged by nuclear physics constraints [40]. These considerations suggest that the amount of bulk-viscous damping will be a sensitive indicator of whether the EOS allows direct-Urca processes at the densities and temperatures prevalent in neutron star mergers. A more precise connection with the EOS will require calculations of the beta equilibration rate that incorporate the effects of temperature, strong interactions, and the gradual opening of phase space above the direct Urca threshold.

We now estimate the dissipation time for compression oscillations. The energy density for a baryon number-density oscillation of amplitude $\Delta n$ around average density $\bar{n}$ is $\varepsilon_\text{comp} \approx K \bar{n} (\Delta n/\bar{n})^2/18$ [41], where $K$ is the nuclear compressibility at that density. If the compression varies on a timescale $t_{\text{damp}}$, then, in a material with bulk viscosity $\zeta$, the dissipated power per unit volume is [42] $(d\varepsilon/dt)_\text{bulk} \approx 2\pi^2 \zeta (\Delta n/\bar{n})^2/\tau_{\text{damp}}^2$. Hence, the time required for bulk viscosity to have a significant impact on the oscillations of the system is

$$\tau_\zeta \equiv \varepsilon_{\text{comp}}/(d\varepsilon/dt)_\text{bulk} \approx K \bar{n} t_{\text{damp}}^2/(36\pi^2 \zeta).$$

Expecting bulk viscosity to reach its maximum value $\zeta_{\text{max}}$ [Eq. (5)] at typical neutron-star merger temperatures [Eq. (6)], we can use Eq. (5) in (7) to find that, when the direct-Urca channel is not open, the minimum timescale for bulk viscosity to impact the oscillations is

$$\tau_{\zeta,\text{min}} \approx 3 \text{ ms} \left( \frac{t_{\text{damp}}}{1 \text{ ms}} \right) \left( \frac{K}{250 \text{ MeV}} \right) \left( \frac{0.25 \text{ MeV}}{Y_\xi} \right).$$

Stated differently, under conditions of maximum bulk viscosity, the damping timescale is a few times larger than the typical timescale $t_{\text{damp}}$ of density variations.

Since strong emission of gravitational waves occurs from the high-density region of the star during the first $\sim 5$ milliseconds after the merger, when characteristic frequencies $f_1$
and $f_3$ appear in the gravitational-wave spectrum [16, 18, 43], bulk viscous damping is most likely to have observable con-
sequences if, during that early time, there are density os-
cillations occurring on a millisecond timescale in parts of
the high-density region where the bulk viscosity is maximal
($T \sim 4 - 7\, \text{MeV}$).

To test whether such conditions are met, we show in Figs. 3
and 4 results from a state-of-the-art simulation of a symmetric
merger of $M = 2 \times 1.35M_\odot$, consistent with GW170817
[2], using the LS220 EOS [35], where $t = 0$ is the time of
merger [43]. Figure 3 uses a colorcode to show the ex-
pansion flow timescale $t_{\text{flow}} \equiv 1/⟨(∇ \cdot \vec{v})⟩ = \rho/D_tρ$ where
( ) represents a time average over a 2 ms time window and
where $D_t$ is the Lagrangian time derivative in Newtonian hy-
drodynamics [44]. This quantity is easily measured and,
for a harmonic density oscillation, it is related to Eqs. (7) and
(8) by $t_{\text{dens}} \approx (4\Delta n/\bar{n})t_{\text{flow}}$. Figure 3 reports $t_{\text{flow}}$ 2.4 ms
after the merger, where the post-merger object is in its vi-
olent and shock-dominated transient phase, (see [43] for a
toy model of this phase). Inside the green contour, the rest-
mass density is above nuclear saturation. The red and gray
lines are temperature contours at $4\, \text{MeV}$ and $7\, \text{MeV}$, re-
spectively. Overall, Fig. 3 shows that there are significant regions
where Eq. (8) is a valid estimate of the dissipation time be-
cause the density is high and the temperature is in the range
that maximizes bulk viscosity [Eq. (6)]. Since in these re-
gions $t_{\text{flow}} \sim 0.1 - 1\, \text{ms}$ and $\Delta n/\bar{n} \sim 1$, we conclude that
$t_{\text{dens}} \approx (4\Delta n/\bar{n})t_{\text{flow}} \sim t_{\text{flow}}$, is indeed in the millisecond
range.

This conclusion is reinforced by Fig. 4, which shows the evolu-
tion of various local properties of representative tracer par-
ticles in the inner region of the merger product [45]. From
the top panel, which reports the evolution of the temperature,
we see that all tracers pass through the temperature range of
large bulk viscosity (dark and light-gray shaded areas, show-
ing the regions of maximum and up to an order of magnitude
smaller dissipation) during the first few milliseconds. The sec-
ond panel reports the evolution of the normalized rest-mass
density and shows that at early times ($t \lesssim 5\, \text{ms}$) there are vari-
ations of order 100% in the rest-mass density on a timescale of
milliseconds, confirming that $t_{\text{dens}}$ is in that range. The third
panel shows the average of $t_{\text{flow}}$ for the tracers, which is in the
$0.1 - 1\, \text{ms}$ range, as expected from Fig. 3. Finally, the bottom
panel of Fig. 4 is a spectrogram averaging the power spectral
densities of the normalized rest-mass densities in the second
panel and showing how, throughout the first $20\, \text{ms}$, the merger
product has oscillation with a significant power at frequencies
in the kHz range.

The results shown in Figs. 3 and 4, combined with
Eq. (8), suggest that if direct Urca processes remain sup-
pressed, then significant bulk viscous dissipation may occur
on timescales of a few milliseconds, which is fast enough to
affect the flow of nuclear material, and hence the emitted grav-
itational signal. Full numerical-relativity simulations account-
ing for bulk viscosity are necessary to quantify the amount of
such dissipation and its impact on the gravitational-wave sig-
nal.

Conclusions. Material properties can only play a significant
role in neutron-star mergers if the relevant dissipation time is
comparable with or shorter than the survival time of the post-
merger object. Using typical values found in numerical simu-
lations, we find that shear viscosity and thermal conduc-
tivity are not likely to play a major role in post-merger dynamics
unless neutrino trapping occurs, which requires $T \gtrsim 10\, \text{MeV}$,
and $\tau_{\text{typ}} \lesssim 0.01\, \text{km}$. On the other hand, if direct-Urca pro-
cesses remain suppressed, leaving modified-Urca processes
to establish flavor equilibrium, then bulk viscous dissipation
could provide significant damping of the high-amplitude den-
sity oscillations observed right after merger. We conclude that
viscous dissipative processes deserve more careful investiga-
tion since they may well affect the spectral properties of the
post-merger gravitational-wave signal, especially the $f_1$ and
$f_3$ peaks that are produced right after the merger and that are
dissipated rapidly [14–16, 18, 46]. Since these peaks are rou-
tinely employed to infer the properties of the EOS [47, 48],
a more realistic treatment is particularly important. In addi-
tion, if viscous dissipation is active after the merger, it will
also heat the merger product, possibly stabilizing it on longer
timescales via the extra thermal pressure [10, 49–51]. If fu-
ture gravitational-wave observations indicate that the actual
dissipation is much smaller than what is suggested by Eq. (8),
e.g. if merger material transforms to quark matter, this would put limits on the fraction of matter for which direct-Urca processes are suppressed.

There are various directions in which our research can be further developed. First, the effects of bulk viscosity should be consistently included in future merger simulations. This has not been attempted before and requires a formulation of the relativistic-hydrodynamic equations that is hyperbolic and stable (see Chap. 6 of [44] for the associated challenges). Second, the bulk viscous effects discussed so far may be amplified by nonlinear suprathermal enhancement [30, 52–55] (which to a weaker extent also affects neutrino cooling, see e.g. [52, 56]), or by the even stronger phase-conversion dissipation [57]. Third, because the role played by shear viscosity depends on the typical scale-height of the fluid flow, investigations of the development of turbulent motion in the post-merger phase will be essential. Finally, given the role they play in determining the strength of thermal transport and of shear/bulk dissipation, neutrino trapping and direct-Urca processes motivate additional work to constrain the conditions under which these phenomena occur. We plan to consider some of these topics in our future work.

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