Detectability of thermal neutrinos from binary neutron-star mergers and implications for neutrino physics

Koutarou Kyutoku[1,2,3,4] and Kazumi Kashiyama[5]

[1]*Theory Center, Institute of Particle and Nuclear Studies, KEK, Tsukuba 305-0801, Japan*
[2]*Department of Particle and Nuclear Physics, the Graduate University for Advanced Studies (Sokendai), Tsukuba 305-0801, Japan*
[3]*Interdisciplinary Theoretical Science (iTHES) Research Group, RIKEN, Wako, Saitama 351-0198, Japan*
[4]*Center for Gravitational Physics, Yukawa Institute for Theoretical Physics, Kyoto University, Kyoto 606-8502, Japan 606-8502, Japan*
[5]*Department of Physics, the University of Tokyo, Bunkyo, Tokyo 113-0033, Japan*

We propose a long-term strategy for detecting thermal neutrinos from the remnant of binary neutron-star mergers with a future M-ton water-Cherenkov detector such as Hyper-Kamiokande. Monitoring $\gtrsim 2500$ mergers within $\lesssim 200$ Mpc, we may be able to detect a single neutrino with a human time-scale operation of ≈ 80 Mtyears for the merger rate of 1 Mpc^{-3} Myr^{-1}, which is slightly lower than the median value derived by the LIGO-Virgo Collaboration with GW170817. Although the number of neutrino events is minimal, contamination from other sources of neutrinos can be reduced efficiently to ≈ 0.03 by analyzing only ≈ 1 s after each merger identified with gravitational-wave detectors if gadolinium is dissolved in the water. The contamination may be reduced further to ≈ 0.01 if we allow the increase of waiting time by a factor of ≈ 1.7. The detection of even a single neutrino can pin down the energy scale of thermal neutrino emission from binary neutron-star mergers and could strongly support or disfavor formation of remnant massive neutron stars. Because the dispersion relation of gravitational waves is now securely constrained to that of massless particles with a corresponding limit on the graviton mass of $\lesssim 10^{-22}$ eV/c^2 by binary black-hole mergers, the time delay of a neutrino from gravitational waves can be used to put an upper limit of $\lesssim O(10)$ meV/c^2 on the absolute neutrino mass in the *lightest* eigenstate. Large neutrino detectors will enhance the detectability, and, in particular, 5 Mt Deep-TITAND and 10 Mt MICA planned in the future will allow us to detect thermal neutrinos every ≈ 16 and 8 years, respectively, increasing the significance.

I. INTRODUCTION

The discovery of GW170817 marked an opening of multimessenger astronomy with binary neutron-star mergers [1,2]. Gravitational waves driving the mergers are important targets for ground-based detectors such as Advanced LIGO and Advanced Virgo and enable us to study the equation of state of supranuclear density matter [1]. Differently from binary black holes, binary neutron stars can also become bright in electromagnetic channels [3,4]. The merger remnants are the prime candidate for the central engine of short-hard gamma-ray bursts [5–8] as has already been suggested by association of GRB 170817A with GW170817 [9] (but see Refs. [10,11]). Binary neutron stars should also eject neutron-rich material during the merger and post-merger phases, and this material will be synthesized to heavy neutron-rich nuclei, namely the *r*-process elements [12,13]. This scenario is supported by detections of electromagnetic counterparts consistent with the macronova/kilonova, an optical-infrared transient powered by decay of the *r*-process elements [14], after GW170817 (see, e.g., Ref. [2]).

Signals from binary neutron-star mergers are not limited to gravitational and electromagnetic radiation. Because the violent merger heats up the high-density material, thermal neutrinos with $\gtrsim 10$ MeV should also be emitted from the remnant of binary neutron-star mergers [15–17].

While direct detections of thermal neutrinos could be an important step toward understanding the realistic merger process as well as their impact on gamma-ray bursts [18,19] and *r*-process nucleosynthesis [20,21], they are quite challenging. On one hand, as we will see later, the detection is hopeless for a single merger at a distance $\gtrsim 100$ Mpc, where the mergers are expected to occur more than once a year [22,23]. On the other hand, "the diffuse neutron-star-merger neutrino background," i.e., superposition of neutrinos from all the binary neutron-star mergers throughout the Universe, is inevitably hidden by the diffuse supernova neutrino background also known as supernova relic neutrinos (see Ref. [24] for reviews), because the rate of supernovae must be higher by a few orders of magnitude than that of binary neutron-star mergers. However, we would like to stress again that detecting thermal neutrinos

will be important to understand binary neutron-star mergers accurately, as theoretical models of supernova explosions are qualitatively confirmed with detections of neutrinos from SN 1987A [25,26].

The chance of detections lies in the gap between these two standard ideas (see also Refs. [27–29] for a third idea on detecting supernova neutrinos). In this paper, we propose a long-term strategy to detect thermal neutrinos from the remnant of binary neutron-star mergers by monitoring many mergers identified by gravitational-wave detectors. Stacking multiple mergers is necessary except for serendipitous nearby mergers, but careless analysis will easily bury MeV neutrinos from binary neutron-star mergers in those from other sources such as the diffuse supernova neutrino background and atmospheric neutrinos including invisible muons (see Ref. [30] for stacking of high-energy neutrinos). A remarkable point is that gravitational-wave observations may determine the time of merger to an accuracy of ≈ 1 ms. Indeed, this precision has already been realized for observed binary black-hole mergers [31].[1] By analyzing the data of neutrino detectors only during ≈ 1 s from each merger, which is much shorter than 1000 s adopted in current counterpart searches with MeV-neutrino detectors [32,33], we can substantially reduce contamination from other sources of neutrinos. A similar idea has been proposed for detecting gravitational waves informed by gamma-ray bursts [34].

If the arrival time of such neutrinos with respect to the time of merger identified by gravitational waves is successfully determined with modest delay of 0.1–1 s, we may be able to put an upper limit on the absolute neutrino mass in the *lightest* eigenstate [35,36] (see also Ref. [37]). Because neutrinos have finite masses of $\lesssim 0.1$ eV/c^2, their travel speed is necessarily slower than the speed of light c, where the precise value depends on the mass eigenstate and the energy. Importantly, detections of gravitational waves from binary black-hole mergers have successfully shown that the dispersion relation of gravitational waves is accurately described by that of massless particles with a corresponding limit on the gravitons mass of $\lesssim 10^{-22}$ eV/c^2 [31,38], practically negligible. Thus, the difference of arrival times between these two messengers, or relative time of flight, will allow us to infer the mass of neutrinos. Although chances are not necessarily large, the range of accessible masses seems worth of serious consideration.

This paper is organized as follows. First, we summarize current understanding of neutrino emission from binary

---

[1]Because the gravitational-wave frequency at mergers of binary neutron stars is likely to be higher than the values accessible by current detectors, we might have to rely on theoretical models to determine the time of merger in realistic situations. The time of merger can vary by up to a few ms depending on the equation of state of neutron-star matter. We expect that the situation will improve by constraining the equation of state using gravitational-wave observations themselves.

neutron-star mergers in Sec. II. Next, we describe our strategy to detect thermal neutrinos from binary neutron-star mergers in Sec. III, and the expected level of contamination is examined in Sec. IV. Implication of detecting thermal neutrinos is discussed in Sec. V. Section VI is devoted to a summary. While we focus only on binary neutron stars in this work due to their expected dominance, it is straightforward to enhance our discussion to include black hole–neutron star binaries.

II. THERMAL NEUTRINO FROM BINARY NEUTRON-STAR MERGERS

We first summarize characteristics of thermal neutrinos emitted from the remnant of binary neutron-star mergers. In this work, we focus on the case that the lifetimes of remnant neutron stars are longer than ~ 1 s and thus substantial neutrino emission can be expected. This is not for optimistic simplification, but the scenario that we would like to verify or reject by detecting thermal neutrinos. The lifetime of remnant massive neutron stars depends on various details and can in fact be very short [17,39]. Fortunately, the prompt collapse to a black hole is not very likely in light of the maximum mass of spherical neutron stars exceeding ≈ 2 $M_\odot$ [40].

Various numerical simulations have shown that the remnant massive neutron stars are heated up to several tens of MeV at the collision unless the merger results in a prompt collapse [15–17]. Then, thermally-produced electron-positron pairs are captured on nucleons to emit a copious amount of electron neutrinos ν_e and antineutrinos $\bar{\nu}_e$ with the rise time of $\lesssim 10$ ms from the merger. In the case of binary neutron-star mergers, $\bar{\nu}_e$ is brighter than ν_e due to the neutron richness. The peak luminosity of electron antineutrinos reaches $1-3 \times 10^{53}$ erg s^{-1} with the typical energy 10–30 MeV depending on binary parameters and unknown equations of state for supranuclear-density matter [41–44]. Pair processes such as the electron-positron annihilation also emit muon and tau (anti)neutrinos. We note that neutrino oscillations in the source region such as the Mikheyev-Smirnov-Wolfenstein effect, bipolar oscillations, and matter-neutrino resonance begin to be explored only recently for remnant massive neutron stars [45–48]. They could reduce effective luminosity of electron antineutrinos to some extent.

So far, little is known about the realistic spectrum of neutrinos from binary neutron-star mergers despite their importance for quantifying the detectability. Monte-Carlo neutrino-transport simulations suggest that the spectrum can be approximated by pinched Fermi-Dirac distribution for binary neutron-star mergers as in the case of supernova explosions [49]. However, it is difficult to determine the degree of distortion at this stage.

The duration of neutrino emission is expected to be a few to ten seconds unless it is shut down by the collapse of the remnant during the hot phase due to angular-momentum

redistribution [17], whereas detailed long-term calculations are not available beyond ≈ 0.5 s [19]. The total energy given to neutrinos is determined by hydrodynamic interactions and will be similar to, or moderately less than, that in supernova explosions (see Ref. [50] for reviews).

III. DETECTING THERMAL NEUTRINOS

A. Expected number of neutrino events

We aim at detecting thermally-produced neutrinos from remnant massive neutron stars formed after binary neutron-star mergers with a water-Cherenkov detector such as planned ≈ 0.37 Mt Hyper-Kamiokande [51,52].[2] Water-Cherenkov detectors are efficient at detecting electron antineutrinos via the inverse β decay, $p + \bar{\nu}_e \rightarrow n + e^+$. The expected number of neutrino events for a single merger is estimated by

$$N_\nu = N_T \int_{t_i}^{t_f} \int_{E_{\min}}^{E_{\max}} \phi(E,t)\sigma(E)dEdt, \qquad (1)$$

where N_T is the number of target protons in the detector, E is the energy of electron antineutrinos, $\phi(E,t)$ is the number flux of electron antineutrinos per unit energy, $\sigma(E)$ is the capture cross section of an electron antineutrino on a proton. The number of target protons is given in terms of the (effective) mass of water M_T as $N_T \approx (M_T/m_{\rm p}) \times (2/18) = 6.7 \times 10^{34}(M_T/1 \text{ Mt})$ with $m_{\rm p}$ the proton mass. The cross section is calculated to various levels of approximations [53,54], and in this study we adopt Eq. (7) of Ref. [24],

$$\sigma(E) = 9.5 \times 10^{-42} \text{ cm}^2 \left(\frac{E - 1.3 \text{ MeV}}{10 \text{ MeV}}\right)^2 \left(1 - \frac{7E}{m_{\rm p}c^2}\right). \qquad (2)$$

Note that the positrons from inverse β decay are distributed nearly isotropically [53]. Threshold energies $\{E_{\min}, E_{\max}\}$ should be determined to span the range relevant to neutrinos from binary neutron-star mergers while suppressing contamination from other sources of neutrinos. Threshold times $\{t_i, t_f\}$ should be determined to detect intense neutrino emission around the peak time, while the level of contamination should be kept as low as possible.

In this work, we adopt the Fermi-Dirac distribution with temperature T and zero chemical potential, for which the average energy of neutrinos is given by $\langle E \rangle \approx 3.15 k_{\rm B} T$ with $k_{\rm B}$ the Boltzmann constant. Thus, by ignoring the time dependence, the spectrum or number flux per unit energy takes the form

$$\phi(E) = \frac{c}{2\pi^2(\hbar c)^3} \frac{E^2}{\exp[E/(k_{\rm B}T)] + 1}, \qquad (3)$$

where $\hbar$ is the reduced Planck constant. The expected rate of neutrino events is obtained by integrating Eq. (1) over a given energy interval, and the result is usefully characterized by the typical energy, $\langle E \rangle$, and the leading-order cross section (called "naive" in Ref. [54])

$$\sigma_{\rm LO}(E) = 9.5 \times 10^{-42} \text{ cm}^2 \left(\frac{E}{10 \text{ MeV}}\right)^2, \qquad (4)$$

as

$$\frac{dN_\nu}{dt} \approx f_E N_T \frac{L_\nu}{4\pi D^2} \frac{\sigma_{\rm LO}(\langle E \rangle)}{\langle E \rangle}, \qquad (5)$$

where L_ν is the luminosity of electron antineutrinos and D is the distance to the source. A factor $f_E \sim 1$ is a number determined by $\{E_{\min}, E_{\max}\}$ (for the assumed spectrum), and we show contribution from different energy ranges in Table I for various values of $\langle E \rangle$. Note that $\sigma_{\rm LO}(E)$ is never used to compute the expected number of neutrino events, and we use it only for normalizing the result keeping the leading-order dependence on $\langle E \rangle$ transparent. This table shows that anywhere between 10 MeV and 50 MeV can contribute appreciably to detections. Possible choices of $\{E_{\min}, E_{\max}\}$ will be discussed later in Sec. IV along with contamination from other sources of neutrinos.

The expected number of neutrino events is obtained by integrating Eq. (5) in time, but time evolution of the luminosity is not understood in detail, particularly on a time scale of $\gtrsim 1$ s. Here, we focus on $\Delta t_{\rm obs} \approx 1$ s from the merger and denote the total energy of electron antineutrinos emitted during $\Delta t_{\rm obs}$ as $E_{\Delta t} = \int L_\nu dt$, which may be $\approx 3 \times 10^{52}$ erg for moderately compact remnant neutron stars [19].[3] This restriction in time is due partly to the lack of knowledge about long-term evolution of the neutrino luminosity, but this is mainly intended to reduce contamination from other sources of neutrinos in realistic observations as described later in Sec. IV. The total energy of neutrinos may be increased by a factor of 2–3 if we take $\Delta t_{\rm obs} \approx 10$ s as the available energy budget suggests [17].

By regarding the typical energy of neutrinos, $\langle E \rangle$, as an appropriate time average, the expected number of neutrino events for a single merger is found to be

$$N_\nu \approx 1.0 \times 10^{-3} \times f_E f_{\rm se} f_{\rm osc} \left(\frac{M_T}{1 \text{ Mt}}\right) \left(\frac{E_{\Delta t}}{3 \times 10^{52} \text{ erg}}\right)$$
$$\times \left(\frac{\langle E \rangle}{10 \text{ MeV}}\right) \left(\frac{D}{100 \text{ Mpc}}\right)^{-2}. \qquad (6)$$

---

[2]See also http://lib-extopc.kek.jp/preprints/PDF/2016/1627/1627021.pdf for a recent design of Hyper-Kamiokande.

[3]Note that Ref. [19] do not incorporate viscous heating, and thus the neutrino luminosity is likely to be underestimated.

TABLE I.   Contribution to f_E from different energy ranges $[E_{\min} : E_{\max}]$ for a variety of the typical energy of neutrinos, $\langle E \rangle$. Recall f_E denotes the ratio of the number of neutrino interactions obtained by integrating the product of the Fermi-Dirac distribution with typical energy $\langle E \rangle \approx 3.15 k_B T$ and the cross section, Eq. (2), to the number for monoenergetic neutrinos with $\langle E \rangle$ and the leading-order cross section, Eq. (4). The final column, 10–50 MeV, is the range expected to be utilizable with Hyper-Kamiokande with Gd dissolution [51,55].

$\langle E \rangle$	0–10 MeV	10–20 MeV	20–30 MeV	30–40 MeV	40–50 MeV	10–50 MeV
10 MeV	0.16	0.53	0.20	0.03	<0.01	0.77
15 MeV	0.04	0.33	0.33	0.15	0.05	0.87
20 MeV	0.02	0.17	0.29	0.22	0.12	0.80

A fudge factor $f_{\rm se} \leq 1$ is the event selection efficiency for the inverse β decay, and we expect $f_{\rm se} \approx 0.9$ and 0.67 for Hyper-Kamiokande without and with gadolinium (Gd), respectively [51]. Another fudge factor $f_{\rm osc} \lesssim 1$ represents the effect of neutrino oscillations. As discussed in Sec. II, the oscillation in the source region are not fully understood yet. Therefore, we leave incorporation of these effects for future study and discuss only the vacuum oscillation during the propagation. In a tri-bimaximal mixing approximation [56], the survival probability of electron (anti)neutrinos is given by 5/9 and the appearance probability from muon and tau (anti)neutrinos is 2/9 each. Taking the lower luminosity and higher typical energy of muon and tau (anti)neutrinos from the remnant of binary neutron-star mergers [17,19], we expect that $f_{\rm osc}$ is not very far from unity. We discuss time dilation associated with the finite masses later in Sec. V B.

This estimate, Eq. (6), clearly shows that it is hopeless to detect thermal neutrinos from a single binary neutron-star merger except for extremely lucky stars at $\lesssim 3$ Mpc. This fact has already been found in previous work [17,42], and we just confirm it with slightly detailed calculations (see also Ref. [57,58] for relevant work on accretion flows).

B. Monitoring multiple mergers

Even though the expected number of neutrinos from a single merger is very low, superposition of many mergers gives us a fair chance of detections. As a qualitative order-of-magnitude estimate, we expect to receive a single thermal neutrino with probability $1 - (1 - N_\nu)^{1/N_\nu} \approx 1 - e^{-1} = 63\%$ with $1/N_\nu \approx 1000$ mergers at 100 Mpc. The problem is that we will have to wait longer than a decade to collect such a large number of mergers, so that neutrinos from other sources completely overwhelm thermal neutrinos from binary neutron-star mergers. However and most importantly, if we focus only on $\Delta t_{\rm obs} \approx 1$ s after each merger using timing information from gravitational-wave detectors, we can efficiently reduce the length of data from neutrino detectors to ≈ 1000 s in total. As we discuss later in Sec. IV, the expected number of contamination events can be reduced to much less than unity for this short duration. We note that neutrino detectors do not require low-latency alerts from gravitational-wave detectors to perform this analysis [32,33].

To assess the effectiveness of this strategy in a quantitative manner, we need to take the spatial distribution of binary neutron-star mergers into account. In this study, we assume the effective range $D_{\rm eff}$ of a gravitational-wave detector for binary neutron stars to be 200 Mpc, which is approximately the design sensitivity of Advanced LIGO [22]. The detectable volume is given by $4\pi D_{\rm eff}^3/3$ by definition, and the detection rate of mergers is given by multiplying this volume by the merger rate per unit volume per unit time of binary neutron stars, $\mathcal{R}$, which is derived to be $1.54^{+3.2}_{-1.22}$ Mpc^{-3} Myr^{-1} by the LIGO-Virgo Collaboration with GW170817 [1].

The period P or exposure PM_T that we have to wait with monitoring multiple mergers while a thermal neutrino is detected is estimated by the condition

$$1 = P \mathcal{R} f_\Omega \int_0^{D_{\rm eff}} N_\nu(D) \times 4\pi D^2 dD, \qquad (7)$$

where $N_\nu(D)$ is Eq. (6) regarded as a function of D. Here, we neglect the cosmological redshift and evolution history that should be insignificant at $\lesssim 200$ Mpc. A factor f_Ω reflects the antenna pattern of gravitational-wave detectors, i.e., dependence of the sensitivity on the sky position and orientation of sources. We can show that $f_\Omega \approx 0.8$ for a single detector (see Appendix) and expect this to approach unity for a network of detectors. Here, we neglect angular dependence of neutrino emission. It should be cautioned that this waiting time, P, denotes not only the operation time of the neutrino detector but also requires coincident operations of gravitational-wave detectors.

Solving Eq. (7), we finally obtain

$$PM_T = 80 \text{ Mt years}$$
$$\times \left(\frac{f_{\rm all}}{0.5}\right)^{-1} \left(\frac{E_{\Delta t}}{3 \times 10^{52} \text{ erg}}\right)^{-1} \left(\frac{\langle E \rangle}{10 \text{ MeV}}\right)^{-1}$$
$$\times \left(\frac{D_{\rm eff}}{200 \text{ Mpc}}\right)^{-1} \left(\frac{\mathcal{R}}{1 \text{ Mpc}^{-3} \text{ Myr}^{-1}}\right)^{-1}, \qquad (8)$$

where $f_{\rm all} \equiv f_E f_{\rm se} f_{\rm osc} f_\Omega$, as a waiting exposure PM_T for detecting nonzero events of thermal neutrinos from binary neutron-star mergers with probability 63%. If the observation period is taken to be xP for a given value of M_T,

TABLE II. Factors appearing in the estimate of the waiting time, Eq. (8), and their origins. We also present their expected values.

Symbol	Origin	Expected value
f_E	Energy range	≈ 0.8 (see Table I)
f_{se}	Selection efficiency	0.9 (no Gd) or 0.67 (Gd)
f_{osc}	Neutrino oscillation	$0.5 \sim 1$
f_Ω	Antenna pattern	$0.8 \sim 1$

detection probability is modified to $1 - e^{-x}$. Here, $E_{\Delta t}$ and $\langle E \rangle$ should be regarded as values averaged over astrophysical populations of binary neutron stars. The meaning of factors f is summarized in Table II. We note that the number of mergers during 80 years is ≈ 2700 and that the expected number of nearby mergers at $\lesssim 3$ Mpc is less than 0.01. As we describe in detail in Sec. IV, focusing on $\Delta t_{obs} \approx 1$ s reduces the length of data from neutrino detectors by a factor of 2700 s/80 years $\approx 10^{-6}$ compared to a blind search, dramatically suppressing contamination from other sources.

The period given by Eq. (8) is not very short but in the human time scale for a ~ 1 Mt detector. As a reference, Hyper-Kamiokande is expected to achieve 0.37 Mt in the near future, and another Hyper-Kamiokande is planned to be built in Korea with 0.26 Mt [52]. One benchmark for the acceptable waiting time is provided by Galactic supernovae, which are expected to occur once in 30–100 years. Therefore, the detection of thermal neutrinos from binary neutron stars may be as likely as that from Galactic supernovae, whereas the number of neutrino events are drastically different. Another (but related) difference is that binary neutron-star mergers will be observed steadily by gravitational waves, while a Galactic supernova is intrinsically rare (see also Ref. [27]). Furthermore, the prospect for constraining the neutrino mass can be higher for binary neutron-star mergers than for supernovae due to longer distances as we will describe in Sec. V B.

In this estimation, as well as in Eq. (6), we made several assumptions on astrophysical inputs such as the total energy, typical energy, spectrum, and merger rate. The most important assumption may be that remnant massive neutron stars do not collapse before sizable emission of neutrinos, and this is what we would like to verify or reject by detecting thermal neutrinos. The total energy is uncertain by a factor of order unity even within a long-lived remnant scenario and also depends on the duration of each observation, which we assume to be $\Delta t_{obs} \approx 1$ s. The typical energy is also uncertain by a factor of ≈ 2. Spectral deformation is likely to be a minor correction that can be absorbed in the variation of f_E (see Table I). While the merger rate is highly uncertain even after the discovery of GW170817, the fiducial value adopted here is on the conservative side (note also that this value was denoted as "realistic" in Ref. [22]). In any case, the merger rate will
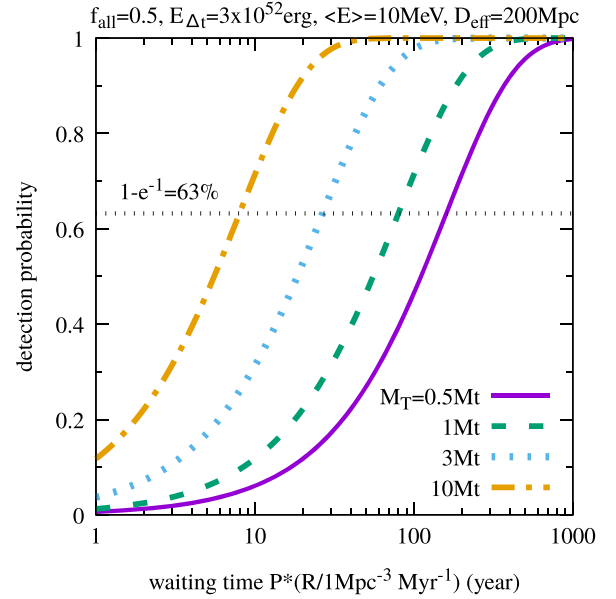

FIG. 1. Detection probability of nonzero events of thermal neutrinos as a function of time for various detector volumes. The waiting time is normalized to our fiducial merger rate of binary neutron stars. Other parameters are taken to be our fiducial values shown in Eq. (8). The horizontal dotted line indicates $1 - e^{-1} = 63\%$.

be understood in the near future by ongoing gravitational-wave observations. If some of these parameters conspire, the waiting time, P, could be shortened by a factor of $\gtrsim 5$.

Ultimately, a large effective volume of water- or ice-Cherenkov detectors is highly desired to increase the likelihood for detecting thermal neutrinos. Figure 1 shows the detection probability of nonzero events of thermal neutrinos as a function of time for various detector volumes with our fiducial parameters. The effective volume can be increased not only by a large detector but also by additional detectors such as Hyper-Kamiokande in Korea. It is conceivable that a more-than-M-ton detector like Deep-TITAND [28] and MICA [29] can be constructed within 30–100 years considered in our strategy. The waiting time can be reduced to less than 10 years with $M_T \gtrsim 8$Mt for our fiducial parameters, and then detections of thermal neutrinos from multiple binary neutron-star mergers (although one for each) could become possible within a realistic operating time of telescopes. For example, Baksan Underground Scintillation Telescope has been running longer than 30 years [59]. Even for such a large detector, and, in fact, irrespective of the detector volume, $\Delta t_{obs} \approx 1$ s has to be chosen to detect thermal neutrinos from binary neutron-star mergers with high significance as we discuss below.

IV. CONTAMINATION

Figure 2 shows the expected spectrum of thermal neutrinos from binary neutron-star mergers at water-Cherenkov detectors with the same level of contamination as planned
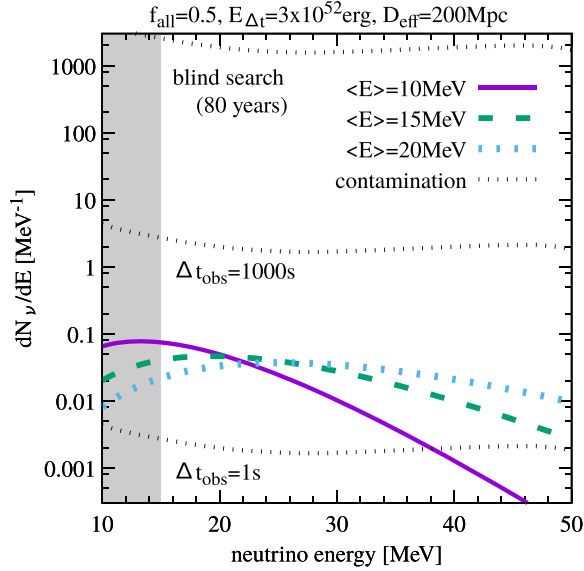
FIG. 2. Spectrum of detected thermal neutrinos normalized to a single event between 10 and 50 MeV, i.e., $\int_{10\text{ MeV}}^{50\text{ MeV}}(dN_\nu/dE)\,dE = 1$. Purple-solid, green-dashed, and blue-dotted curves show the expected spectra at the detector for $\langle E \rangle = 10$, 15, and 20 MeV, respectively, where we adopt Eq. (2) as the cross section. Other parameters are taken to be our fiducial values adopted in Eq. (8). Black dotted curves show the expected spectrum of contamination during a blind search of 80 years expected to be required for a single detection (top), currently adopted $\Delta t_{\text{obs}} = 1000$ s for all the mergers during this period (middle), and $\Delta t_{\text{obs}} = 1$ s we proposed in this work (bottom). Specifically, we consider invisible muons, atmospheric antineutrinos, neutral-current quasielastic scattering, and diffuse supernova background as the sources of contamination assuming Gd dissolution [51,55]. The shaded area on the left approximately represents the energy range unavailable due to the spallation background, where the precise location of the threshold will change.

Hyper-Kamiokande with Gd dissolution [51,55]. To clarify dependence of the background level on observation strategies, we plot the expected spectrum of contamination events per single thermal neutrino for (i) a blind search of 80 years, (ii) $\Delta t_{\text{obs}} = 1000$ s for each merger, and (iii) $\Delta t_{\text{obs}} = 1$ s for each merger proposed in this work. Specifically, we show the

sum of decay electrons from invisible muons, atmospheric antineutrinos, neutral-current quasielastic scattering, and diffuse supernova neutrino background [51,60]. All these spectra are independent of the detector volume as far as the number of contamination events scales linearly with the detector volume, except for the change in the waiting time $P \approx 80$ years for our fiducial parameters. If the merger rate is changed, the level of contamination for the blind search scales linearly with the waiting time, P, while the spectra for $\Delta t_{\text{obs}} = 1000$ s and 1 s are unchanged because of the identical number of mergers.

This figure clearly shows that the chance of detecting thermal neutrinos arises only when we focus on a short time interval of $\Delta t_{\text{obs}} \approx 1$ s right after the merger. Otherwise, for example with the currently adopted $\Delta t_{\text{obs}} = 1000$ s [32,33], thermal neutrinos from binary neutron-star mergers are heavily obscured by contamination from other sources of neutrinos. In the following, we discuss the expected level of contamination more quantitatively.

A. Quantitative assessment

We examine how severe contamination from other sources of neutrinos is to detect thermal neutrinos from binary neutron-star mergers focusing on Hyper-Kamiokande based on Ref. [51]. Characteristics of target neutrinos are very similar to those from supernova explosions, and thus backgrounds are basically the same as those encountered in searches of the diffuse supernova neutrino background [24]. This fact implies that Gd dissolution will significantly increase the prospect for detecting thermal electron antineutrinos from binary neutron-star mergers via tagging neutrons from inverse β decay [55]. Important numbers are summarized in Table III and the final paragraph of this subsection.

Without Gd, we have to cope not only with electron antineutrinos that induce inverse β decay but also with various sources of Cherenkov radiation. On one hand, the lower energy threshold, E_{min}, will be required to be $\gtrsim 20$ MeV to avoid spallation products and solar neutrinos. As a reference, the number of solar neutrino events in 9.0–9.5 MeV is reported to be 1350 for 0.09 Mt years in Super-Kamiokande [51]. This corresponds to ≈ 1.3 events for 2700 Mt s relevant for detecting a single thermal

TABLE III. Background and the expected number of events for 2700 Mt s relevant for detecting a single thermal neutrino from binary neutron-star mergers (but note that the waiting exposure depends on the energy window via f_E). All the data are taken from Ref. [51], except for neutral-current quasielastic scattering taken from Ref. [60] by assuming that Gd does not change the level of this background. We do not include solar and reactor neutrinos severe at $E \lesssim 10$ MeV.

	Without Gd (20–30 MeV)	With Gd (10–50 MeV)
Decay electron from invisible muons	0.03	0.03
Atmospheric antineutrino	0.003	0.003
Neutral-current quasielastic scattering	$<10^{-3}$	0.003
Diffuse supernova neutrino background	0.004	0.01

neutrino, and the rate for spallation products is higher at least by a factor of 5 than this. Even though solar neutrinos become much weaker at higher energy (see also discussions in Ref. [61]), spallation products will serve as severe contaminants for Hyper-Kamiokande at a shallow site (but see Refs. [62–64] for possible order-of-magnitude reduction informed by shower physics). On the other hand, the higher threshold $E_{\max}$ has to be chosen to avoid decay electrons from invisible muons. The rate of events from invisible muons (and atmospheric electron antineutrinos) is expected to be ≈ 220 for 20–30 MeV and 0.56 Mt years [51]. Thus, it will produce only ≈ 0.03 event for 2700 Mt s, while the number will increase to ≈ 0.08 and 0.15 with $E_{\max} = 40$ MeV and 50 MeV, respectively. Therefore, a reasonable choice may be $E_{\min} \approx 20$ MeV and $E_{\max} \approx 30$ MeV, which results in $f_E \approx 0.2$–0.33 (see Table I). This could be acceptable but is not very satisfactory. In particular, the number of mergers required for detecting a thermal neutrino is proportional to f_E^{-1}, and thus contamination events increase accordingly.

If Gd is dissolved successfully, the situation improves substantially. First, spallation backgrounds as well as solar neutrinos can be suppressed via neutron tagging. The lower threshold may be reduced to $E_{\min} \lesssim 10$ MeV limited by electron antineutrinos from nuclear reactors. A caveat comes from $\beta^- n$ decay of isotopes such as ^{9}Li and ^{8}He, which mimics the inverse β decay. Still, it is suggested that these isotopes could be removed by identifying preceding neutrons produced during the spallation [62], and here we optimistically assume that these isotopes can also be efficiently removed. Because neutrinos with $E < 10$ MeV are minor (see Table I), it will be sufficient if we could reduce $E_{\min}$ to 10 MeV. Next, Gd also reduces the events from invisible muons at high energy by a factor of ~ 5, and this allows us to choose $E_{\max} = 40$–50 MeV with keeping the rate of contamination to be $\lesssim 0.02$–0.04 for 2700 Mt s taking the reduction of selection efficiency, $f_{\rm se}$, into account. These threshold values give us $f_E \approx 0.8$.

Recently, quasielastic scattering of neutrinos by oxygen nuclei via neutral-current interactions has been recognized as a significant source of contamination at low energy (see, e.g., Ref. [65]). Because this interaction could mimic inverse β decay via neutron ejection, Gd cannot be used to suppress this background in a straightforward manner. This contamination could dominate invisible muons at $\lesssim 15$ MeV, and the number of events is estimated to be ≈ 0.003 for 10–50 MeV and 2700 Mt s according to Ref. [60]. We note that this contamination has not yet been studied extensively, and further reduction is discussed for detecting the diffuse supernova neutrino background [65].

One difference from the search of the diffuse supernova neutrino background is that these neutrinos themselves serve as contamination in our search. The event rate is estimated to be ≈ 83 in 10–30 MeV for 0.56 Mt years with Gd dissolution [51], which corresponds to ≈ 0.01 for 2700 s

with a 1 Mt detector, and the contribution from $E > 30$ MeV is minor. While the uncertainty is large, it is not very likely that the realistic diffuse background is very intense taking present nondetection into account.

The expected number of contamination events per single thermal neutrino from binary neutron-star mergers r and the required energy window are summarized as follows (see also Table III). If Gd is not dissolved, we have to choose 20–30 MeV to avoid spallation products at low energy and decay electrons from invisible muons at high energy. Taking the increase of required exposure for a single detection by a factor of 2–3, this will result in $r \sim 0.05$–0.1. If Gd is dissolved, we may be able to achieve $r \approx 0.03$–0.05 adopting 10–40 or 50 MeV. The lower threshold is determined by reactor neutrinos, and the higher threshold is determined by invisible muons now suppressed by a factor of ~ 5.

B. Toward high significance

The contamination event $r \sim 0.03$ for one detection of thermal neutrinos in 80 years is not hopeless but not very comfortable. Straightforward improvement comes from a large detector that will enable us to detect multiple thermal neutrinos. Here, we would like to discuss other directions to reduce the contamination further.

For this purpose, dependence of the number of contamination events on various parameters should be examined. Generally, strong neutrino emission per merger reduces the required number of mergers and increase the significance of detections. Thus, large values of $f_{\rm all}$, $E_{\Delta t}$, and $\langle E \rangle$ reduce the number of contamination events. At the same time, a large number of mergers and a long observing time window will increase the number of contamination events. Specifically, the expected number of contamination events per single thermal neutrino from binary neutron-star mergers is given by

$$r \propto \frac{\Delta t_{\rm obs} D_{\rm eff}^2}{f_{\rm all} E_{\Delta t} \langle E \rangle}. \qquad (9)$$

Here, dependence on $D_{\rm eff}$ is given by competition between the volume $\propto D_{\rm eff}^3$ and the period required for detecting a single neutrino $P \propto D_{\rm eff}^{-1}$ [see Eq. (8)].

One parameter we can actively choose is $\Delta t_{\rm obs}$, which also affects $E_{\Delta t}$. Because the neutrino luminosity is higher in the earlier epoch, focusing on a short time window after the merger is advantageous for increasing the significance. For example, we may be able to choose $\Delta t_{\rm obs} = 0.1$ s while reducing $E_{\Delta t}$ only by a factor of ≈ 3. This gives us $r \approx 0.01$ with an obvious price of increasing the waiting time by the same factor. Accurate numerical simulations of neutrino emission will be helpful to determine an optimal time interval, $\Delta t_{\rm obs}$, and energy thresholds, $\{E_{\min}, E_{\max}\}$, for detecting thermal neutrinos with high significance.

Another parameter we can actively choose is D_{eff}, the distance to which we try to observe thermal neutrinos from binary neutron-star mergers, or equivalently the threshold signal-to-noise ratio for gravitational-wave detections. If we discard distant mergers with weak neutrino emission, the average fluence of neutrinos per merger increases. Again, the price is the increase of the waiting time, P. Because rP^2 is approximately independent of D_{eff}, the number of contamination events can be suppressed relatively efficiently by restricting the range of D_{eff} with only a modest increase of the waiting time. Specifically, the waiting time increases only by $\sqrt{3} \approx 1.7$ when r is reduced from ≈ 0.03 to ≈ 0.01. The reason for this is that we selectively keep nearby binary neutron-star mergers with large fluences of neutrinos. Therefore, we expect that observing binary neutron-star mergers within $D_{\text{eff}} \lesssim 120$ Mpc with $\Delta t_{\text{obs}} \approx 1$ s may be close to the optimal strategy. This consideration on D_{eff} immediately means that high-sensitivity gravitational-wave detectors such as the Einstein Telescope [66] and Cosmic Explorer [67] will not necessarily be helpful to detect thermal neutrinos, because the significance can be kept high only when we focus on nearby mergers.

The level of contamination will be independent of the merger rate, $\mathcal{R}$, because it is irrelevant to the neutrino energy emitted during Δt_{obs}. While the merger rate critically affects the waiting time, P, its uncertainty does not affect the significance of neutrino detections achieved with our strategy. We also do not expect that the value of M_T changes the significance of each detection of a thermal neutrino.

V. PHYSICS IMPLICATION

Even though we may be able to detect only a single thermal neutrino, it offers a unique opportunity to extract various information. In this section, we describe its possible implication to physics and astrophysics of neutrinos.

A. Energy scale of the neutrino emission

We can infer the energy scale of the neutrino emission by counting the number of mergers that we collect to detect a single thermal neutrino. Figure 3 shows the confidence interval of the neutrino energy that typical binary neutron-star mergers emit during Δt_{obs}, namely $E_{\Delta t}$. We may be able to narrow down the energy scale to about an order of magnitude with 68% confidence even accounting for a factor of ≈ 2 uncertainty in the typical energy, $\langle E \rangle$.

While this constraint is crude, this gives us a unique opportunity to characterize the emission of thermal neutrinos from binary neutron-star mergers in a quantitative manner. In particular, this constraint is the most direct route to infer the nature of the merger remnant, e.g., formation of a long-lived neutron star. Emission with $\gtrsim 10^{52}$ erg supports formation of massive neutron stars as a typical
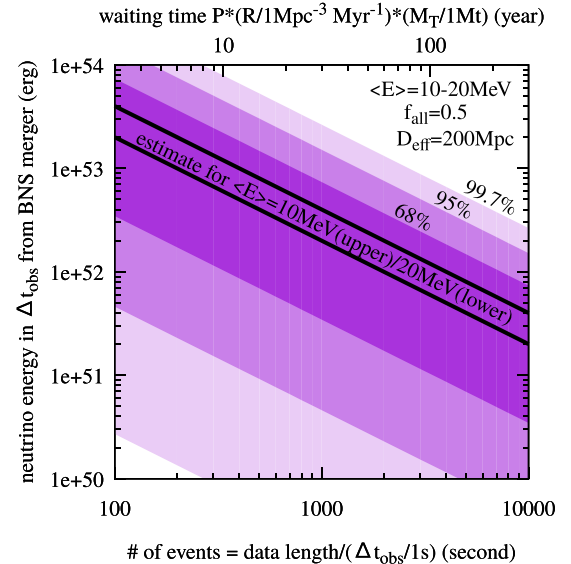

FIG. 3. Confidence interval of the neutrino energy emitted in Δt_{obs} as a function of the number of binary neutron-star mergers observed until the first detection. The upper and lower solid lines denote the estimates under the assumption of $\langle E \rangle = 10$ MeV and 20 MeV, respectively.

remnant of binary neutron stars rather than a prompt collapse, whereas formation of a massive black hole accretion disk with $\gtrsim 0.1\ M_{\odot}$ is not strictly excluded. The upper limit on $E_{\Delta t}$ will reject extraordinary emission of thermal neutrinos. If $E_{\Delta t}$ is smaller than, say 10^{52} erg, it would indicate that either the neutrino emission is relatively weak or the merger remnant rapidly collapses to a black hole. The former will be true if the neutron-star equation of state is very stiff so that the collision does not increase the temperature very much or if the viscosity does not play an important role to heat up the remnant material. The latter should indicate that the maximum mass of the neutron star is not very large, say very close to $2\ M_{\odot}$.

B. Constraining the neutrino mass

Once we detect a thermal neutrino by monitoring multiple mergers, we can identify the progenitor gravitational-wave source from time coincidence in a straightforward manner. One possible concern is that neutrinos may be delayed substantially from gravitational waves due to their finite masses so that the time coincidence becomes loose or completely lost (say, time delay longer than a day). Specifically, the velocity of neutrinos with the momentum p is given by

$$\frac{v}{c} \approx 1 - \frac{m_\nu^2 c^2}{2p^2} \qquad (10)$$

for a neutrino mass $m_\nu \ll p/c \approx E/c^2$. By contrast, because the dispersion relation of gravitational waves, or gravitons, is securely constrained to be that of massless

particles with an upper limit of $\approx 10^{-22}$ eV/c^2 on the corresponding graviton mass [38], we can safely assume that gravitational waves propagate with the speed of light.[4] Thus, the expected time delay of neutrinos relative to gravitational waves is written as

$$\Delta t_d = \left(1 - \frac{v}{c}\right)\frac{D}{c} \tag{11}$$

$$\approx 0.51\,\text{s}\left(\frac{D}{100\,\text{Mpc}}\right)\left(\frac{m_\nu c^2}{0.1\,\text{eV}}\right)^2\left(\frac{E}{10\,\text{MeV}}\right)^{-2}. \tag{12}$$

This expression implies that the time delay can become problematic only for low-energy neutrinos from a very distant merger. Hereafter, the three mass eigenvalues are denoted as m_1, m_2, and m_3. Because m_1 ($\approx m_2$) should be smaller than 0.1 eV/c^2 taking the squared mass difference $|\Delta m_{31}^2| \approx |\Delta m_{32}^2| \approx 2.5 \times 10^{-3}\text{eV}^2/c^4 = (0.05\ \text{eV}/c^2)^2$ [68] and an upper limit on the sum of three mass eigenvalues $\sum_{i=1,2,3} m_i \lesssim 0.2$ eV/c^2 inferred from the *Planck* measurement combined with baryon acoustic oscillation measurements [69], the realistic time delay should be much smaller than the assumed duration of each analysis, $\Delta t_{\text{obs}} \approx 1$ s, particularly for normal hierarchy. Even if the hierarchy is inverted, the dominant part with $m_1 \approx m_2$ is marginally able to produce $\Delta t_d \gtrsim 1$ s with a "worst" combination of parameters. This fact means that the time delay will not substantially degrade the performance of our strategy for neutrino detections. Accordingly, we do not have to worry seriously about the reduction of f_{osc} due to broadening in time of neutrino light curves and corresponding decrease of the flux caused by the mass differences.

Conversely, if we could detect a neutrino and measure its time delay Δt_d relative to the merger, or relative time of flight, we can impose an upper limit on the absolute neutrino mass of the lightest eigenstate[5] from the condition that the mass should not produce time delay longer than Δt_d. Quantitatively, we immediately derive

$$m_\nu c^2 \lesssim \sqrt{\frac{2c\Delta t_d}{D}}E \tag{13}$$

---

[4]It should be cautioned that frequency-independent modification of the speed of gravity cannot be constrained meaningfully from current gravitational-wave observations. We neglect this anomalous case. The constraint obtained by comparing GW170817 and GRB 170817A is not sufficient for our purpose [9].

[5]Because m_3 is coupled only very weakly to the electron-type neutrino [68], we may be able to constrain the mass of $m_1 \approx m_2$, which is not the lightest for inverted hierarchy. We try to be conservative here, because the neutrino oscillation is the source region is not fully understood.

$$\approx 44\ \text{meV}\left(\frac{\Delta t_d}{0.1\ \text{s}}\right)^{1/2}\left(\frac{D}{100\ \text{Mpc}}\right)^{-1/2}\left(\frac{E}{10\ \text{MeV}}\right). \tag{14}$$

Here, we adopt $\Delta t_d = 0.1$ s somewhat optimistically, and we believe that we have a good chance to obtain this value because of the higher luminosity in the earlier epoch [17,19]. In principle, $\Delta t_d \approx 1$ ms can be achieved, where the limitation comes from the timing accuracy of gravitational-wave detectors. Hyper-Kamiokande will determine the arrival time of neutrinos much more accurately than 1 ms.

The measurement error of the distance, D, could degrade the constraint significantly. Gravitational-wave detectors will determine the distance within an error of $\approx 50\%$ [31], and thus the constraint on the neutrino mass will be loosened by $\approx 25\%$. The accuracy of the distance measurement can be improved by an order of magnitude and thus become negligible if we detect electromagnetic counterparts, which can be searched for after prompt identification of a coincident neutrino.

Equation (14) suggests that we might be able to constrain the absolute neutrino mass of the lightest eigenstate to $\lesssim O(10)$ meV/c^2 by detecting a thermal neutrino from binary neutron-star mergers. It would be worthwhile to compare this value with other proposals for constraints. *(i) Supernova.* This limit is tighter by an order of magnitude than eV-scale constraints envisioned for supernova observations both without gravitational waves [70,71] and with gravitational waves [72]. The primary reason of this improvement is that binary neutron stars merge at cosmological distances of $\gtrsim 100$ Mpc, which should be compared with the length scale of our Galaxy, 10 kpc. *(ii) Direct measurements.* The KATRIN experiment is now planning to measure directly the effective mass of electron neutrinos down to $m_{\nu_e,\text{eff}} \approx 0.2$ eV/c^2 via the β decay of the tritium [73], although it is not fair to compare future observations on a time scale of ≈ 30–100 years considered in this study with ongoing experiments. Double β decay experiments can also constrain the mass of neutrinos to sub-eV if they are Majorana particles, but this limit does not apply to Dirac neutrinos (see Ref. [74] for reviews). *(iii) Cosmology.* Our constraint is comparable to current cosmological constraints [69,75] (see also Ref. [76,77]). Taking the potential uncertainty of cosmological models into account (see, e.g., Ref. [78]), the independent constraint from the relative time-of-flight will be invaluable.

VI. SUMMARY

We presented a long-term strategy to detect thermal neutrinos emitted from the remnant of binary neutron-star mergers with a future M-ton water-Cherenkov detector such as Hyper-Kamiokande [51,52]. Although the detection from a single merger is not expected and the diffuse neutron-star-merger neutrino background will be hidden by

other neutrinos, monitoring multiple mergers only for $\Delta t_{\rm obs} \approx 1$ s each by using timing information from gravitational-wave detectors could give us a chance of detection with a human time-scale operation of ≈ 80 Mt years. Contamination from other sources of neutrinos may be reduced to ≈ 0.03 with Gd dissolution. We may be able to reduce the contamination further to ≈ 0.01 with an increase of the waiting time by only a factor of ≈ 1.7 by focusing only on slightly nearby mergers. Ultimately, the chance of detections can be increased by a large effective volume of neutrino detectors, and possible more-than-M-ton class detectors such as Deep-TITAND [28] and MICA [29] will enable us to detect thermal neutrinos from multiple binary neutron-star mergers.

The direct detection will qualitatively confirm the formation of a hot remnant after the merger of binary neutron stars and verify current theoretical pictures. Moreover, the energy scale of the neutrino emission can be constrained from the number of mergers that we collect to detect a single neutrino, and the formation of remnant massive neutron stars could be strongly supported or disfavored. Because distances to binary neutron-star mergers are expected to be cosmological ($\gtrsim 100$ Mpc), we could obtain meaningful upper limits, $\lesssim O(10)$ meV/c^2, on the absolute neutrino mass of the *lightest* eigenstate from the time delay relative to gravitational waves, which are now securely considered to propagate with the speed of light.

APPENDIX: COMPUTATION OF f_Ω

The strain received by a gravitational-wave detector like Advanced LIGO is given by (see, e.g., Ref. [79])

$$h(\theta, \varphi, \iota, \psi) = F_+(\theta, \varphi, \psi) h_+(\iota) + F_\times(\theta, \varphi, \psi) h_\times(\iota).$$
(A1)

The antenna pattern functions $\{F_+, F_\times\}$ depend on the position of the source on the sky (θ, φ) and the so-called

polarization angle ψ that dictates the orientation of the source in the sky plane with respect to the detector as

$$F_+ = \frac{1}{2}(1 + \cos^2\theta) \cos 2\varphi \cos 2\psi - \cos\theta \sin 2\varphi \sin 2\psi,$$
(A2)

$$F_\times = \frac{1}{2}(1 + \cos^2\theta) \cos 2\varphi \sin 2\psi + \cos\theta \sin 2\varphi \cos 2\psi.$$
(A3)

For a quadrupolar gravitational-wave source with the inclination angle ι, we have

$$h_+ = h_0 \frac{1 + \cos^2\iota}{2} \cos(\omega t), \qquad h_\times = h_0 \cos\iota \sin(\omega t),$$
(A4)

where h_0 and ω are the amplitude and frequency of gravitational waves, respectively. To separate geometrical parameters from intrinsic properties of the source, it is useful to define

$$w \equiv \sqrt{\frac{(1 + \cos^2\iota)^2}{4} F_+^2(\theta, \varphi, \psi) + \cos^2\iota F_\times^2(\theta, \varphi, \psi)},$$
(A5)

where $w \leq 1$. The inequality is saturated for face-on binaries ($\iota = 0$ or π) along the normal direction to the detector plane ($\theta = 0$ or π). The horizon distance D_H is defined as the maximal distance at which the signal is detectable with a threshold signal-to-noise ratio and is realized for $w = 1$.

The detectable volume is given in terms of the horizon distance, D_H, by averaging over the binary orientation (ι, ψ) and integrating over the sky position (θ, φ) as

$$V \equiv \frac{1}{4\pi} \int_{\iota\psi} \int_{\theta\varphi} \int_0^{D_H w} r^2 dr d\Omega_{\theta\varphi} d\Omega_{\iota\psi},$$
(A6)

and the effective range of the detector is related to the detectable volume by $V = 4\pi D_{\rm eff}^3/3$. The integral can be evaluated numerically as

$$\frac{1}{(4\pi)^2} \int_{\iota\psi} \int_{\theta\varphi} w^3 d\Omega_{\theta\varphi} d\Omega_{\iota\psi} \approx \frac{1}{(2.26)^3},$$
(A7)

and this shows that $D_{\rm eff} \approx D_H/2.26$. This should be compared with $D_{\rm eff} = D_H$ for a hypothetical case with $w = 1$ in all the directions and orientations.

The neutrino flux and fluence are proportional to D^{-2}, and the average fluence of neutrinos from all the mergers detectable by gravitational waves is given by

$$S_{\text{ave}} = \frac{E_{\Delta t}}{4\pi V} \times \frac{1}{4\pi} \int_{\iota\psi} \int_{\theta\varphi} \int_0^{D_H w} dr d\Omega_{\theta\phi} d\Omega_{\iota\varphi}. \qquad (A8)$$

For a hypothetical case with $w = 1$ in all the directions and orientations, this gives us $S_{\text{ave}} = E_{\Delta t} D_H / V = E_{\Delta t} D_{\text{eff}} / V$. For the realistic antenna pattern, we numerically obtain

$$\frac{1}{(4\pi)^2} \int_{\iota\psi} \int_{\theta\varphi} w d\Omega_{\theta\varphi} d\Omega_{\iota\psi} \approx 0.352, \qquad (A9)$$

and the average fluence of neutrinos is found to be reduced by a factor of $0.352 \times 2.26 \approx 0.797$ compared to the hypothetical case for a given value of D_{eff}.

---

[1] B. P. Abbott, R. Abbott, T. D. Abbott, F. Acernese, K. Ackley, C. Adams, T. Adams, P. Addesso, R. X. Adhikari, V. B. Adya *et al.*, Phys. Rev. Lett. **119**, 161101 (2017).

[2] B. P. Abbott, R. Abbott, T. D. Abbott, F. Acernese, K. Ackley, C. Adams, T. Adams, P. Addesso, R. X. Adhikari, V. B. Adya *et al.*, Astrophys. J. **848**, L12 (2017).

[3] B. D. Metzger and E. Berger, Astrophys. J. **746**, 48 (2012).

[4] T. Piran, E. Nakar, and S. Rosswog, Mon. Not. R. Astron. Soc. **430**, 2121 (2013).

[5] B. Paczyński, Astrophys. J. **308**, L43 (1986).

[6] J. Goodman, Astrophys. J. **308**, L47 (1986).

[7] D. Eichler, M. Livio, T. Piran, and D. N. Schramm, Nature (London) **340**, 126 (1989).

[8] R. Narayan, B. Paczyński, and T. Piran, Astrophys. J. **395**, L83 (1992).

[9] B. P. Abbott, R. Abbott, T. D. Abbott, F. Acernese, K. Ackley, C. Adams, T. Adams, P. Addesso, R. X. Adhikari, V. B. Adya *et al.*, Astrophys. J. **848**, L13 (2017).

[10] K. P. Mooley *et al.*, Nature (London) **554**, 207 (2018).

[11] J. J. Ruan, M. Nynka, D. Haggard, V. Kalogera, and P. Evans, Astrophys. J. **853**, L4 (2018).

[12] J. M. Lattimer and D. N. Schramm, Astrophys. J. **192**, L145 (1974).

[13] E. Symbalisty and D. N. Schramm, Astrophys. Lett. **1982**, 143 (1982).

[14] L.-X. Li and B. Paczyński, Astrophys. J. **507**, L59 (1998).

[15] M. Ruffert, H.-T. Janka, K. Takahashi, and G. Schäfer, Astron. Astrophys. **319**, 122 (1997).

[16] S. Rosswog and M. Liebendörfer, Mon. Not. R. Astron. Soc. **342**, 673 (2003).

[17] Y. Sekiguchi, K. Kiuchi, K. Kyutoku, and M. Shibata, Phys. Rev. Lett. **107**, 051102 (2011).

[18] O. Just, M. Obergaulinger, H.-T. Janka, A. Bauswein, and N. Schwarz, Astrophys. J. **816**, L30 (2016).

[19] S. Fujibayashi, Y. Sekiguchi, K. Kiuchi, and M. Shibata, Astrophys. J. **846**, 114 (2017).

[20] S. Wanajo, Y. Sekiguchi, N. Nishimura, K. Kiuchi, K. Kyutoku, and M. Shibata, Astrophys. J. **789**, L39 (2014).

[21] S. Goriely, A. Bauswein, O. Just, E. Pllumbi, and H.-T. Janka, Mon. Not. R. Astron. Soc. **452**, 3894 (2015).

[22] J. Abadie, B. P. Abbott, R. Abbott, M. Abernathy, T. Accadia, F. Acernese, C. Adams, R. Adhikari, P. Ajith, B. Allen *et al.*, Classical Quantum Gravity **27**, 173001 (2010).

[23] M. Dominik, E. Berti, R. O'Shaughnessy, I. Mandel, K. Belczynski, C. Fryer, D. Holz, T. Bulik, and F. Pannarale, Astrophys. J. **806**, 263 (2015).

[24] J. F. Beacom, Annu. Rev. Nucl. Part. Sci. **60**, 439 (2010).

[25] K. Hirata *et al.*, Phys. Rev. Lett. **58**, 1490 (1987).

[26] R. M. Bionta *et al.*, Phys. Rev. Lett. **58**, 1494 (1987).

[27] S. Ando, J. F. Beacom, and H. Yüksel, Phys. Rev. Lett. **95**, 171101 (2005).

[28] M. D. Kistler, H. Yüksel, S. Ando, J. F. Beacom, and Y. Suzuki, Phys. Rev. D **83**, 123008 (2011).

[29] S. Böser, M. Kowalski, L. Schulte, N. L. Strotjohann, and V. Markus, Astropart. Phys. **62**, 54 (2015).

[30] S. S. Kimura, K. Murase, P. Mészáros, and K. Kiuchi, Astrophys. J. **848**, L4 (2017).

[31] B. P. Abbott, R. Abbott, T. D. Abbott, M. R. Abernathy, F. Acernese, K. Ackley, C. Adams, T. Adams, P. Addesso, R. X. Adhikari *et al.*, Phys. Rev. X **6**, 041015 (2016).

[32] A. Gando *et al.*, Astrophys. J. **829**, L34 (2016).

[33] K. Abe *et al.*, Astrophys. J. **830**, L11 (2016).

[34] C. S. Kochanek and T. Piran, Astrophys. J. **417**, L17 (1993).

[35] A. Nishizawa and T. Nakamura, Phys. Rev. D **90**, 044048 (2014).

[36] K. Langæble, A. Meroni, and F. Sannino, Phys. Rev. D **94**, 053013 (2016).

[37] T. Baker and M. Trodden, Phys. Rev. D **95**, 063512 (2017).

[38] B. P. Abbott, R. Abbott, T. D. Abbott, F. Acernese, K. Ackley, C. Adams, T. Adams, P. Addesso, R. X. Adhikari, V. B. Adya *et al.*, Phys. Rev. Lett. **118**, 221101 (2017).

[39] K. Hotokezaka, K. Kiuchi, K. Kyutoku, T. Muranushi, Y.-i. Sekiguchi, M. Shibata, and K. Taniguchi, Phys. Rev. D **88**, 044026 (2013).

[40] K. Hotokezaka, K. Kyutoku, H. Okawa, M. Shibata, and K. Kiuchi, Phys. Rev. D **83**, 124008 (2011).

[41] Y. Sekiguchi, K. Kiuchi, K. Kyutoku, and M. Shibata, Phys. Rev. D **91**, 064059 (2015).

[42] C. Palenzuela, S. L. Liebling, D. Neilsen, L. Lehner, O. L. Caballero, E. O'Connor, and M. Anderson, Phys. Rev. D **92**, 044045 (2015).

[43] F. Foucart, R. Haas, M. D. Duez, E. O'Connor, C. D. Ott, L. Roberts, L. E. Kidder, J. Lippuner, H. P. Pfeiffer, and M. A. Scheel, Phys. Rev. D **93**, 044019 (2016).

[44] Y. Sekiguchi, K. Kiuchi, K. Kyutoku, M. Shibata, and K. Taniguchi, Phys. Rev. D **93**, 124046 (2016).

[45] Y.-L. Zhu, A. Perego, and G. C. McLaughlin, Phys. Rev. D **94**, 105006 (2016).

[46] M. Frensel, M.-R. Wu, C. Volpe, and A. Perego, Phys. Rev. D **95**, 023011 (2017).

[47] A. Chatelain and C. Volpe, Phys. Rev. D **95**, 043005 (2017).

[48] M.-R. Wu and I. Tamborra, Phys. Rev. D **95**, 103007 (2017).

[49] S. Richers, D. Kasen, E. O'Connor, R. Fernández, and C. D. Ott, Astrophys. J. **813,** 38 (2015).

[50] H.-T. Janka, in *Handbook of Supernovae*, edited by A. W. Alsabti and P. Murdin (Springer, New York, 2016), p. 1.

[51] K. Abe *et al.*, arXiv:1109.3262.

[52] Hyper-Kamiokande Proto-Collaboration, arXiv:1611.06118.

[53] P. Vogel and J. F. Beacom, Phys. Rev. D **60,** 053003 (1999).

[54] A. Strumia and F. Vissani, Phys. Lett. B **564,** 42 (2003).

[55] J. F. Beacom and M. R. Vagins, Phys. Rev. Lett. **93,** 171101 (2004).

[56] P. F. Harrison, D. H. Perkins, and W. G. Scott, Phys. Lett. B **530,** 167 (2002).

[57] O. L. Caballero, G. C. McLaughlin, and R. Surman, Phys. Rev. D **80,** 123004 (2009).

[58] T. Liu, B. Zhang, Y. Li, R.-Y. Ma, and L. Xue, Phys. Rev. D **93,** 123004 (2016).

[59] R. V. Novoseltseva, M. M. Boliev, I. M. Dzaparova, M. M. Kochkarov, Y. F. Novoseltsev, V. B. Petkov, V. I. Volchenko, G. V. Volchenko, and A. F. Yanin, Phys. Part. Nucl. **47,** 968 (2016).

[60] H. Sekiya, J. Phys. Conf. Ser. **888,** 012041 (2017).

[61] S. Ando, K. Sato, and T. Totani, Astropart. Phys. **18,** 307 (2003).

[62] S. W. Li and J. F. Beacom, Phys. Rev. C **89,** 045801 (2014).

[63] S. W. Li and J. F. Beacom, Phys. Rev. D **91,** 105005 (2015).

[64] S. W. Li and J. F. Beacom, Phys. Rev. D **92,** 105033 (2015).

[65] A. Priya and C. Lunardini, J. Cosmol. Astropart. Phys. 11 (2017) 031.

[66] B. Sathyaprakash *et al.*, Classical Quantum Gravity **29,** 124013 (2012).

[67] B. P. Abbott, R. Abbott, T. D. Abbott, M. R. Abernathy, K. Ackley, C. Adams, P. Addesso, R. X. Adhikari, V. B. Adya, C. Affeldt *et al.*, Classical Quantum Gravity **34,** 044001 (2017).

[68] F. Capozzi, E. Di Valentino, E. Lisi, A. Marrone, A. Melchiorri, and A. Palazzo, Phys. Rev. D **95,** 096014 (2017).

[69] Planck Collaboration, Astron. Astrophys. **596,** A107 (2016).

[70] T. Totani, Phys. Rev. Lett. **80,** 2039 (1998).

[71] J. F. Beacom, R. N. Boyd, and A. Mezzacappa, Phys. Rev. Lett. **85,** 3568 (2000).

[72] N. Arnaud, M. Barsuglia, M.-A. Bizouard, F. Cavalier, M. Davier, P. Hello, and T. Pradier, Phys. Rev. D **65,** 033010 (2002).

[73] R. G. Hamish Robertson, arXiv:1307.5486.

[74] F. T. Avignone III, S. R. Elliott, and J. Engel, Rev. Mod. Phys. **80,** 481 (2008).

[75] N. Palanque-Delabrouille *et al.*, J. Cosmol. Astropart. Phys. 11 (2015) 011.

[76] E. Giusarma, M. Gerbino, O. Mena, S. Vagnozzi, S. Ho, and K. Freese, Phys. Rev. D **94,** 083522 (2016).

[77] S. Vagnozzi, E. Giusarma, O. Mena, K. Freese, M. Gerbino, S. Ho, and M. Lattanzi, Phys. Rev. D **96,** 123503 (2017).

[78] A. G. Riess, L. M. Macri, S. L. Hoffmann, D. Scolnic, S. Casertano, A. V. Filippenko, B. E. Tucker, M. J. Reid, D. O. Jones, J. M. Silverman, R. Chornock, P. Challis, W. Yuan, P. J. Brown, and R. J. Foley, Astrophys. J. **826,** 56 (2016).

[79] B. S. Sathyaprakash and B. F. Schutz, Living Rev. Relativity **12,** 2 (2009).