High Energy Neutrinos from Choked Gamma-Ray Bursts in AGN Accretion Disks

JIN-PING ZHU,1 KAI WANG,2 BING ZHANG,3 YUAN-PEI YANG,4 YUN-WEI YU,5 AND HE GAO6

1Department of Astronomy, School of Physics, Peking University, Beijing 100871, China; zhujp@pku.edu.cn
2Department of Astronomy, School of Physics, Huazhong University of Science and Technology, Wuhan 430074, China
3Department of Physics and Astronomy, University of Nevada, Las Vegas, NV 89154, USA; zhang@physics.unlv.edu
4South-Western Institute for Astronomy Research, Yunnan University, Kunming, Yunnan, People’s Republic of China
5Institute of Astrophysics, Central China Normal University, Wuhan 430079, China
6Department of Astronomy, Beijing Normal University, Beijing 100875, China

ABSTRACT

Both long-duration gamma-ray bursts (LGRBs) from core collapse of massive stars and short-duration GRBs (SGRBs) from mergers of binary neutron star (BNS) or neutron star–black hole (NSBH) are expected to occur in the accretion disk of active galactic nuclei (AGNs). We show that GRB jets embedded in the migration traps of AGN disks are promised to be choked by the dense disk material. Efficient shock acceleration of cosmic rays at the reverse shock is expected, and high-energy neutrinos would be produced. We find that these sources can effectively produce detectable TeV–PeV neutrinos through $p\gamma$ interactions. From a choked LGRB jet with isotropic equivalent energy of $10^{53}$ erg at 100 Mpc, one expects $\sim 2 (7)$ neutrino events detectable by IceCube (IceCube-Gen2). The contribution from choked LGRBs to the observed diffuse neutrino background depends on the unknown local event rate density of these GRBs in AGN disks. For example, if the local event rate density of choked LGRBs in AGN disk is $\sim 5\%$ that of low-luminosity GRBs ($\sim 10$ Gpc$^{-3}$ yr$^{-1}$), the neutrinos from these events would contribute to $\sim 10\%$ of the observed diffuse neutrino background. Choked SGRBs in AGN disks are potential sources for future joint electromagnetic, neutrino, and gravitational wave multi-messenger observations.

Keywords: Cosmological neutrinos (338); Gamma-ray bursts (629); Neutron stars (1108); Black holes (16); Active galactic nuclei (16); Gravitational waves (678)

1. INTRODUCTION

Massive stars can be born in situ in an AGN accretion disk or be captured from the nuclear star cluster around the AGN (e.g., Artymowicz et al. 1993; Collin & Zahn 1999; Wang et al. 2011; Fabj et al. 2020; Dittmann et al. 2021). When they die, these massive stars will make supernovae (SNe) and leave behind NS and BH remnants inside the disk. Most of them might have extreme high spins and easily make LGRBs (Jernyn et al. 2021). Such embedded massive stars and stellar remnants would migrate inwards the trapping orbits within the disk (e.g., Bellovary et al. 2016; Tagawa et al. 2020). Abundant compact objects within the orbit would likely collide and merge (Cheng & Wang 1999; McKernan et al. 2020), which are the promising astrophysical gravitational wave (GW) sources for LIGO (Abbott et al. 2009). A possible SN (Assef et al. 2018) and a candidate binary black hole merger induced electromagnetic transient (Graham et al. 2020) have been reported recently in association with AGN disks.

GRBs, both long-duration ones associated with core collapse of massive stars (Woosley 1993; MacFadyen & Woosley 1999; Zhang et al. 2003; Zhang & Mészáros 2004; Galama et al. 1998; Hjorth et al. 2003; Stanek et al. 2003) and short-duration ones associated with neutron star mergers (Paczynski 1986, 1991; Eichler et al. 1989; Narayan et al. 1992; Abbott et al. 2017a,b,c) have been suggested as sources of astrophysical high-energy cosmic rays and neutrinos (Waxman & Bahcall 1997). On the other hand, non-detection of GRB neutrino signals (Icecube Collaboration et al. 2012; Aartsen et al. 2015a), likely related to a large emission radius from the central engine (Zhang & Kumar 2013), suggested GRB-associated neutrinos can only account for at most $\lesssim 1\%$ of the diffuse neutrino fluence (Murase 2008; Wang & Dai 2009). The so-called low-luminosity GRBs are more abundant than successful ones (Liang et al. 2007; Sun et al. 2015) and can give significant contribution to the neutrino background (Murase & Nagataki 2006; Gupta & Zhang 2007). One possibility is that they originate...
from massive stars that launched jets that are choked in the stellar envelope (Bromberg et al. 2011a; Nakar 2015). Such choked jets from the death of massive stars or even from neutron star mergers could be more promising sites for efficient neutrino production, which may contribute to a considerable fraction of the diffuse neutrino background (Murase & Ioka 2013; Senno et al. 2016; Xiao & Dai 2014; Kimura et al. 2018; Fasano et al. 2021).

The presence of massive stars and compact binaries in AGN disks indicates that both LGRBs and SGRBs could possibly occur in such a high density environment. Very recently, Zhu et al. (2021); Perna et al. (2021) suggested that GRBs jet in the AGN disks are likely choked\(^1\), instead of making and resulting in observable signals from optical to \(\gamma\)-ray. In this work, we consider GRB jets choked inside the AGN disks as hidden sources of high energy cosmic neutrinos, which can ease the tension between the diffuse extragalactic gamma-ray background and the diffuse background of TeV–PeV neutrinos.

2. JET DYNAMICS IN AGN ACCRETION DISKS

SNe and neutron star mergers are expected to occur in the migration traps (Bellovary et al. 2016) plausibly located at \( a \sim 10^3 r_g \), where \( r_g \equiv GM_{\text{BH}}/c^2 \), \( G \) is the gravitational constant, \( M_{\text{BH}} \) is the AGN supermassive BH mass, and \( c \) is the speed of light. For a gas-pressure-dominated disk, the disk density is \( \rho(z) = \rho_0 \exp(-z/H) \), where \( \rho_0 \) is the midplane density, \( z \) is the vertical distance, and \( H = 1.5 \times 10^{14} M_{\text{BH},8}\odot \left(\frac{H/a}{0.01}\right) \left(\frac{a}{10^4 r_g}\right) \) is the disk height with the disk aspect ratio \( H/a \sim (10^{-3}, 0.1) \) (Goodman & Tan 2004; Thompson et al. 2005). Near the migration traps, the midplane density is almost \( \rho_0 \sim O(10^{-10}) \) g cm\(^{-3}\), and the disk height is \( H \sim O(10^{14}) \) cm. For a disk with an exponentially decaying density profile, the density \( \rho \approx \rho_0 \) for \( z < H \) while \( \rho \) decreases rapidly for \( z > H \). Therefore, for simplicity we approximately adopt a uniform density profile for the AGN disk for \( z < H \). Hereafter, the convention \( Q_x = Q/10^x \) is adopted in cgs units.

Figure 1 illustrates the physical processes for a jet traveling through the progenitor star and the AGN accretion disk. When a jet initially propagates inside the progenitor star, the collision between the jet and the stellar gas medium leads to the formation of a forward shock sweeping into the medium and a reverse shock entering the jet material (e.g., Matzner 2003; Bromberg et al. 2011b; Yu 2020). Such a structure is known as the jet head. The hot material that enters the head flows sideways and produces a powerful cocoon to drive a collimation shock into the jet material. The jet is collimated inside the star and gets accelerated to a relativistic velocity after it breaks out from the progenitor star.

After the jet entering into the AGN disk, the jet head velocity is given by (Matzner 2003): \( \beta_h = \beta_j/(1 + L^{-1/2}) \), where \( \beta_j \approx 1 \) and \( L \equiv L_\text{iso}/\pi r_j^2 \rho_0 c^3 \approx (10^{14} \text{cm}/r_j)^2 L_{50} \rho_0^{-1} \approx 10^{33} \) is the critical parameter that determines the evolution of the jet (Bromberg et al. 2011b), \( L_j \) is the jet luminosity, and \( r_j \) is the radius of the jet from the central engine. Since this jet radius should be \( r_j < 10^{14} \) cm, one gets \( L > \theta_0^{-4/3} \approx 1 \), where \( \theta_0 \approx 0.2 \). In this case, the jet head would travel with a relativistic speed. The cocoon pressure is too weak to affect the geometry of the jet so that the jet is uncollimated (i.e., \( \theta_j \approx \theta_0 \))(Bromberg et al. 2011b).

Therefore, the critical parameter can be expressed as \( \bar{L} = L_{\text{iso}}/2\pi r_j^2 \rho_0 c^3 \), where \( L_{\text{iso}} \approx 2L_j/\theta_0^2 \) is the isotropic-equivalent one-side jet luminosity and \( r_h \approx r_j/\theta_0 \) is the distance between the jet head and the central engine. The Lorentz factor of the jet head is given by \( \Gamma_h \approx \bar{L}^{1/4}/\sqrt{2} \). The internal energy and density evolution of the forward shock and reverse shock can be described by the shock jump conditions (Blandford & McKee 1976; Sari & Piran 1995; Zhang 2018):

\[
e'_\ell/n'_t m_p c^2 = \Gamma_h - 1, \quad n'_t/n_h = (\dot{\gamma}_1 \Gamma_h + 1)/(\dot{\gamma}_1 - 1) = 4\Gamma_h, \\
e'_r/n'_t m_p c^2 = \bar{\Gamma}_h - 1, \quad n'_t/n'_h = (\dot{\gamma}_2 \bar{\Gamma}_h + 1)/(\dot{\gamma}_2 - 1) = 4\bar{\Gamma}_h, 
\]

where \( \dot{\gamma}_1 = (4\Gamma_h + 1)/3\Gamma_h, \dot{\gamma}_2 = (4\bar{\Gamma}_h + 1)/3\bar{\Gamma}_h \), the subscripts “\( a \)”, “\( \ell \)”, “\( r \)” and “\( j \)” represent regions of the

\(^1\) Kimura et al. (2021) recently presented that compact binaries can accrete, produce radiation-driven outflows, and create cavities in the AGN disks before the merger, so that aligned SGRB jets could successfully break out from the AGN disks.
unshocked AGN material, the jet head’s forward shock, the jet head’s reverse shock, and the unshocked jet, and \( \hat{\Gamma}_h = \Gamma_j / (1 - \beta_j \beta_h) \approx \Gamma_j \hat{L}^{-1/4}/\sqrt{2} \) is the Lorentz factor of the unshocked jet measured in the jet head frame. Note we distinguish among three reference frames: \( Q \) for the AGN rest frame, \( Q' \) for the jet head comoving frame, and \( Q'' \) for the jet comoving frame.

We assume that a luminous jet can easily break out from the progenitor star, which has a similar parameter distribution as classical GRBs. Observationally, the average duration for LGRBs is \( t_s \approx 10^{1.5} \, \text{s} \) (Kouveliotou et al. 1993; Horváth 2002) while the median isotropic energy release is \( E_{\text{iso}} \approx 10^{53} \, \text{erg} \) for LGRB jets (the isotropic equivalent luminosity \( L_{\text{iso}} = E_{\text{iso}}/t_s \)) (Kumar & Zhang 2015). The jet could be choked in the AGN disk when the central engine quenches so that the jet head radius at \( t_j \) can be defined as the stalling radius, i.e., \( r_{\text{stall}} = r_h(t_j) \approx 2 \Gamma_h^2 c t_j = 2.0 \times 10^{13} E_{\text{iso},53}^{1/4} \rho_0^{-1/4} \, \text{cm} \), with \( \Gamma_h = 2.7 E_{\text{iso},53}^{-3/8} \rho_0^{-1/8} \). For SGRB jets, the average duration and the median isotropic energy are \( t_j \approx 10^{-0.1} \, \text{s} \) and \( E_{\text{iso}} \approx 10^{53} \, \text{erg} \) (Fong et al. 2015), respectively. The stalling radius for SGRB jet is \( r_{\text{stall}} = 2.5 \times 10^{12} E_{\text{iso},51j}^{1/4} \rho_0^{-1/4} \, \text{cm} \) and the jet head Lorentz factor when it chokes is \( \Gamma_h = 6.1 E_{\text{iso},51j}^{-3/8} \rho_0^{-1/8} \). Because \( r_{\text{stall}} \ll H \), both LGRB and SGRB jets can be easily choked in an AGN accretion disk.

### 3. Neutrino Production

We assume that the Fermi acceleration operates and accelerated protons have a power-law distribution in energy: \( dn_p/d\epsilon_p \propto \epsilon_p^{-s} \) with \( s = 2 \) (Achterberg et al. 2001; Keshet & Waxman 2005). Here, the thermal photons from the jet head are treated as the only background photon field for hadronic interactions since the number densities of other types of radiation (e.g., the classical keV–MeV emission of GRBs) are typically much lower (Senni et al. 2016). Based on Eq. (1), the internal energy and proton energy density of the jet head can be expressed as \( \epsilon_p' = (\hat{\Gamma}_h - 1) n'_p m_p c^2 \) and \( n'_p = 4 \hat{\Gamma}_h n''_p \), where \( m_p \) is the proton mass and \( n''_p = L_{\text{iso}}/4 \pi \Gamma_j^{3/2} m_p c^3 \) is the jet density. The photon temperature of the jet head is \( k_B T'_\gamma = (15 \epsilon_p'^3 c \epsilon_p'/\pi^2)^{1/4} \), where \( \epsilon_p' \approx 0.1 \) is the electron energy fraction. For classical parameters of LGRB (SGRB) jets, \( k_B T'_\gamma = 0.21(0.32) \, \text{keV} \). One can obtain the average thermal photon energy \( \epsilon'_\gamma = 2.7 k_B T'_\gamma \) and the average thermal photon density \( n'_\gamma = \epsilon'_\gamma \epsilon'_p/n'_p \).

For the reverse shock, efficient Fermi acceleration can occur only if the radiation constraint (Murase & Ioka 2013) is satisfied, i.e., \( \tau_T = n'_p \sigma_T r_{\text{stall}} / \Gamma_j \lessgtr \min[1, 0.1 C^{-1} \Gamma_j] \), where \( C = 1 + 2 \ln \Gamma_h^2 \). By considering \( \Gamma_j = 300 \) (500) for LGRB (SGRB) jets, we get \( \tau_T = 6.8 \times 10^{-3} E_{\text{iso},53}^{3/4} T_{\gamma,j}^{-5/4} \rho_0^{-10} \leq 0.32 C_{1,2}^{-1} \Gamma_h,1.7 \leq 1 \) (\( \tau_T = 4.6 \times 10^{-3} E_{\text{iso},51j}^{-5/4} \rho_0^{-10} \Gamma_j^{-3} \leq 0.26 C_{1,2}^{-1} \Gamma_h,1.6 \leq 1 \)), which means that Fermi acceleration is always effective. Note that different from Senni et al. (2016) whose protons are from the internal shocks, the interacted protons in our calculations are accelerated from the reverse shock.

With the assumption of perfectly efficient acceleration, the acceleration timescale is given by \( t_{p,\text{acc}} = \epsilon'_p/(e B' c) \), where the jet head comoving magnetic field strength is \( B' = \sqrt{8 \pi \epsilon_B ' c} \approx (\hat{\Gamma}_h / \Gamma_j)(8 \epsilon_B L_{\text{iso}} / r_{\text{stall}} c)^{1/2} \) and the magnetic field energy fraction is assumed as \( \epsilon_B = 0.1 \).

A high energy proton loses its energy through radiative, hadronic, and adiabatic processes. The radiative cooling mechanisms contain synchrotron radiation with cooling timescale

\[
t'_{p,\text{syn}} = \frac{6 m_p^4 c^3}{\sigma_T m_e^2 B'^2} \epsilon_p',
\]

and inverse-Compton scattering with cooling timescale

\[
t'_{p,\text{lC}} = \left\{ \begin{array}{ll}
3 m_p c^3 & \epsilon_p' < m_p^2 c^4, \\
\frac{3 c \epsilon_p'}{4 \pi m_e^2 c^5 n'_\gamma} & \epsilon_p' > m_p^2 c^4.
\end{array} \right.
\]

The hadronic cooling mechanisms mainly include the inelastic hadronuclear scattering (pp), the Bethe-Heitler pair production (\( p\gamma \rightarrow e^+ e^- \)), and the photomeson production (\( p\gamma \)). High energy neutrinos are expected to be produced via pp and p\gamma processes. The cooling timescale of pp scattering is given by \( t_{p,\text{pp}} = 1/\epsilon_p^2 n'_p k_{\text{pp}} \), where the inelasticity is set as \( k_{\text{pp}} \approx 0.5 \) and the cross section \( \sigma_{\text{pp}} \) is obtained from Kelner et al. (2006). The energy loss rate of \( p\gamma \) production is calculated by the formula given in Stecker (1968); Murase (2007), i.e.,

\[
t'_{p,\text{p}\gamma} = \frac{c}{2 n_p} \int_{\epsilon_{\text{th}}}^{\infty} d\epsilon \sigma_{\gamma p}(\epsilon) \epsilon \epsilon_{\gamma p}(\epsilon) \int_{\epsilon/2\epsilon_{p}}^{\infty} d\epsilon' e^{-2dn/d\epsilon} d\epsilon,
\]

where \( \epsilon_{\text{th}} \) represents the photon energy in the rest frame of the proton, \( \epsilon_{\text{th}} \approx 145 \, \text{MeV} \) is the threshold energy, \( \gamma_p = \epsilon_p'/m_p c^2 \), and \( dn/d\epsilon \) is the photon number density in the energy range of \( \epsilon \) to \( \epsilon + d\epsilon \). The inelasticity \( \kappa_{\gamma p} \) and the cross section \( \sigma_{\gamma p} \) are taken from Stecker (1968); Patrignani et al. (2016). The energy loss rate of Bethe-Heitler process \( t'_{\text{BHT}}^{-1} \) can be also estimated based on Eq. (4) by using \( \kappa_{\text{BH}} \) and \( \sigma_{\gamma p} \) instead of \( \kappa_{\gamma p} \) and \( \sigma_{\gamma p} \).
the acceleration and cooling timescales of a choked LGRB and SGRB jet in an AGN disk in Figure 2. For both cases, pp scattering would dominate the cooling process for low-energy protons. Bethe-Heitler process leads and suppresses neutrino production if the energy of protons falls within the range of \(0.5 \text{ TeV} \lesssim \epsilon_p' \lesssim 20 \text{ TeV}\). At higher energies, the dominant cooling mechanism for protons is \(p\gamma\) interaction, which also limits the maximum proton energy to \(\epsilon_{p,\max} \sim 10 \text{ PeV}\).

Pions and kaons created through pp and p\gamma processes decay into muons and muon neutrinos. Pions and kaons are subject to hadronic scattering, \(t_{i,\text{had}}' = 1/c\sigma_{i,K}\text{had}n'_p\kappa_h\), where \(\sigma_{i,K}\) \(\approx 5 \times 10^{-26} \text{ cm}^2\) and \(\kappa_h \approx 0.8\) (Olive & Particle Data Group 2014). The intermediate muons then decay to muon neutrinos, electron neutrinos and electrons. Similar to protons, pions, kaons, and muons also experience radiative processes and adiabatic cooling. One can calculate the synchrotron and IC cooling timescales of pions, kaons, and muons by Eq. (2) and Eq. (3) with \(\epsilon_p' \rightarrow \epsilon_i'\) and \(m_p \rightarrow m_i\), where \(i = \pi, K, \mu\), and \(\mu_{\pi, K}\) are the parent particles for the neutrinos. The energy fractions from a proton to intermediate particles are calculated according to Denton & Tamborra (2018). By comparing these cooling timescales with the decay timescales of intermediate particles, i.e., \(t_{i,\text{dec}} = \gamma_i\tau_i\) (where \(\gamma_i = \epsilon_i'/m_i c^2\) and \(\tau_i\) are the Lorentz factor and the rest frame lifetime, respectively), the final neutrino spectrum can be obtained.

Since both pp and p\gamma processes produce neutrinos while other proton cooling processes suppress the final neutrino spectrum, the suppression factor taking into account various proton cooling processes is expressed as (Murase 2008; Wang & Dai 2009; Xiao et al. 2016) \(\eta_{p,\text{sup}}(\epsilon_{\nu_i}) = \left(t_{i,\text{opt}}^{-1} + t_{i,\text{sup}}^{-1}/t_{i,\text{cool}}^{-1}\right)^{-1}\), where \(t_{i,\text{opt}}^{-1} = t_{i,\text{dec}}^{-1} + t_{i,\text{had}}^{-1} + t_{i,\text{syn}}^{-1} + t_{i,\text{IC}}^{-1} + t_{i,\text{ada}}^{-1}\). Similarly, the suppression factor due to meson cooling can be written as \(\eta_{\nu_i}(\epsilon_{\nu_i}) = t_{i,\text{dec}}^{-1}/t_{i,\text{cool}}^{-1}\), where \(t_{i,\text{cool}}^{-1} = t_{i,\text{dec}}^{-1} + t_{i,\text{had}}^{-1} + t_{i,\text{syn}}^{-1} + t_{i,\text{IC}}^{-1} + t_{i,\text{ada}}^{-1}\). One can obtain the neutrino spectrum in each neutrino production channel for a single event,

\[
\epsilon_{\nu_i}^2F_{\nu_i} = \frac{N_{\nu_i}\sigma_{\text{iso}}\epsilon_{\nu_i}^{0.6}\eta_{p,\text{sup}}(\epsilon_{\nu_i})\eta_{\nu_i}(\epsilon_{\nu_i})}{4\pi D_L^2\ln(\epsilon_{\nu_i}^{\max}/\epsilon_{\nu_i}^{\min})},
\]

where \(N_\pi = N_{\mu_\pi} = 0.12\), \(N_K = 0.009\), and \(N_{\mu_K} = 0.003\), the neutrino energy is \(\epsilon_{\nu_i} = a_{\nu_i}\Gamma_\nu\epsilon_p'\) with \(a_{\nu_\pi} = 0.05\), \(a_K = 0.10\), and \(a_{\mu_K} = 0.033\), where \(D_L\) is the luminosity distance, and \(\ln(\epsilon_{\nu_i}^{\max}/\epsilon_{\nu_i}^{\min})\) is the normalized factor with \(\epsilon_{\nu_i}^{\min} \approx \Gamma_1 m_n c^2\). Highly efficient acceleration is assumed here (i.e., the acceleration efficiency \(\zeta_p \approx 1\)) so that \(\zeta_{\text{acc}} \approx \zeta_p(1 - \epsilon_\gamma - \epsilon_B) \approx 0.8 \sim 1\) which is in accord with the fiducial value of the baryon loading parameter \(\zeta_{\text{acc}} \approx \epsilon_{\text{acc}}/\epsilon_e \approx 10\) (Murase 2007).

Figure 3 shows the all-flavor fluence of a single burst at \(D_L = 100\) Mpc. The fluence is mainly determined by the isotropic energy. The dip around a few TeV is caused by the suppression of neutrino production due to the Bethe-Heitler process. Low-energy neutrinos are dominated by pp interactions, and the neutrino spectrum above \(\sim 1 \text{ TeV}\) that we are interested in mainly attributes to \(p\gamma\) interactions. Both pp and p\gamma processes are efficient as shown in Figure 3, since both the pp optical depth \(n_p'\sigma_{pp}\tau_{\text{opt}}\) and the p\gamma optical depth \(n_p'\sigma_{p\gamma}\tau_{\text{opt}}\) for a classical LGRB (SGRB) are quite large (Murase 2008) (consid-
The effective areas of IceCube for the up- and down-going events from Aartsen et al. (2017). The effective volume of the Gen2 is larger than that of IceCube by a factor of $\sim 10$, corresponding to a factor of $\sim 10^{2/3}$ times larger in the effective area (Aartsen et al. 2017). The number of detected $\nu_\mu$ (e.g., Harrison et al. 2002) from a single event located at 100 Mpc are shown in Table 1 (after considering neutrino oscillation). If a classical LGRB occurs in an AGN disk at 100 Mpc, we expect $\sim 2 \times (7)$ neutrino events from a single event detected by IceCube (Gen2). The neutrino flux from a SGRB is lower, and the detection for a single choked SGRB is possible only with Gen2 given that the SGRB has a high energy and occurs in the Northern Hemisphere.

We simulate the joint GW+neutrino detection rate for neutron star mergers occurring in AGN disks. McKernan et al. (2020) showed that the local event rate densities $\rho_0$ for BNS and NSBH mergers in the AGN channel are $\rho_0_{\text{BNS}} \sim f_{\text{AGN}} [0.2, 400] \text{Gpc}^{-3} \text{yr}^{-1}$ and $\rho_0_{\text{NSBH}} \sim f_{\text{AGN}} [10, 300] \text{Gpc}^{-3} \text{yr}^{-1}$, respectively, where $f_{\text{AGN}}$ is the fraction of the observed BBH in the AGN channel. The redshift distribution $f(z)$ we adopt is a model invoking a log-normal delay timescale distribution with respect to star formation history, as applied to the study of SGRBs (Wanderman & Piran 2015; Virgili et al. 2011; Sun et al. 2015). The GW horizon distances for BNS and NSBH mergers in different eras are obtained from (Abbott et al. 2018; Maggiore et al. 2020; Hild et al. 2011; Zhu et al. 2020). By considering the beaming correction factor $f_b \approx (\theta_0 + 1/\Gamma_b)^2/2$ and assuming that all BNS and 20% NSBH mergers (McKernan et al. 2020) in AGN disks can power classical SGRBs, we show the results of joint GW+neutrino detection rates in Table 1. Such joint detections appear difficult with the current GW and neutrino detectors, but could be possible in the future with next generation GW and neutrino detectors.

5. NEUTRINO DIFFUSE EMISSION

The diffuse neutrino fluence can be estimated as (e.g., Razzaque et al. 2004)

$$\epsilon^2_{\nu,\text{obs}} \Phi_{\nu,\text{obs}} = \epsilon^2_{\nu,\text{obs}} \int_0^{2}\! dz \rho_0(z) F_{\nu}(\epsilon_{\nu,\text{obs}}) \frac{dV}{dz},$$

where $\epsilon_{\nu,\text{obs}} = \epsilon_{\nu}/(1+z)$, and $dV/dz = 4\pi D_L^2 c/(1+z)|dt/dz|$ is the comoving volume element with $(dt/dz)^{-1} = -H_0(1+z)\sqrt{\Omega_\Lambda + \Omega_m(1+z)^3}$. The standard $\Lambda$CDM cosmology with $H_0 = 67.8 \text{km s}^{-1} \text{Mpc}^{-1}$, $\Omega_{\Lambda} = 0.692$, and $\Omega_m = 0.308$ (Planck Collaboration et al. 2016) is applied.

In view that the energy of a typical SGRB is much smaller than that of a LGRB and that their event rate density may not be much greater than that of long

---

### Table 1. Neutrino Bursts Detection

| Model ($E_{\text{iso}}/\text{erg}$) | IceCube (Up) | IceCube (Down) | Gen2(Up) |
|---------------------------------|--------------|----------------|----------|
| LGRB ($10^{52}$)               | 0.16         | 0.016          | 0.76     |
| LGRB ($10^{53}$)               | 1.5          | 0.15           | 6.8      |
| LGRB ($10^{54}$)               | 13           | 1.4            | 61       |
| SGRB ($10^{50}$)               | $1.2 \times 10^{-3}$ | $1.4 \times 10^{-4}$ | $5.7 \times 10^{-3}$ |
| SGRB ($10^{51}$)               | 0.011        | $1.3 \times 10^{-3}$ | 0.050    |
| SGRB ($10^{52}$)               | 0.094        | 0.012          | 0.44     |

### Joint GW + Neutrino Detection Rate

| Era/GW/Neutrino | Detection Rate (yr$^{-1}$) |
|-----------------|-----------------------------|
| O4/IceCube      | $f_{\text{AGN}}[0.006, 1.3] \times 10^{-1}$ |
| O5/IceCube      | $f_{\text{AGN}}[0.007, 1.3] \times 10^{-1}$ |
| Voyager/Gen2    | $f_{\text{AGN}}[0.041, 8.0] \times 10^{-1}$ |
| ET/Gen2         | $f_{\text{AGN}}[0.042, 8.8] \times 10^{-1}$ |

---

ering $\sigma_{pp} \approx 5 \times 10^{-26}$ cm$^2$ and $\sigma_{p\gamma} \approx 5 \times 10^{-28}$ cm$^2$ (Particle Data Group et al. 2004, for rough estimations).

4. NEUTRINO BURSTS DETECTION

The expected number of muon neutrinos $\nu_\mu$ from an on-axis GRB event detectable by IceCube and IceCube-Gen2 (Gen2) can be calculated by

$$N(\epsilon_{\nu} > 1 \text{ TeV}) = \int_{1 \text{ TeV}}^{\epsilon_{\nu,\text{max}}} d\epsilon_{\nu} F_{\nu}(\epsilon_{\nu}) A_{\text{eff}}(\epsilon_{\nu}),$$

where $A_{\text{eff}}$ is the effective area of the detector. We obtain the effective areas of IceCube for the up- and down-
Figure 4. Expected all-flavor diffuse neutrino fluence contributed from choked LGRBs in AGN disks as a function of neutrino energy $\epsilon_\nu$. Three local event rates are considered: 100 (solid line), 10 (dashed line), and 1 Gpc$^{-3}$ yr$^{-1}$ (dashed-dot line). The pink circles are observed diffuse neutrino fluence measured by IceCube (Aartsen et al. 2015b).

GRBs, the contribution of choked SGRBs in AGN disks to the neutrino background should be much lower than that of choked LGRBs. We thus only consider the contribution from the latter. Since the cosmic evolution of AGN and star formation rate is not significant (e.g., Madau & Dickinson 2014), we assume that LGRBs in AGN disks are classical LGRBs which closely track the star formation history. We adopt the $f(z)$ distribution based on Yüksel et al. (2008).

We show in Figure 4 the diffuse neutrino fluence by considering three values of the local event rate density due to choked LGRBs in AGN disks are poorly constrained (since these events do not show up as classical GRBs Zhu et al. 2021; Perna et al. 2021). In an extreme case, if $\dot{\rho}_0 = 100$ Gpc$^{-3}$ yr$^{-1}$, i.e., comparable to that of low-luminosity GRBs ($\sim 164^{+98}_{-65}$ Gpc$^{-3}$ yr$^{-1}$, Sun et al. 2015), most of the observed neutrino background fluence could be interpreted by choked LGRBs in AGN disks. If $\dot{\rho}_0 = 10$ Gpc$^{-3}$ yr$^{-1}$, which is $\sim 5\%$ that of low-luminosity GRBs, the neutrinos from such events can contribute up to $\sim 10\%$ of the observed diffuse neutrino background fluence. Future observations of shock breakout transients from AGN disks (Zhu et al. 2021; Perna et al. 2021) will better constrain the cosmological event rate density of these choked GRBs in the AGN channel, leading to a better estimation of their contribution to the neutrino background.

6. DISCUSSION

Choked LGRBs and SGRBs in AGN disks are ideal targets for multi-messenger observations. Besides neutrino emission discussed in this paper, they can also produce electromagnetic signals from optical to $\gamma$-ray bands (Zhu et al. 2021; Perna et al. 2021). For choked LGRBs, associated SNe could be directly discovered by time-domain survey searches (Assef et al. 2018). Within several years of operation by IceCube and Gen2, neutrino bursts from single events would be possible to be directly detected. Choked SGRBs in AGN disks are emitters of electromagnetic, neutrino and GW signals. Future joint observations of electromagnetic, neutrino, and GW signals can reveal the existence these hitherto speculated transient population in AGN disks, shedding light into the interplay between AGN accretion history and the star formation and compact binary merger history in the universe.

ACKNOWLEDGMENTS

We thank Zhuo Li, Di Xiao, Ming-Yang Zhuang, Yuan-Qi Liu, and Di Zhang for valuable comments. The work of J.P.Z is partially supported by the National Science Foundation of China under Grant No. 11721303 and the National Basic Research Program of China under grant No. 2014CB845800. K.W is supported by the National Natural Science Foundation of China under grants 12003007 and the Fundamental Research Funds for the Central Universities (No. 2020kfyXJJ0939). Y.P.Y is supported by National Natural Science Foundation of China grant No. 12003028 and Yunnan University grant No.C176220100087. Y.W.Y is supported by the National Natural Science Foundation of China under Grant No. 11822302 and Yunnan University Grant No.C176220100087. H.G. is supported by the National Natural Science Foundation of China under Grant No. 11722324, 11690024, 11633001, the Strategic Priority Research Program of the Chinese Academy of Sciences, Grant No. XDB23040100 and the Fundamental Research Funds for the Central Universities.

REFERENCES

Aartsen, M. G., Ackermann, M., Adams, J., et al. 2015a, ApJL, 805, L5, doi: 10.1088/2041-8205/805/1/L5

Aartsen, M. G., Abraham, K., Ackermann, M., et al. 2015b, ApJ, 809, 98, doi: 10.1088/0004-637X/809/1/98
Aartsen, M. G., Ackermann, M., Adams, J., et al. 2017, ApJ, 843, 112, doi: 10.3847/1538-4357/aa7569
Abbott, B. P., Abbott, R., Adhikari, R., et al. 2009, Reports on Progress in Physics, 72, 076901, doi: 10.1088/0034-4885/72/7/076901
Abbott, B. P., Abbott, R., Abbott, T. D., et al. 2017a, PhRvL, 119, 161101, doi: 10.1103/PhysRevLett.119.161101
—. 2017b, ApJL, 848, L13, doi: 10.3847/2041-8213/aa920c
—. 2017c, ApJ, 848, L12, doi: 10.3847/2041-8213/aa91c9
—. 2018, Living Reviews in Relativity, 21, 3, doi: 10.1007/s41114-018-0012-9
Achterberg, A., Gallant, Y. A., Kirk, J. G., & Guthmann, A. W. 2001, MNRAS, 328, 393, doi: 10.1046/j.1365-8711.2001.04851.x
Artymowicz, P., Lin, D. N. C., & Wampler, E. J. 1993, ApJ, 409, 592, doi: 10.1086/172690
Assef, R. J., Prieto, J. L., Stern, D., et al. 2018, ApJ, 866, 26, doi: 10.3847/1538-4357/aadd7
Bellovary, J. M., Mac Low, M.-M., McKernan, B., & Ford, K. E. S. 2016, ApJL, 819, L17, doi: 10.3847/2041-8205/819/2/L17
Blandford, R. D., & McKee, C. F. 1976, Physics of Fluids, 19, 1130, doi: 10.1063/1.861619
Bromberg, O., Nakar, E., & Piran, T. 2011a, ApJL, 739, L55, doi: 10.1088/2041-8205/739/2/L55
Bromberg, O., Nakar, E., Piran, T., & Sari, R. 2011b, ApJ, 740, 100, doi: 10.1088/0004-637X/740/2/100
Cheng, K. S., & Wang, J.-M. 1999, ApJ, 521, 502, doi: 10.1086/307572
Chodorowski, M. J., Zdziarski, A. A., & Sikora, M. 1992, ApJ, 400, 181, doi: 10.1086/171984
Collin, S., & Zahn, J.-P. 1999, A&A, 344, 433
Denton, P. B., & Tamborra, I. 2018, ApJ, 855, 37, doi: 10.3847/1538-4357/aaa5a4
Dittmann, A. J., Cantiello, M., & Jermyn, A. S. 2021, arXiv e-prints, arXiv:2102.12484. https://arxiv.org/abs/2102.12484
Eichler, D., Livio, M., Piran, T., & Schramm, D. N. 1989, Nature, 340, 126, doi: 10.1038/340126a0
Fabj, G., Nasim, S. S., Caban, F., et al. 2020, MNRAS, 499, 2608, doi: 10.1093/mnras/staa3044
Fasano, M., Celli, S., Guetta, D., et al. 2021, arXiv e-prints, arXiv:2101.03502. https://arxiv.org/abs/2101.03502
Fong, W., Berger, E., Margutti, R., & Zauderer, B. A. 2015, ApJ, 815, 102, doi: 10.1088/0004-637X/815/2/102
Galama, T. J., Vreeswijk, P. M., van Paradijs, J., et al. 1998, Nature, 395, 670, doi: 10.1038/27150
Kimura, S. S., Murase, K., Bartos, I. 2018, PhRvD, 98, 043020, doi: 10.1103/PhysRevD.98.043020
Kouveliotou, C., Meegan, C. A., Fishman, G. J., et al. 1993, ApJ, 413, L101, doi: 10.1086/186969
Kumar, P., & Zhang, B. 2015, PhR, 561, 1, doi: 10.1016/j.physrep.2014.09.008
Liang, E., Zhang, B., Virgili, F., & Dai, Z. G. 2007, ApJ, 662, 1111, doi: 10.1086/517959
MacFadyen, A. I., & Woosley, S. E. 1999, ApJ, 524, 262, doi: 10.1086/307790
Madau, P., & Dickinson, M. 2014, ARA&A, 52, 415, doi: 10.1146/annurev-astro-081811-125615
MacFadyen, A. I., & Woosley, S. E. 2002, ApJ, 574, 1104, doi: 10.1086/344172
Maggiore, M., Van Den Broeck, C., Bartolo, N., et al. 2020, JCAP, 2020, 050, doi: 10.1088/1475-7516/2020/03/050
Matzner, C. D. 2003, MNRAS, 345, 575, doi: 10.1046/j.1365-8711.2003.06969.x
McKernan, B., Ford, K. E. S., & O’Shaughnessy, R. 2020, MNRAS, 498, 4088, doi: 10.1093/mnras/staa2681
Murase, K. 2007, PhRvD, 76, 123001, doi: 10.1103/PhysRevD.76.123001
—. 2008, PhRvD, 78, 101302, doi: 10.1103/PhysRevD.78.101302
Graham, M. J., Ford, K. E. S., McKernan, B., et al. 2020, PhRvL, 124, 251102, doi: 10.1103/PhysRevLett.124.251102
Gupta, N., & Zhang, B. 2007, Astroparticle Physics, 27, 386, doi: 10.1016/j.astropartphys.2007.01.004
