Choked jet model for the neutrino emission associated with Tidal Disruption Events

JIAN-HE ZHENG,1,2 RUO-YU L I U,1,2 AND XIANG-YU WANG1,2

1 School of Astronomy and Space Science, Nanjing University, Nanjing 210023, People’s Republic of China
2 Key laboratory of Modern Astronomy and Astrophysics (Nanjing University), Ministry of Education, Nanjing 210023, People’s Republic of China

ABSTRACT

Three tidal disruption event (TDE) candidates (AT2019dsg, AT2019fdr, AT2019aalc) have been found to be coincident with high energy astrophysical neutrinos in multi-messenger follow-ups. Recent studies suggest the presence of a quasi-spherical, optically thick envelope around the supermassive black holes in TDEs, resulted from stellar debris after the disruption. We study whether the neutrino signal can be explained by choked relativistic jets inside the envelope. While powerful jets, such as that in Swift J1644+57, can successfully break out the envelope, those with relatively weak power could be choked by the envelope. Choked jets can still accelerate cosmic rays through internal shocks or reverse shocks deep in the envelope, which further produce high-energy neutrinos via interaction with the thermal photons in the envelope. We explore the parameter space of the jets that can produce detectable neutrino flux while being choked. Under reasonable assumption about the envelope mass, we find that the cumulative neutrino numbers of three TDEs are consistent with the expected range imposed by observations. Compared with other proposed models, the relativistic bulk motion of the jets in our model can magnify the neutrino flux by Lorentz boosting. The neutrino time delay relative to the optical peak time of TDEs can be explained as the jet propagation time in the envelope before being choked. The discovery of TDE-associated neutrino events may suggest that jets might have been commonly formed in TDEs, as expected from super-Eddington accretion, but most of them are too weak to break out of the envelopes.

Keywords: Neutrino astronomy (1100) — Relativistic jets(1390) — Tidal disruption(1696) — Cosmic rays(329)

1. INTRODUCTION

The origin of extragalactic high-energy neutrino is one of the main puzzles in neutrino astronomy. Neutrino alert-triggered follow-up searches in electromagnetic data have proven successful in identifying individual blazars as sources, with the most prominent case being TXS 0506+056, which was found to be in a gamma-ray flaring state during the neutrino emission (IceCube Collaboration et al. 2018). In time-integrated point source searches, a clustering of neutrinos from the direction of the starburst galaxy NGC 1068 was found (Aartsen et al. 2020). Recently, optical follow-up observations of neutrino alerts using the Zwicky Transient Facility (ZTF) have identified two optical flares from the centers of galaxies coincident with 100 TeV-scale neutrinos: AT2019dsg associated with the IceCube neutrino event IC191001A (Stein et al. 2021) and AT2019fdr associated with IC200530A (Reusch et al. 2022). The former belongs to the class of spectroscopically-classified tidal disruption events (TDEs) from quiescent black holes, while the latter originates from an unobscured active galactic nucleus (AGN). Afterwards, it was noticed that these TDEs were accompanied by an IR echo, and this connection then led to the identification of a third TDE, AT2019aalc, as the counterpart of IC191119A (van Velzen et al. 2021a). There have been arguments that multiple source populations may contribute to the astrophysical diffuse neutrino flux, based on neutrino event detections and population statistics (Bartos et al. 2021), as well as from spectral shape and directional information (Palladino & Winter 2018).
TDEs are phenomena in which a massive star passes close enough to a supermassive black hole (SMBH) to be ripped apart by its tidal forces. Relativistic jets in TDEs have been proposed as ultra-high energy cosmic ray sources (Farrar & Gruzinov 2009; Farrar & Piran 2014). Neutrino emission is predicted to be produced in both successful jets (Wang et al. 2011; Wang & Liu 2016; Dai & Fang 2017; Lunardini & Winter 2017; Senno et al. 2017) and choked jets (Wang & Liu 2016; Senno et al. 2017). After the identification of AT2019dsg/IC191001A, TDE jets (Winter & Lunardini 2021; Liu et al. 2020), corona and hidden wind (Murase et al. 2020) as well as outflow-cloud interactions (Wu et al. 2022) have been proposed to explain the neutrino emission. While a successful jet has the advantage that it can provide the necessary power for the neutrino emission (see discussion in Winter & Lunardini (2022a)), no convincing direct jet signatures for AT2019dsg have been observed (Mohan et al. 2022). For AT2019fdr, similarly, various plausible cosmic ray acceleration sites have been proposed, such as the corona, a subrelativistic wind, or a relativistic jet (Reusch et al. 2022). The neutrino production site is therefore uncertain, and Winter & Lunardini (2022b) recently discussed unified time-dependent interpretations of these events, considering three models in which quasi-isotropic neutrino emission is due to the interactions of accelerated protons of moderate, medium, and high energy with X-rays, OUV, and IR photons, respectively.

Recent studies suggest the presence of an extended, quasi-spherical, optically thick envelope around the SMBH in TDEs (Loeb & Ulmer 1997; Coughlin & Begelman 2014; Roth et al. 2016; Metzger 2022). The presence of the gas envelope can solve the puzzle that the temperatures (few 10^4 K) found in optically discovered TDEs are significantly lower than the predicted thermal temperature (> 10^5 K) of the accretion disk (Gezari et al. 2012; Wevers et al. 2019). The gas at large radii can absorb UV photons produced by the inner accretion disk and re-emits it at lower temperatures.

As no convincing jet signatures for the three TDEs have been observed, in this paper, we study whether the three neutrino events can be interpreted as arising from choked jets in TDEs. Numerical simulations by De Colle et al. (2012) show that powerful jets can successfully cross the envelope. However, jets without enough power could be choked in the dense envelope (Wang & Liu 2016; De Colle & Lu 2020). Three jetted TDEs, Swift J1644+57, Swift J2058+05, and Swift J1112-8238 all have relativistic jets with a non-thermal X-ray luminosity above 10^{48} ergs^{-1}. Since the black-hole accretion is expected to produce a continuous distribution of jet luminosity, as has been seen in AGN jets, the fact that no less powerful TDE jets have been detected so far is a puzzle. One possibility is that jets with luminosity much less than 10^{48} ergs^{-1} have been choked in the envelopes.

2. MODEL DESCRIPTION

2.1. Envelope formation and the condition for jets being choked

As a result of the disruption process, about half of the stellar mass remains bound. After the first passage of the star, the debris forms an elongated stream of gas that later spreads out, intersects itself, and dissipates its orbital kinetic energy into heat (e.g. Rees 1988; Evans & Kochanek 1989). Near the pericenter, shocks may organize the high-entropy debris into a thick rotating torus, which would be dominated by radiation pressure. Farther away from the pericenter, rotation is dynamically less important and the strong radiation pressure could disperse the marginally bound gas into a quasi-spherical configuration (Loeb & Ulmer 1997). The outer geometry of the envelope would therefore be close to spherical, whereas its base near the torus would deviate from sphericity due to rotation. The hydrodynamic study by Coughlin & Begelman (2014) also shows that the infalling gas traps accretion energy until it inflates into a weakly bound, quasi-spherical structure with gas extending to the poles. A quasi-spherical emission surface is also supported by recent spectropolarimetry observations of some TDEs (e.g., Patra et al. 2022).

Another approach to forming an envelope is through the super-Eddington wind outflow. Gas inflowing at super-Eddington rates cannot radiatively cool on the dynamical or accretion time-scale, in which case photons are trapped and advected with the fluid. Low initial binding energy coupled with this globally adiabatic evolution suggests that only a fraction of the infalling debris is likely to remain gravitationally bound to the SMBH, with the remainder unbound in an outflow (Strubbe & Quataert 2009; Metzger & Stone 2016; Vinkó et al. 2015; Miller 2015).

We assume the envelope is spherically symmetric and quasi-static. We model the envelope with a power-law density profile, \( \rho = \rho_{\text{out}} (r/R_{\text{out}})^{-n} \), extending from an inner radius, \( R_{\text{in}} \), to an outer radius, \( R_{\text{out}} \). An exponent \( n = 3 \) applies to a radiation pressure-supported envelope (Loeb & Ulmer 1997), while \( n = 2 \) applies to a steady state wind outflow (Roth et al. 2016). The hydrodynamic solutions of Coughlin & Begelman (2014)
find values of \( n = 1.5 \) to 3. The density at the outer surface \( \rho_{\text{out}} \) of the envelope is given by

\[
\rho_{\text{out}} = \frac{\xi M_*}{4\pi R_{\text{out}}^3} \left\{ \begin{array}{ll}
\frac{3 - n}{1 - (R_{\text{in}}/R_{\text{out}})^{3-n}} & n \neq 3 \\
\frac{1}{\ln (R_{\text{out}}/R_{\text{in}})} & n = 3
\end{array} \right.,
\]

where \( M_* \) is the mass of the disrupted star, \( \xi = M_{\text{env}}/M_* \) is the mass ratio between the envelope and the disrupted star, \( R_{\text{out}} \) and \( R_{\text{in}} \) are the outer radius and the inner radius, respectively. Generally, the inner radius is \( R_{\text{in}} = \max \{ R_t, R_{\text{ISCO}} \} \), where \( R_t = R_c(M_{\text{BH}}/M_*)^{1/3} \) is the tidal disrupted radius of TDEs and \( R_{\text{ISCO}} \) is the innermost stable circular orbit (ISCO) of SMBHs. If the mass of SMBH and disrupted star are both large, the tidal disrupted radius is greater than the radius of ISCO.

The dynamics of the jet propagating in the TDE envelope has been discussed in (Wang & Liu 2016; De Colle & Lu 2020). The envelope material changes the dynamics of the jet launched by the SMBH in a way similar to the jet propagation in the collapsing star in a long gamma-ray burst (Begelman & Cioffi 1989; Bromberg et al. 2011). As the jet advances in the surrounding envelope, the jet drives a bow shock ahead of it. The jet is capped by a termination shock, and a reverse shock propagates back into the jet, where the jet is decelerated and heated. The velocity of the jet head can be related to the isotropically equivalent jet kinetic luminosity \( L_{\text{kin,iso}} \) and the density of the envelope (Bromberg et al. 2011), \( \bar{L} = L_{\text{kin,iso}}/(4\pi r^2 \rho c^3) \). In the limit \( \bar{L} \ll 1 \), the dimensionless velocity of the jet head is \( v_h \approx \bar{L}^{1/2} \), i.e.,

\[
v_h = \beta_c c = \left( \frac{L_{\text{kin,iso}}}{4\pi r^2 \rho c^3} \right)^{1/2} \left( \frac{M_*}{M_{\odot}} \right)^{1/2} L_{\text{kin,iso,46}}^{1/2} \tag{2}
\]

in the observer’s frame (Begelman & Cioffi 1989; Wang & Liu 2016).

In the free-fall situation, matter returns to the region near the pericenter radius at a rate \( \dot{M} \propto (t/\tau)^{-5/3} \), where \( \tau \) is the characteristic timescale for initiation of this power-law accretion rate, which is roughly the orbital period of the most bound debris, i.e.,

\[
\tau = \frac{\pi}{M_*} \sqrt{\frac{\eta M_{\text{BH}} R^3}{2G}} = 41 M_{\text{BH,6}}^{1/2} \left( \frac{M_*}{M_\odot} \right)^{-0.1} \text{ day,} \tag{3}
\]

where \( \eta \) is a correction factor depending on the structure of stars, which is close to \( \eta \approx 1 \) for solar-mass stars (Phinney 1989; Gezari 2021). We adopt the mass-radius relation of stars \( R \propto M_*^{0.6} \) from Kippenhahn et al. (2012). As the jet power may scale with the accretion rate as \( L_j \propto \dot{M} \) in the super-Eddington accretion phase (Krolik & Piran 2012; Piran et al. 2015), the characteristic lifetime of the jet is also \( \tau \) and \( L_j \propto (t/\tau)^{-5/3} \) after that. Noting that the observed power indices \( \alpha \) in TDEs are diverse (van Velzen et al. 2021b), we take the power index as a free parameter assuming the jet luminosity follows \( L_j = L_{j,0}(1 + t/\tau)^{-\alpha} \), where \( L_{j,0} \) is the initial jet luminosity. The isotropic luminosity is related to the beam-corrected jet luminosity by \( L_{j,\text{iso}} = 2L_j/\theta_j^2 \), where \( \theta_j \) is the half opening-angle of the jet. The value of the half opening-angle of the jet is assumed to be \( \theta_j \approx 1/\Gamma \), where \( \Gamma \) is the bulk Lorentz factor of the jet. The typical value of the Lorentz factors is expected to be \( \Gamma \approx 10 \), as inferred from jetted TDEs (Metzger et al. 2016). Since a significant part of the jet energy may be transferred into reverse shocks and internal shocks and only a portion of the energy remains in the advancing jet, we assume \( L_{\text{kin,iso}} = 0.5 L_{j,\text{iso}} \).

The breakout time is

\[
t_{\text{br}} = \tau \left\{ \frac{[\alpha + 2] - (R_{\text{in}}/R_{\text{out}})^{2-n/2}}{2C_0 \tau} + 1 \right\}^{-1} \tag{4}
\]

where \( C_0 = \sqrt{L_{\text{kin,iso,0}}/4\pi c \rho_{\text{out}} R_{\text{out}}^3} \). For a constant-luminosity jet with \( \alpha = 0 \) and \( n = 3 \), the breakout time reduces to \( t_{\text{br}} \approx 2/C_0 \propto R_{\text{out}}^{3/2} L_{\text{kin,iso}}^{1/2} M_*^{-1/2} \). If the breakout time is longer than the lifetime of the jet, the jet will be choked by the envelope. By comparing the jet break-out time with the lifetime \( \tau \) of the jet, we can obtain the condition for jets being choked. The dashed lines in Figure 1 show the critical values of the jet luminosity and the bulk Lorentz factor for choked jets in three TDEs.

### 2.2. Neutrino production

Internal shocks may occur due to the internal collisions within the jets, resulted from the inhomogeneity in the jet velocity. Temporal variability with \( \delta t \approx 100 \) s has been seen in the X-ray emission of the jetted TDE Swift J1644+57, which is thought to be generated by internal shocks (Burrows et al. 2011). The initial Lorentz factor of the jet is \( \Gamma \approx 3 - 10 \), and hence the collision occurs at \( R_{\text{int}} \approx \Gamma^2 c \delta t = 3 \times 10^{14} \text{cm} \Gamma^2 \). We assume the jet is matter-dominant when internal shock occurs. The density of jets \( n_j \approx 1.8 \times 10^7 \text{cm}^{-3} L_{\text{iso,46}} R_{\text{int,14}}^{-2} \Gamma^{-2} \) is much lower than that of envelopes \( n \approx 10^{13} \text{cm}^{-3} M_{\odot}^{-1} R_{\text{int,14}}^{-3} \). Thus, the low-density jet forms a cavity inside the envelope, while the jets are propagating. The envelope contains dense thermal photons, which diffuse out of the
Figure 1. The accumulative neutrino numbers (the color map) as a function of the jet luminosity and bulk Lorentz factor assuming envelope masses of three TDEs given in case A. The black solid lines represent the parameter values corresponding to the minimally required neutrino numbers, which are $N_\nu = 0.008$ for AT2019dsg and $N_\nu = 0.007$ for AT2019fdrl and AT2019aacl (Stein et al. 2021; Reusch et al. 2022). The black dash line represents the critical parameter values for jets being choked and the direction of the arrow means that the permitted parameter space for choked jets is below this line. The top and bottom panels represent the cases of $n = 3$ and $n = 2$ for the envelope density profile, respectively. Internal shock radius are taken to be $R_{\text{out}} = 3 \times 10^{14}$ cm for AT2019aacl and AT2019fdrl, and $R_{\text{out}} = 10^{14}$ cm for AT2019dsg. $\alpha = 5/3$, $\beta = 5/12$, $\Gamma = 3$ and $f_{\text{ph}} = 0.05$ are used in the calculation. Other parameters are listed in Table 1.

optically thick envelope and enter the optically-thin cavity. The radiation diffusion time in the envelope is

$$t_{\text{diff}} = \frac{r^2 \kappa \rho(r)}{c} \text{ at a particular radius } r (Roth et al. 2016).$$

Using $\kappa R_{\text{out}} \rho_{\text{out}} = 1$, the diffusion time is $t_{\text{diff}} = \frac{R_{\text{out}}}{c}$ in the case of $n = 2$ and $t_{\text{diff}} = (\frac{R_{\text{out}}}{c})(\frac{R_{\text{out}}}{r})$ in the case of $n = 3$. Compared to the jet propagation time ($r/v_h$), the number density of photons in the cavity should be suppressed by a factor of $f_{\text{ph}} = \min\{\frac{(\epsilon_{\gamma})}{(\epsilon_{\text{out}})}, 1\}$ in the case of $n = 2$ and $f_{\text{ph}} = \min\{\frac{(\epsilon_{\gamma})}{(R_{\text{out}})^2}, 1\}$ in the case of $n = 3$. Internal shocks that propagate into the low-density jets are collisionless, although they locate inside the optically thick envelope (Wang & Liu 2016). It has been shown that shocks in TDE jets can accelerate cosmic rays to ultrahigh energies (Farrar & Gruzinov 2009; Wang et al. 2011).

The accelerated protons will produce neutrinos by $pp$ and $p\gamma$ reactions. The dense envelope and high temperature provide a thick target for neutrino production. The cooling time of $pp$ interaction is

$$t_{pp}^{-1} = \kappa_{pp} n_p \sigma_{pp} c,$$

where $\kappa_{pp} \approx 0.5$ is the inelasticity of protons and $\sigma_{pp} \approx 3 \times 10^{26}$ cm$^2$ is the cross section. Meanwhile, the cooling time of $p\gamma$ interaction is

$$t_{p\gamma}^{-1} = \frac{c}{2 \gamma_p} \int_{\sigma_{\text{th}}}^{\infty} \sigma_{p\gamma}(\gamma) \kappa_{p\gamma}(\gamma) \frac{\sigma_{\text{th}}}{\sigma_{p\gamma}(\gamma)} \int_{\epsilon/2\gamma_p}^{\infty} \epsilon^{-2} \frac{dn}{d\epsilon} d\epsilon,$$

where $\gamma_p$ is Lorentz factors of protons, $dn/d\epsilon$ is the number density of seed photons, $\epsilon_{\text{th}} \approx 145$ MeV is the threshold energy of $p\gamma$ interaction, $\epsilon$ and $\gamma$ are photons’ energy in the observer’s frame and center of mass frame of protons, respectively. Seed photons of $p\gamma$ interaction are mainly dominated by thermal photons of the envelope.

To calculate the neutrino flux, we need to know the temperature of the thermal photons in the envelope. The photosphere temperature can be obtained from observations. To obtain the interior temperature, we need to solve the energy transferring equation, which is

$$\frac{du}{dr} = -\frac{3 \kappa \rho(r) L_{\text{bol}}}{4 \pi c r^2},$$

where $L_{\text{bol}}$ is the bolometric luminosity.
where $L_{\text{bol}}$ is the bolometric luminosity, $u$ is the energy density and $\kappa_{\text{es}} = 0.33 \text{cm}^2/\text{g}$ is free electron scattering opacity. Adopting temperature of ideal photon gas $T = (uc/4\sigma)^{1/4}$ where $\sigma$ is Stefan-Boltzmann constant, the solution of temperature inside the envelope is (Roth et al. 2016)

$$T(r) = T_{\text{ph}} \left[ \frac{3\tau_{\text{out}}}{4(n + 1)} \left( \frac{r_{\text{out}}^{n+1}}{r^{n+1}} - 1 + \frac{4(n + 1)}{3\tau_{\text{out}}} \right) \right]^{1/4},$$

where $\tau_{\text{out}} \equiv \kappa_{\text{es}}\rho_{\text{out}}R_{\text{out}}$. The temperature follows $T \propto r^{-(n+1)/4}$ when $r \ll R_{\text{out}}$. The large number density of thermal photons leads to $t_{p\gamma}^{-1} \gg t_{p\gamma}^{-1}$. Therefore, the main production channel of neutrinos is $p\gamma$ interaction.

The typical energy of protons interacting with thermal photons is $\epsilon_p = 0.15 \text{GeV}^2/3kT = 0.6 - 6 \text{PeV}$ for $T = 10^5 - 10^6 \text{K}$ and the neutrino energy is correspondingly $\epsilon_{\nu} = 0.05\epsilon_p = 30 - 300 \text{TeV}$. We use a power-law proton spectrum $dN/dE \propto E^{-p}$ with $p = 2$. The maximum energy of protons is obtained by equating the acceleration time $t_{\text{acc}} = \gamma_p m_pc/\eta_{\text{acc}}eB$ and the cooling time $t_{p\gamma}$, where $\eta_{\text{acc}} = 0.1$ is the acceleration efficiency, and the magnetic field is given by $B = \sqrt{2\epsilon_B L_{\text{iso}}/R_{\text{int}}^2}$. The maximum energy is determined by the balance between the cooling time of $p\gamma$ interactions and the acceleration. The fluence of neutrinos is

$$\epsilon_{\nu} F_{\nu} \approx \int_{t_{\text{tri}}/(1+z)}^{t_{\text{acc}}/(1+z)} \frac{L_p f_{p\gamma} f_{\mu,\text{sup}}}{32\pi D_E^2 \ln(\epsilon_{\nu,\text{max}}/\epsilon_{\nu,\text{min}})} e^{-\epsilon_{\nu}/\epsilon_{\nu,\text{min}}} \, dt,$$

where $t_{\text{tri}}$ is the neutrino trigger time since the discovery of TDEs in the observer’s frame, $L_p = \epsilon_p L_{\text{j,iso}}$ is the isotropic proton luminosity and $f_{\mu,\text{sup}} = 1 - \exp(-t_{\mu,\text{dec}}/t_{\mu,\text{syn}})$ is the suppression factor for muon decay. $t_{\mu,\text{dec}} = \gamma_{\mu} t_{\mu,\text{dec}}$ is the decay time of relativistic muons where $\tau_{\mu,\text{dec}}$ is the mean lifetime of muons in the rest frame and $t_{\mu,\text{syn}}$ is the synchrotron cooling time of relativistic muons. The minimum energy of protons is $\epsilon_{p,\text{min}} = \Gamma m_pc^2$. We take the equipartition factors $\epsilon_B = 0.1$ and $\epsilon_p = 0.2$ in the calculation.

Then we calculate the cumulative muon neutrino number between 10 TeV and 1 PeV.

$$N_{\nu} = \int_{10 \text{ TeV}}^{1 \text{ PeV}} A_{\text{eff}}(\epsilon_{\nu}) F_{\nu} d\epsilon_{\nu}.$$  

The effective area of detectors is $A_{\nu}$ effective area taken from Blaufuss et al. (2019).

3. APPLICATION TO AT2019FDR, AT2019DSG AND AT2019AALC

To calculate the break-out time of jets and neutrino flux in the three TDEs, we need first to know the enve-
lope mass of each TDE. We suggest two approaches to estimate the envelope mass, one uses the photosphere radius given by $R_{\text{ph}} \simeq 10^{15}\text{cm} (M_{\text{env}}/0.5M_{\odot})^{1/2}$ for a radiation-pressure supported envelope or a steady state wind outflow (Loeb & Ulmer 1997; Roth et al. 2016), and the other assumes the envelope possesses a power-law radial density profile with a characteristic radius and a sharp outer edge, where the photosphere radius scales as $R_{\text{ph}} \propto M_{\text{env}}M_{\text{BH}}^{2/3}$ (see Eq. 7 in Metzger (2022); $M_{\text{BH}}$ is the black hole mass). We define the above two cases as Case A and Case B, respectively.

AT2019dsg has a photosphere radius of $R_{\text{ph}} = 5.3 \times 10^{14}\text{cm}$, which is common in TDEs, while AT2019fdr and AT2019aalc are rare events with photosphere radii an order of magnitude larger. In case A, the inferred envelope masses are $0.1M_{\odot}$ for AT2019dsg, $21.7M_{\odot}$ for AT2019fdr and $9.3M_{\odot}$ for AT2019aalc. The decrease of the event rate with the photosphere radius has been interpreted as being due to the lower number of high-mass stars, as implied by the initial mass functions of stars (van Velzen et al. 2021b). Here we have assumed that the opacity is the same for three TDEs. In reality, the opacity of AT2019fdr and AT2019aalc could be higher since their photosphere temperatures are lower (Roth et al. 2016), and then their envelope masses will be correspondingly lower than the above estimates.

In case B, the new scaling relation suggests envelope masses of about $1M_{\odot}$, $5M_{\odot}$ and $4M_{\odot}$ for AT2019dsg, fAT2019fdr and AT2019aalc, respectively. Large masses for the disrupted stars in AT2019fdr and AT2019aalc are natural considering the overall larger energy budget and longer duration of the flares in these two TDEs, since the available energy may be proportional to a fixed fraction of the star’s mass.

The black hole mass $M_{\text{BH}}$ is estimated from the optical spectrum of the host galaxy or the flare (see Table 1 of van Velzen et al. (2021a)). We summarize the masses of the disrupted stars and the black holes, as well as other relevant parameters, for three TDEs in Table 1. The jet lifetime $\tau$ can then be obtained using Eq. 3, which is also listed in Table 1. We assume that the jet luminosity and envelope temperature decrease with time following the relations $L_{\text{j}} = L_{\text{j,0}}(1 + t/\tau)^{-\alpha}$ and $T_{\text{ph}} = T_{\text{ph,0}}(1 + t/\tau)^{-\beta}$. The fiducial values are $\alpha = 5/3$ and $\beta = 5/12$, where $\alpha = 5/3$ corresponds to the freefall accretion and $\beta = 5/12$ is taken from the simulation (Krolik et al. 2020).

The cumulative neutrino numbers of three TDEs are shown by the color intensity in Figures 1 and 2, respectively, for case A and case B. The solid lines show the parameter values corresponding to the minimum neutrino number of each TDE required by observations at 90% confidence (Stein et al. 2021). The space between the solid line and dashed line (along the direction of the arrow) represents the allowed range for the jet luminosity and the Lorentz factor of the jet, where neutrinos above the minimum number will be produced while the jets are choked before break-out.

AT2019fdr is an exceptionally luminous TDE candidate with a peak optical/UV bolometric luminosity of $10^{45}\text{ergs}^{-1}$ and a long duration ($\geq 1000\text{days}$). It can be seen from Figure 1 that the cumulative neutrino number in the choked jet model can reach $N_{\nu} = 0.02$ for a large envelope mass (Case A). With a lower envelope mass (Case B), as shown in Figure 2, the cumulative neutrino number is at most $N_{\nu} = 0.005$, which falls slightly below the minimum required value. Given that the minimum required value of neutrinos corresponds to a 90% confidence, the choked jet model is still possible for AT2019fdr. We also find that the maximum cumulative neutrino number in the $n = 2$ case is slightly larger than that in the $n = 3$ case, which is because the break-out time is longer for the $n = 2$ profile according to Eq. 4.

For AT2019dsg, its small photosphere radius implies a very small mass of the envelope ($0.1M_{\odot}$) in Case A. The jet can break out easily in this case and only very weak jets will be choked. This leads to a cumulative neutrino number of $N_{\nu} = 0.001$ in the optimistic case,

| Table 1. Properties and parameters |
|-----------------------------------|
|                                   | AT2019fdr | AT2019dsg | AT2019aalc |
|---|---|---|---|
| redshift | 0.267 | 0.051 | 0.0356 |
| $\epsilon_{\nu}$ (TeV) | 82 | 217 | 176 |
| $t_{\text{ev}}$ (days) | 400 | 180 | 300 |
| $L_{\text{peak}}$ ($10^{48}\text{ergs/s}$) | $14^a$ | $1.8^b$ | $2.7^c$ |
| $T_{\text{ph}}$ (K) | $13526^a$ | $39000^b$ | $11000^c$ |
| $R_{\text{ph}}$ ($10^{14}\text{cm}$) | 78$^a$ | 5.3 | 51 |
| $M_{\text{BH}}$ ($10^6M_{\odot}$) | 35$^a$ | 5$d$ | 20$d$ |
| $\tau$ (days) | 180 | 40$^{90b, c}$ | 140 |
| $M_{\text{env}}$ ($M_{\odot}$) in case A | 21.7 | 0.1 | 9.3 |
| $M_{\text{env}}$ ($M_{\odot}$) in case B | 5 | 1 | 4 |

$^a$Reusch et al. (2022)
$^b$van Velzen et al. (2021b)
$^c$Winter & Lunardini (2022b)
$^d$van Velzen et al. (2021a)
$^e$40-days is obtained by fitting observations of bolometric light curves (van Velzen et al. 2021b), while 90-days is calculated from Eq. 3. We use $\tau=40$ days in our calculations.
Figure 3. Upper panel: Neutrino spectra of three TDE events in Case A. The jet luminosities are taken to be $L_{j,0} = 4 \times 10^{43}$ erg/s for AT2019dsg, $L_{j,0} = 5 \times 10^{45}$ erg/s for AT2019fdr and $L_{j,0} = 10^{45}$ erg/s for AT2019aalc. Lower panel: Neutrino spectra of three TDE events in Case B. The jet luminosities are taken to be $L_{j,0} = 3 \times 10^{44}$ erg/s for AT2019dsg, $L_{j,0} = 10^{45}$ erg/s for AT2019fdr and $L_{j,0} = 5 \times 10^{44}$ erg/s for AT2019aalc. The solid and dot-dashed lines represent the neutrino spectra in the $n = 3$ and $n = 2$ profiles, respectively.

which falls below the expected range $(0.008 \lesssim N_\nu \lesssim 0.76)$(Stein et al. 2021). Thus, the choked jet model is disfavored for AT2019dsg if the envelope mass is indeed as small as $0.1 M_\odot$. In Case B, with an envelope mass of $1 M_\odot$, the optimal neutrino number of AT2019dsg can reach $N_\nu = 0.015$, satisfying the requirement of observations (Stein et al. 2021).

Similar to AT2019fdr, AT2019aalc is also a long-lasting TDE ($\geq 700$ days) and has a large photosphere radius. The neutrino number of AT2019aalc can reach $N_\nu = 0.5$ in Case A and $N_\nu = 0.3$ in Case B.

Figure 4. Time evolution of cumulative neutrino numbers. The symbols and parameter values are identical with those in Fig. 3.

The neutrino spectra for three TDEs are shown in Figure 3 for both Case A (upper panel) and Case B (bottom panel). The high temperature of $\sim 10^5$ K inside the envelope leads to a neutrino energy of the order of hundreds of TeV, which is consistent with the observed values of three neutrinos. Because of the high $p\gamma$ efficiency with $f_{p\gamma} \simeq 1$ for neutrino production, the neutrino spectrum follows the proton spectrum and is thus flat above a critical energy.

The evolution of the cumulative neutrino numbers with time is shown in Figure 4. The cumulative neutrino number increases with time up to a point corresponding to about $\sim 2 \tau$ before flattening. This is because although the jet luminosity starts to decrease after $\tau$, the time-integrated neutrino flux still increases. This can naturally explain the time delay between the neutrino arriving time and the optical peak time of the TDEs.
4. CONCLUSIONS AND DISCUSSIONS

We have proposed that the neutrino emission associated with three non-jetted TDEs can be explained by the choked jet model, where relativistic jets are choked inside the quasi-spherical, optically thick envelope formed from stellar debris of TDEs. From the point of view of observations, the presence of an envelope can solve the puzzle that the temperatures (few 10^4 K) found in optically discovered TDEs are significantly lower than the predicted thermal temperature (> 10^5 K) of the accretion disk (Gezari et al. 2012; Arcavi et al. 2014). While powerful jets, such as that in Swift J1644+57, can break out the envelope, less powerful jets would be choked inside the envelope. Cosmic-ray protons accelerated by shocks in choked jets produce high-energy neutrinos via $p\gamma$ interactions with the thermal photons in the envelope. The energy of neutrinos produced by choked jets is typically around hundreds of TeV, consistent with measured energy of three neutrinos. The hundreds of days time-delay between neutrino arrival time and TDE optical peaks can be explained as the propagation time of the jets before being choked in the envelope. The cumulative neutrino numbers in our model are consistent with the expected range for individual TDEs and can reach $N_\nu = 0.5$ in the case of AT2019aalc. One important advantage of the choked jet model is that the neutrino flux is magnified by the beaming effect due to the relativistic bulk motion.

The non-detection or weak radio emission of the three TDEs disfavor the presence of successful relativistic jets (van Velzen et al. 2021a; Reusch et al. 2022), but are consistent with the choked jet model. For the choked TDE jets, as the neutrino production site is within the optically thick region, the associated high-energy gamma-rays cannot escape. Instead, high-energy gamma-rays are absorbed by low-energy electrons and photons in the envelope, depositing their energy finally into the envelope. Therefore, these choked TDE jets are hidden sources of gamma-rays, consistent with the non-detection of GeV emission from these TDEs (van Velzen et al. 2021a). The late-time appearance of thermal X-ray emission in AT2019fdr and AT2019aalc after the neutrino trigger supports that the accretion discs are obscured by some optically thick materials at early time. The X-rays can only leak out after the obscuring gas has become transparent to X-rays (Metzger & Stone 2016; Lu & Kumar 2018). For AT2019dsg, the detection of thermal X-ray emission in the first ~ 40 days and the rapid decrease of the X-ray flux can be explained by an increasing X-ray obscuration (Stein et al. 2021), which could be due to the gradual formation of the envelope. The non-detection of any jet emission at early time in AT2019dsg may be due to that the jet formation is also delayed (Cendes et al. 2022). In addition, the neutrino emission from AT2019dsg could be due to off-axis jets or other processes (Winter & Lunardini 2021; Liu et al. 2020; Murase et al. 2020; Wu et al. 2022).

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