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Probing the tidal disruption flares of massive black holes with high-energy neutrinos

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The recently discovered high-energy transient Swift $J164449.3 + 573451$ (Sw $J1644 + 57$) is thought to arise from the tidal disruption of a passing star by a dormant massive black hole. Modeling of the broadband emission suggests the presence of a powerful relativistic jet, which contributes dominantly to the observed x-ray emission. Here we suggest that protons can be accelerated to ultrahigh energies by internal shocks occurring in the jets, but their flux is insufficient to account for the observed flux of ultrahigh-energy cosmic rays. High-energy protons can produce $\sim$0.1–10 PeV neutrinos through photomeson interactions with x-ray photons. The large x-ray fluence (7×10^{-4} erg cm^{-2}) and high photopion efficiency, together with the insignificant cooling of secondary mesons, result in bright neutrino emission expected from Sw $J1644 + 57$ if the jet composition is matter-dominated. One to several neutrinos may be detected by a Km^3-scale detector from one tidal disruption event similar to Sw $J1644 + 57$, thereby providing a powerful probe of the composition of the jets.

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Massive black holes are believed to reside at the centers of most galaxies and the vast majority of them are considered to be dormant. It was long predicted that, if a star of mass $M_\star$ and radius $R_\star$ passes occasionally within the disruption radius $r_T \simeq R_\star(M_{\rm BH}/M_\star)^{1/3}$ of the dormant black hole (where $M_{\rm BH}$ is the black hole mass), the star will be torn apart by gravitational tidal forces, leading to a transient accretion disk and a bright panchromatic flare [1]. Candidates of such tidal disruption flares (TDFs) have been suggested [2]. Relativistic jets are expected to form in such accretion disk systems and may produce observational phenomena as well [3].

The high-energy transient Swift $J164449.3 + 573451$ (hereafter Sw $J1644 + 57$) was discovered by the Swift satellite on 28 March 2011 at 12:57:45 UT [4]. The fact that it occurred near the center of a galaxy, no archival x-ray emission before detection, and its analogy with blazars in the x-ray emission suggest that it is most likely to be a TDF event [5–7]. At redshift $z = 0.3534$ (corresponding to a luminosity distance of $d_L = 1.8$ Gpc) [6], the early x-ray flare emission reached a peak luminosity of $L_{\rm X} \simeq 4.3 \times 10^{48}$ erg s^{-1} (isotropic equivalent) and then transited to a low state with a median luminosity of $L_{\rm X} \simeq 2.96 \times 10^{47}$ erg s^{-1} in 0.4–13.5 keV over a time $\Delta T \sim 10^6$ s [4]. Correcting for the live-time fraction of the observation, the total unabsorbed fluence is $S_{\rm X} = 7.1 \times 10^{-4}$ erg cm^{-2} in the observed 0.3–10.0 keV band [4]. This gives an estimate of the total isotropic equivalent energy of $E_{\rm X} = 3 \times 10^{53}$ erg in the 0.4–13.5 keV rest frame energy band. The observed minimum x-ray variability time of Sw $J1644 + 57$, $t_v \simeq 100$ s [4], constrains the size of the black hole under the assumption that the size of the central engine determines the shortest variability and suggests an upper limit on the massive black hole mass $M_{\rm BH} \leq 8 \times 10^6 M_\odot$. The observed peak luminosity is super-Eddington and requires a strongly anisotropic radiation pattern with a relativistic jet of a bulk Lorentz factor of $\Gamma \simeq 10$ pointed towards us [4,5]. Modeling of the broadband spectral energy distribution (SED) also requires a powerful relativistic jet with $\Gamma \simeq 10$, which produces dominantly the observed x-ray emission [4,5]. A relativistic jet is also required to explain the radio transient [8]. However, how the jet is launched is not well understood. It is believed that the composition of the jet, whether it is matter-dominated or magnetic-field–dominated, is crucial to unveiling the formation mechanism of the jets. In this paper, we suggest that jets in TDFs can produce bright emission in high-energy neutrinos if the jet is matter-dominated, and thus neutrino observations provide an important tool to diagnose the jet composition.

Internal shocks and particle acceleration.—We consider that a TDF event produces a relativistic matter-dominated wind of luminosity $L_w \sim 10^{49}$ erg s^{-1}, moving with a bulk Lorentz factor $\Gamma \sim 10$. Variability of the source on time scale t_v, resulting in fluctuations in the wind bulk Lorentz factor Γ on similar time scale, would lead to semirelativistic internal shocks [9] in the ejecta at a radius

$$R \simeq 2\Gamma^2 c t_v = 6 \times 10^{14} \Gamma_1^2 t_{v,2} \text{ cm}, \tag{1}$$

which is well above the photosphere radius at $R_{ph} = \sigma_{\rm T} L_w/(4\pi\Gamma^3 m_p c^3) = 1.2 \times 10^{13} L_{w,49}\Gamma_1^{-3}$ cm, where $\sigma_{\rm T}$ is the Thomson cross section. We use cgs units and the denotation $Q = 10^x Q_x$ throughout the paper.

Denoting ϵ_B as the fraction of the wind kinetic energy converted into magnetic fields, we have a magnetic field $B' = (8\pi\epsilon_B L_w/4\pi R^2\Gamma^2 c)^{1/2} = 1.3 \times 10^3 \epsilon_{B,-1}^{1/2} L_{w,49}^{1/2} R_{14.8}^{-1} \Gamma_1^{-1}$ G, where the prime symbol represents quantities measured in the comoving frame of the shock. It is assumed that internal shocks accelerate protons with a spectrum $dn/d\varepsilon_p \sim \varepsilon_p^{-2}$, where ε_p is the proton energy in the observer frame. The maximum proton energy is set by comparing the acceleration time $t'_{\rm acc} = \alpha\varepsilon_p/(e\Gamma B'c) = 860\alpha(\frac{\varepsilon_p}{10^{20}\ {\rm eV}})\epsilon_{B,-1}^{-1/2}L_{w,49}^{-1/2}R_{14.8}$ s with the shock dynamic time $t'_{\rm dyn} = R/\Gamma c = 10^3\Gamma_1 t_{v,2}$ s, where $\alpha\sim$ a few, describing the ratio between the acceleration time and Larmor time. This gives a maximum proton energy of $\varepsilon_{\rm max,dyn} = 2.4\times 10^{20}\alpha^{-1}\epsilon_{B,-1}^{1/2}L_{w,49}^{1/2}\Gamma_1^{-1}$ eV. The maximum energy is also limited by the cooling time of protons. The synchrotron cooling time is $t'_{\rm syn} = 6\pi m_p^4 c^3\Gamma/(\sigma_T m_e^2\varepsilon_p B^2) = 240\epsilon_{B,-1}^{-1}L_{w,49}^{-1}R_{14.8}^2\Gamma_1^3(\frac{\varepsilon_p}{10^{20}\ {\rm eV}})^{-1}$ s, which gives a maximum proton energy $\varepsilon_{\rm max,syn} = 0.5\times 10^{20}\alpha^{-1/2}\epsilon_{B,-1}^{-1/4}L_{w,49}^{-1/4}t_{v,2}^{1/2}\Gamma_1^{5/2}$ eV. Another process that may prohibit the acceleration of protons to ultrahigh energies (UHE) is the photopion cooling loss. UHE protons of energy ε_p interact with soft photons with energy $\epsilon_\gamma = 0.15\ {\rm GeV}^2\Gamma^2/\varepsilon_p = 0.15\Gamma_1^2\varepsilon_{p,20}^{-1}$ eV, which locate at near infrared (NIR) band for typical jet parameters. The number density of NIR photons in the comoving frame is $n'_{\rm NIR} = L_{\rm NIR}/(4\pi R^2\Gamma c\epsilon_\gamma) = 2.5\times 10^{14}L_{\rm NIR,44}\Gamma_1^{-5}t_{v,2}^{-2}(\epsilon_\gamma/0.15\ {\rm eV})^{-1}\ {\rm cm}^{-3}$, where $L_{\rm NIR}\simeq 10^{44}\ {\rm erg\,s^{-1}}$ is the NIR luminosity at times a few days after the initial trigger [4]. The photopion cooling time in the comoving frame is $t'_{p\gamma} = 1/(\sigma_{p\gamma}n'_{\rm NIR}cK_{p\gamma}) = 1200L_{\rm NIR,44}^{-1}\Gamma_1^5 t_{v,2}^2(\epsilon_\gamma/0.15\ {\rm eV})$ s, where $K_{p\gamma}$ is the inelasticity and $\sigma_{p\gamma} = 5\times 10^{-28}\ {\rm cm}^{-2}$ is the peak cross section at the Δ resonance. By equating $t'_{\rm acc}$ with $t'_{p\gamma}$, we get the maximum proton energy limited by the photopion cooling process, $\varepsilon_{{\rm max},p\gamma} = 1.3\times 10^{20}\alpha^{-1}\epsilon_{B,-1}^{1/4}L_{w,49}^{1/4}L_{\rm NIR,44}^{-1/2}\Gamma_1^{5/2}t_{v,2}^{1/2}$ eV. Thus, internal shocks in TDFs can accelerate protons to energies above 10^{19} eV, in support of the earlier suggestion that TDFs can produce ultrahigh-energy cosmic rays (UHECRs) [10]. However, the flux of such UHE protons contributed by TDFs in the universe is insufficient to explain the observed flux of UHECRs, as we show below.

UHECR flux.—The Swift Burst Alert Telescope, with a field of view of $S_{\rm FOV}\simeq 4\pi/7$ sr, has detected one such event in a time $T\simeq 7$ years at a flux that would have been detectable up to a luminosity distance of $d_{\rm max} = 5$ Gpc [4]. Therefore we will assume that the rate of the Sw $J1644+57$-like event is

$$\dot{R} = \frac{4\pi}{S_{\rm FOV}}\frac{1}{T}\frac{1}{(4/3)\pi d_{\rm max}^3} \simeq 2\times 10^{-12}\ {\rm Mpc^{-3}\,yr^{-1}} \quad (2)$$

and the energy injection rate in x rays is $\dot{\varepsilon}_{\rm X} = \dot{R}E_{\rm X} = 6\times 10^{41}\ {\rm erg\,Mpc^{-3}\,yr^{-1}}$. Following Ref. [11], the total energy in accelerated protons E_p can be parameterized by $E_p = \xi_{\rm p}E_{\rm X}$, where the nonthermal baryon loading factor $\xi_{\rm p}$ can be expressed by $\xi_{\rm p} = 10\varsigma_p\eta_e^{-1}(0.1/\epsilon_e)$, ϵ_e is the fraction of the shock internal energy that goes into nonthermal electrons, η_e is the radiative efficiency of these electrons, and ς_p is the proton acceleration efficiency. Modeling of the afterglows of gamma-ray bursts gives a typical value $\epsilon_e\simeq 0.1$ for relativistic shocks [12], so the typical value of $\xi_{\rm p}$ would be $\sim$10. Thus the differential energy injection rate in protons is

$$\varepsilon_p^2\frac{d\dot{n}}{d\varepsilon_p} = \frac{\xi_{\rm p}\dot{R}E_{\rm X}}{\ln(\varepsilon_{p,\rm max}/\varepsilon_{p,\rm min})} \simeq 6\times 10^{41}\xi_{\rm p,1}\ {\rm erg\,Mpc^{-3}\,yr^{-1}}, \quad (3)$$

where $\varepsilon_{p,\rm max}$ and $\varepsilon_{p,\rm min}$ are, respectively, the maximum and minimum energy of accelerated protons and we have used $\ln(\varepsilon_{p,\rm max}/\varepsilon_{p,\rm min}) = 10$ in the last step. For $\xi_{\rm p}$ of the order $\sim$10, this rate is much smaller than the required energy generation rate of cosmic rays per energy decade from $0.5-20\times 10^{44}\ {\rm erg\,Mpc^{-3}\,yr^{-1}}$ deduced by different authors [13]. Note, however, that it was only the presence of short, powerful bursts early on that alerted us to its presence, so we cannot exclude the possibility that many other similar, but rather less variable, events could be undetected.

Pion production.—Now we consider the neutrino emission produced by these protons interacting with the soft photons in the sources. Protons lose $\sim$20% of their energy at each $p\gamma$ interaction, dominated by the Δ resonance. Approximately half of the pions are charged and decay into high-energy neutrinos $\pi^+\rightarrow\mu^+ + \nu_\mu\rightarrow e^+ + \nu_e + \bar{\nu}_\mu + \nu_\mu$, with the energy distributed roughly equally among the decay products. The fraction of energy lost by protons to pions is $f_{p\gamma} = R/(\Gamma ct'_{p\gamma})$. Denoting by $n(\epsilon_\gamma)d\epsilon_\gamma$ the number density of photons in the energy range ϵ_γ to $\epsilon_\gamma + d\epsilon_\gamma$, the cooling time of protons in the shock comoving frame for $p\gamma$ process is given by

$$t'_{p\gamma} = \frac{c}{2\Gamma_p^2}\int_{\epsilon_{\rm th}}^{\infty} d\epsilon\sigma(\epsilon)K(\epsilon)\epsilon\int_{\epsilon/2\Gamma_p}^{\infty} dx x^{-2}n(x), \quad (4)$$

where $\Gamma_p = \varepsilon_p/\Gamma m_p c^2$, σ, and K are, respectively, the cross section and the inelasticity for $p\gamma$ process [14]. The spectrum from infrared to x-ray frequencies of Sw $J1644+57$ can be approximately described by a broken power law with $n(\epsilon_\gamma) = n_b(\epsilon_\gamma/\epsilon_b)^{-\alpha}$ for $\epsilon_\gamma < \epsilon_b$ and $n(\epsilon_\gamma) = n_b(\epsilon_\gamma/\epsilon_b)^{-\beta}$ for $\epsilon_\gamma > \epsilon_b$, where $\epsilon_b\sim 1$ KeV, $\alpha\simeq 2/3$, and $\beta\simeq 2$ [4]. Approximating the integral by the contribution from the resonance we obtain

$$f_{p\gamma}(\varepsilon_p)\simeq 0.35\frac{L_{\rm X,47.5}}{\Gamma_1^4 t_{v,2}\epsilon_{\rm b,KeV}}\begin{cases}(\varepsilon_p/\varepsilon_{pb})^{\beta-1} & (\varepsilon_p<\varepsilon_{pb})\\ (\varepsilon_p/\varepsilon_{pb})^{\alpha-1} & (\varepsilon_p>\varepsilon_{pb})\end{cases} \quad (5)$$

where $\varepsilon_{pb} = 0.15\ {\rm GeV}^2\Gamma^2/\epsilon_b = 1.5\times 10^{16}\Gamma_1^2\epsilon_{\rm b,KeV}^{-1}$ eV is the proton break energy. To include the effect of multipion production and high inelasticity (which increases from $\simeq 0.2$ at energies not far above the threshold to $\sim$0.5–0.6 at energies where multipion production dominates) at high energies [15,16], the above estimate of $f_{p\gamma}$ should be multiplied by factor of ~ 2. As the neutrino energy is $\sim 5\%$ of the proton energy, the neutrino flux will peak at $\varepsilon_{\nu b} \simeq 7.5 \times 10^{14} \Gamma_1^2 \epsilon_{\mathrm{b,KeV}}^{-1}$ eV.

In the modeling of the SED of Sw $J1644+57$, upper limits from Fermi and the Very Energetic Radiation Imaging Telescope Array System require $\Gamma < 20$ in the x-ray emitting region [4]. It is useful to express $f_{p\gamma}$ as a function of the pair production optical depth $\tau_{\gamma\gamma}$. The optical depth for pair production of a photon of energy ϵ_h is $\tau_{\gamma\gamma}(\epsilon_h) = \frac{R}{\Gamma l_{\gamma\gamma}} = \frac{R}{\Gamma} \frac{\sigma_T}{16} \frac{U_\gamma \epsilon_h}{\Gamma (m_e c^2)^2}$, where $l_{\gamma\gamma}$ is the mean free path. The fraction of energy lost by protons to pions is $f_{p\gamma} \simeq \frac{R}{\Gamma} \frac{U_\gamma}{2\varepsilon_p'} \sigma_{p\gamma} K_{p\gamma} (\varepsilon_p / \varepsilon_{pb})^{\alpha - 1}$ for protons with energy $\varepsilon_p > \varepsilon_{pb}$. Thus, we have [17]

$$f_{p\gamma}(\varepsilon_p) \simeq 0.5 \tau_{\gamma\gamma}(100 \text{ MeV}) \left(\frac{\varepsilon_{pb}}{1.5 \times 10^{16} \text{ eV}} \right) \left(\frac{\varepsilon_p}{\varepsilon_{pb}} \right)^{\alpha - 1}. \tag{6}$$

Modeling of the SED of Sw $J1644+57$ requires $\tau_{\gamma\gamma}(100 \text{ MeV}) > 1$ [4], so we conclude that a significant fraction ($> 50\%$) of the energy of protons accelerated to energies larger than the break energy, $\varepsilon_{pb} \sim 10^{16}$ eV, would be lost to pion production.

Meson cooling. — The neutrino production efficiency will be reduced if the secondary mesons suffer from cooling before decaying to neutrinos and other products [18]. The pions and muons suffer from radiative cooling due to both synchrotron emission and inverse-Compton emission. The total radiative cooling time for pions is $t'_{\pi,\mathrm{rad}} = 3 m_\pi^4 c^3 / [4 \sigma_T m_e^2 \epsilon'_\pi U'_B (1 + f_{\mathrm{IC}})] \simeq 2 \times 10^6 (\epsilon'_\pi / 1 \text{ TeV})^{-1} \epsilon_{B,-1}^{-1} L_{w,48}^{-1} R_{14.5}^2 \Gamma_1^2$ s, where U'_B is the energy density of the magnetic filed in the shock region and $f_{\mathrm{IC}} \le 1$ is the correction factor accounting for the inverse-Compton loss. By comparing this cooling time $t'_{\pi,\mathrm{rad}}$ with the lifetime of pions $\tau'_\pi = \gamma_\pi \tau = 1.9 \times 10^{-4} (\epsilon'_\pi / 1 \text{ TeV})$ s in the shock comoving frame, where γ_π and τ are the pion Lorentz factor and proper lifetime, one can find a critical energy (in the observer frame) for pions, above which the effect of radiative cooling starts to suppress the neutrino flux,

$$\varepsilon_{\pi,\mathrm{rad}} = 6.3 \times 10^5 \epsilon_{B,-1}^{-1/2} L_{w,49}^{-1/2} \Gamma_1^4 t_{v,2} \text{ TeV}. \tag{7}$$

Similarly, by comparing the radiative cooling time $t'_{\mu,\mathrm{rad}}$, with the lifetime of muons τ'_μ, one can obtain a critical energy for muons, above which the effect of radiative cooling starts to suppress the antimuon neutrino flux,

$$\varepsilon_{\mu,\mathrm{rad}} = 3.2 \times 10^4 \epsilon_{B,-1}^{-1/2} L_{w,49}^{-1/2} \Gamma_1^4 t_{v,2} \text{ TeV}. \tag{8}$$

The above estimates lead us to conclude that the neutrino flux below $\sim 10^{16}$ eV is not affected by the meson cooling for typical parameters of TDFs. At higher energies, however, pion cooling and muon cooling will suppress the neutrino flux by a factor approximately given, respectively, by [19]

$$\zeta_\pi = \min\{t'_{\pi,\mathrm{rad}} / \tau'_\pi, 1\}, \qquad \zeta_\mu = \min\{t'_{\mu,\mathrm{rad}} / \tau'_\mu, 1\}. \tag{9}$$

The spectrum and flux of the neutrino flare.—The fluence spectrum of the muon neutrino ($\nu_\mu + \bar{\nu}_\mu$) emission from one TDF is

$$\begin{aligned} \varepsilon_\nu^2 \Phi_\nu &= \varepsilon_p^2 \frac{dn_p}{d\varepsilon_p} \frac{f_{p\gamma} \zeta_\pi (1 + \zeta_\mu)}{8} \\ &= \frac{E_{\mathrm{X}}}{32 \pi d_L^2} \frac{\xi_{\mathrm{p}} f_{p\gamma} \zeta_\pi (1 + \zeta_\mu)}{\ln(\varepsilon_{p,\max} / \varepsilon_{p,\min})}, \end{aligned} \tag{10}$$

where $\varepsilon_p^2 \frac{dn_p}{d\varepsilon_p} = E_p / [4 \pi d_L^2 \ln(\varepsilon_{p,\max} / \varepsilon_{p,\min})]$ is the differential proton fluence produced by one TDF of total energy $E_p = \xi_{\mathrm{p}} E_{\mathrm{X}}$ in protons and the factor $1/8$ represents that the neutrinos produced by pion decay carry one-eighth of the energy lost by protons to pion production, since charged and neutral pions are produced with roughly equal probability and muon neutrinos carry roughly one-fourth of the pion energy in pion decay. Figure 1 shows the expected muon neutrino fluence spectra from Sw $J1644+57$, obtained by using the Lorentzian form for the photopion production cross section at the resonance peak plus a component contributed by multipion production at higher energies [20] in calculating $t'_{p\gamma}$ with Eq. (4). The initial rise in the spectrum at low energies is caused by the increasing pion production efficiency with energy, while the mild steepening and sharp steepening seen at higher energies are caused by muon cooling and pion cooling, respectively.

Now we estimate the number of neutrinos that can be detected from one TDF event similar to Sw $J1644+57$. The detection efficiency in water or ice of ultrarelativistic

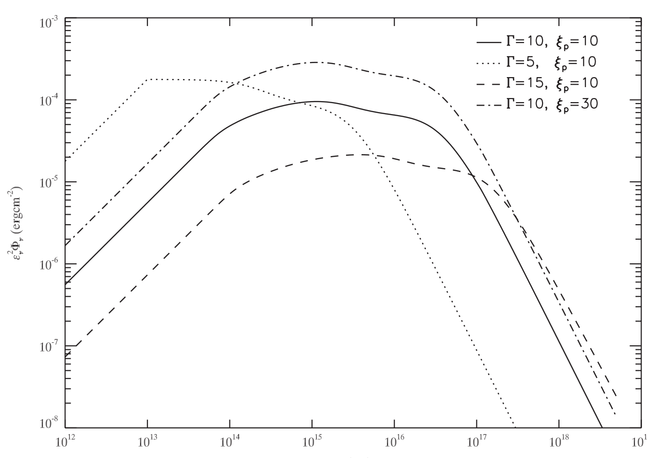

FIG. 1. The expected muon neutrino ($\nu_\mu + \bar{\nu}_\mu$) spectra from Sw $J1644+57$ for different jet parameters. In all lines, $S_{\mathrm{X}} = 7 \times 10^{-4}$ erg cm^{-2}, $L_{\mathrm{X}} = 3 \times 10^{47}$ erg s^{-1}, $t_v = 100$ s, $\epsilon_b = 1$ KeV and $\epsilon_B = 0.1$ are used.

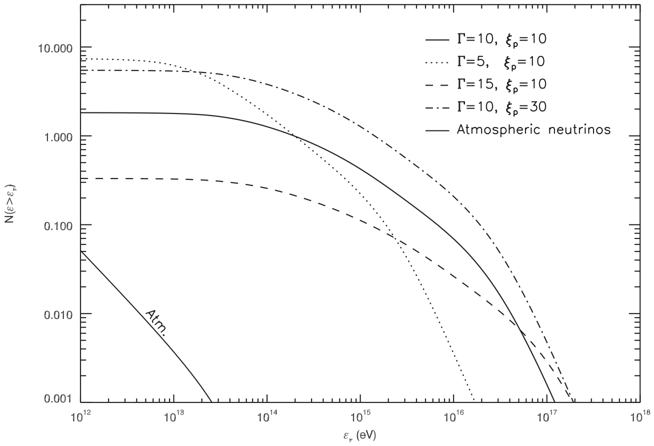

FIG. 2. The expected number of neutrinos ($\nu_\mu + \bar{\nu}_\mu$) above a certain energy detected from Sw $J1644+57$ by Km^3-scale neutrino detectors such as IceCube. The thin solid line represents the background atmospheric neutrinos. The same parameters as in Fig. 1 are used.

upward-going muon neutrinos with energies ε_ν is $P_{\nu\mu} \simeq 7 \times 10^{-5}(\varepsilon_\nu/10^{4.5}\ \mathrm{GeV})^\kappa$, where $\kappa = 1.35$ for $\varepsilon_\nu < 10^{4.5}$ GeV, and $\kappa = 0.55$ for $\varepsilon_\nu > 10^{4.5}$ GeV [21]. For neutrino fluence spectrum parameterized by $\varepsilon_\nu^2 \Phi_\nu$, the number of ν_μ and $\bar{\nu}_\mu$ above a certain energy $\varepsilon_{\nu 0}$ detected by a km^3-scale neutrino detector, such as IceCube, with area $A = 10^{10} A_{10}$ cm^2 is

$$N_\nu(>\varepsilon_{\nu 0} = 10^{4.5}\ \mathrm{GeV}) = \int_{\varepsilon_{\nu 0}}^{\infty} \Phi_\nu P_{\nu\mu} A d\varepsilon_\nu \simeq 2\xi_{\mathrm{p},1} A_{10} f_{p\gamma}(\varepsilon_{pb}) \left(\frac{E_X}{3 \times 10^{53}\ \mathrm{erg}}\right)\left(\frac{d_L}{1.8\ \mathrm{Gpc}}\right)^{-2}, \tag{11}$$

where we have used $\ln(\varepsilon_{p,\max}/\varepsilon_{p,\min}) = 10$, $\zeta_\pi \simeq 1$, $\zeta_\mu \simeq 1$, $\alpha = 2/3$, and $\beta = 2$ in the last step. As $f_{p\gamma} > 0.5$, we expect ≥ 1 neutrinos above 30 TeV could be detected from Sw $J1644+57$ by Km^3-scale detectors for $\xi_{\mathrm{p}} = 10$. A careful calculation of the number of neutrinos above a certain energy as a function of the neutrino energy is shown in Fig. 2. A lower bulk Lorentz factor or a higher wind luminosity (i.e. a larger ξ_{p}) leads to a larger number of neutrinos that can be detected. If a similar event to Sw $J1644+57$ occurs at a closer distance (e.g. at $z = 0.2$) in future, more neutrinos would be detected as well.

Neutrino detection from TDFs can be assured only if the number of background counts is smaller than 1. The number of atmospheric neutrinos above 30 TeV expected in the direction of the source during the flare period is

$$N_{\mathrm{atm}}(>10^{4.5}\ \mathrm{GeV}) = \int_{10^{4.5}\ \mathrm{GeV}}^{\infty} d\varepsilon_\nu \int d\Omega \int dt \frac{F_\nu^{\mathrm{atm}}}{\varepsilon_\nu^2} P_{\nu\mu} A \simeq 3 \times 10^{-3}\left(\frac{\Delta T}{10^6\ \mathrm{s}}\right) A_{10} \left(\frac{\theta}{1^\circ}\right)^2, \tag{12}$$

where ΔT is characteristic duration of the TDF, $\theta \simeq 0.5°$–$0.6°$ is the angular resolution of the neutrino detector at 30 TeV-PeV [22], F_ν^{atm} is the cosmic-ray induced atmospheric neutrino background flux. We fit the atmospheric neutrino data measured by IceCube [23] with a single power-law function, which gives $F_\nu^{\mathrm{atm}} \simeq 4.7 \times 10^{-8}\ \mathrm{erg\,cm^{-2}\,s^{-1}\,sr^{-1}}(\varepsilon_\nu/1\ \mathrm{TeV})^{-\beta}$ with $\beta \simeq 1.44$ in the energy range of 0.1–400 TeV. Since the number of atmospheric neutrinos above 30 TeV expected in the direction of the source during the flare period is much smaller than 1, a detection of two neutrinos at energies above 30 TeV from TDF sources will be highly significant.

Discussions.—Neutrino emission has been predicted to be produced by relativistic jets in gamma-ray bursts [11,17,24], active galactic nuclei [25] and microquasars [26]. Observations of Sw $J1644+57$ suggest that powerful jets are formed in TDFs, which have larger fluence in x/γ-rays than the brightest gamma-ray bursts and have higher x-ray luminosity than active galactic nuclei. There are three factors that are favorable for bright neutrino emission produced in such TDFs: (1) large fluence in the x-ray emission, which suggests large fluence in accelerated protons; (2) high pion production efficiency as implied by the high opacity of high-energy gamma-rays, which leads to a high fraction of the proton energy lost into secondary pions; (3) insignificant radiative cooling of secondary pions and muons, which leads to an almost flat neutrino spectrum up to $\sim 10^{16}$ eV. The main uncertainty lies in the ratio between the energy density of protons and the energy density in x rays. In the magnetic-field–dominated jet model, the proton energy density is subdominant, so the neutrino flux would be low, whereas in the matter-dominated jet model, we expect one to several neutrinos detectable by Km^3-scale neutrino detectors from TDFs similar to Sw $J1644+57$. The Swift Burst Alert Telescope, with a field of view of $4\pi/7$ sr, has detected one TDF in 7 yr, so the all-sky rate of TDFs would be one event in every 1 yr. For neutrino detectors such as IceCube that have a 2π sr field of view, we expect one TFD event in the field of view of IceCube every 2 yr, if the electromagnetic counterparts can be identified. Therefore neutrino observations provide a promising approach to diagnose the composition of the jets resulted from tidal disruption of stars by massive black holes in the galaxy center.

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[1] M. J. Rees, Nature (London) **333**, 523 (1988); C. R. Evans and C. S. Kochanek, Astrophys. J. Lett. **346**, L13 (1989); J. Wang and D. Merritt, Astrophys. J. **600**, 149 (2004); L. E. Strubbe and E. Quataert, Mon. Not. R. Astron. Soc. **400**, 2070 (2009).
[2] S. Komossa and J. Greiner, Astron. Astrophys. **349**, L45 (1999); S. Gezari *et al.*, Astrophys. J. **676**, 944 (2008); S. van Velzen *et al.*, arXiv:1009.1627.
[3] K. S. Cheng and Y. Lu, Mon. Not. R. Astron. Soc. **320**, 235 (2001); K. S. Cheng *et al.*, Astrophys. J. **645**, 1138 (2006); , Astron. Astrophys. **473**, 351 (2007); D. Giannios and B. D. Metzger, arXiv:1102.1429.
[4] D. Burrows *et al.*, arXiv:1104.4787.
[5] J. Bloom *et al.*, Science **333**, 203 (2011).
[6] A. Levan *et al.*, Science **333**, 199 (2011).
[7] U. Barres de Almeida and A. De Angelis, arXiv:1104.2528; L. Shao *et al.*, Astrophys. J. Lett. **L33**, 734 (2011); J. K. Cannizzo *et al.*, arXiv:1105.2816; A. Socrates, arXiv:1105.2557.
[8] A. Zauderer *et al.* (unpublished).
[9] P. Mészáros and M. J. Rees, Mon. Not. R. Astron. Soc. **269**, 41P (1994); B. Paczyński and G. Xu, Astrophys. J. **427**, 708 (1994); M. Spada, G. Ghisellini, D. Lazzati, and A. Celotti, Mon. Not. R. Astron. Soc. **325**, 1559 (2001); M. Böttcher and C. D. Dermer, Astrophys. J. **711**, 445 (2010).
[10] G. R. Farrar and A. Gruzinov, Astrophys. J. **693**, 329 (2009).
[11] K. Murase and S. Nagataki, Phys. Rev. D **73**, 063002 (2006); Phys. Rev. Lett. **97**, 051101 (2006).
[12] A. Panaitescu and P. Kumar, Astrophys. J. **560**, L49 (2001).
[13] B. Katz, R. Budnik, and E. Waxman, J. Cosmol. Astropart. Phys. 3 (2009) 020; S. T. Scully and F. W. Stecker, Astropart. Phys. **16**, 271 (2002); V. Berezinsky, Adv. Space Res. **41**, 2071 (2008).
[14] F. W. Stecker, Phys. Rev. Lett. **21**, 1016 (1968).
[15] A. Mücke *et al.*, Pub. Astron. Soc. Aust. **16**, 160 (1999).
[16] A. M. Atoyan and C. D. Dermer, Astrophys. J. **586**, 79 (2003).
[17] E. Waxman and J. Bahcall, Phys. Rev. Lett. **78**, 2292 (1997).
[18] J. P. Rachen and P. Mészáros, Phys. Rev. D **58**, 123005 (1998).
[19] S. Razzaque, P. Mészáros, and E. Waxman, Phys. Rev. Lett. **93**, 181101 (2004).
[20] A. Mücke *et al.*, Nucl. Phys. B, Proc. Suppl. **80**, 8 (2000).
[21] T. K. Gaisser, F. Halzen, and T. Stanev, Phys. Rep. **258**, 173 (1995); S. Razzaque, P. Mészáros, and E. Waxman, Phys. Rev. D **69**, 023001 (2004); K. Ioka *et al.*, Astrophys. J. **633**, 1013 (2005).
[22] J. Ahrens *et al.*, Astropart. Phys. **20**, 507 (2004).
[23] R. Abassi *et al.*, Phys. Rev. D **83**, 012001 (2011).
[24] C. D. Dermer and A. Atoyan, Phys. Rev. Lett. **91**, 071102 (2003); S. Razzaque, P. Mészáros, and E. Waxman, Phys. Rev. Lett. **90**, 241103 (2003); X. Y. Wang and Z. G. Dai, Astrophys. J. **691**, L67 (2009).
[25] K. Mannheim, T. Stanev, and P. L. Biermann, Astron. Astrophys. **260**, L1 (1992); C. D. Dermer and A. Atoyan, Phys. Rev. Lett. **87**, 221102 (2001); F. W. Stecker, Phys. Rev. D **72**, 107301 (2005); J. K. Becker and P. L. Biermann, Astropart. Phys. **31**, 138 (2009).
[26] A. Levinson and E. Waxman, Phys. Rev. Lett. **87**, 171101 (2001); F. Aharonian *et al.*, J. Phys. Conf. Ser. **39**, 408 (2006).