Probing Microquasars with TeV Neutrinos

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The jets associated with galactic microquasars are believed to be ejected by accreting stellar mass black holes or neutron stars. We show that if the energy content of the jets in the transient sources is dominated by electron-proton plasma, then a several hour outburst of 1–100 TeV neutrinos produced by photomeson interactions should precede the radio flares associated with major ejection events. Several neutrinos may be detected during a single outburst by a 1 km^2 detector, thereby providing a powerful probe of microquasar jet physics.

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Microquasars are galactic jet sources associated with some classes of x-ray binaries involving both neutron stars and black hole candidates (albeit with some notable differences between these two classes) [1–3]. During active states, the x-ray flux and spectrum can vary substantially between different substates, with a total luminosity that, during the so-called very high states, often exceeds the Eddington limit (typically a few times 10^{38} ergs s^{-1} in those sources). Their activity involves ejection of jets with kinetic power that appears to constitute a considerable fraction of the liberated accretion energy (in some cases the minimum jet power inferred exceeds the peak x-ray luminosity [4,5]), and that gives rise to intense radio and IR flares. Radio monitoring of some x-ray transients has revealed superluminal motions in currently three of the sources [6–8], indicating that at least in these objects the jets are relativistic. The Lorentz factors of the radio emitting blobs have been estimated to be (coincidently) $\Gamma \sim 2.5$ in the two superluminal sources, GRS 1915 + 105 [6] and GRO J1655-40 [7], and somewhat smaller ($\Gamma \sim 2$) in the third one, XTE J1819 [8]. Whether these are representative values or merely the result of selection effects is unclear at present; the class of microquasars may contain sources with much larger Γ, rendered invisible by beaming away from us [4].

The temporal behavior of microquasars appears to be rather complex. They exhibit large amplitude variations over a broad range of time scales and frequencies, with apparent connections between the radio, IR, and soft/hard x-ray fluxes [9–12]. The characteristics of the multi-waveband behavior depend on the state of the source, that is, whether the source is in a very high, soft/high, or low/hard state [3]. The ejection episodes are classified into several classes according to the brightness of synchrotron emission produced in the jet and the characteristic time scale of the event [13]. The duration of major ejection events (class A) is typically on the order of days, while that of less powerful flares (classes B and C) is correspondingly shorter (minutes to hours). The correlations between the x-ray and synchrotron emission clearly indicates a connection between the accretion process and the jet activity. Whether radio and IR outbursts represent actual ejection of blobs of plasma or, alternatively, formation of internal shocks in a quasisteady jet is unclear (cf. Ref. [14]). In any case, since the overall time scale of outbursts (minutes to days) is much longer than the dynamical time of the compact object (milliseconds), it is likely that shocks will continuously form during the ejection event over a range of time scales that encompasses the dynamical time, owing to fluctuations in the parameters of the expelled wind, leading to dissipation of a substantial fraction of the bulk energy at relatively small radii. The extremely rapid variations of the x-ray flux often seen in these sources supports this view (e.g., Ref. [15]). If a fraction of at least a few percent of the jet power is tapped for acceleration of electrons to very high energies, then emission of high-energy gamma rays is anticipated, in addition to the observed radio and IR emission [4,5]. Energetic Gamma Ray Experiment Telescope (EGRET) upper limits for some x-ray novae have been considered in this context [16]. A recently discovered microquasar, which appears to have a persistent radio jet [17], seems to coincide with an unidentified EGRET source having a total luminosity in excess of its x-ray luminosity.

The content of jets in microquasars is yet an open issue. The synchrotron emission both in the radio and in the IR is consistent with near equipartition magnetic field, which is also implied by minimum energy considerations [4]. However, the dominant energy carrier in the jet is presently unknown (with the exception of the jet in SS433). Scenarios whereby energy extraction is associated with spin down of a Kerr black hole favor $e^{\pm}$ composition (although baryon entrainment is an issue). However, the pair annihilation rate inferred from the estimated jet power implies electromagnetic domination on scales smaller than roughly 10^9 cm in the superluminal sources and requires a transition from electromagnetic to particle dominated flow above the annihilation radius by some unspecified mechanism [4]. Alternatively, in scenarios in which an initial rise of the x-ray flux leads to ejection of the inner part of the accretion disk, as widely claimed to be suggested by the anticorrelation between the x-ray and radio flares seen during major ejection events (cf. Ref. [4] for a different interpretation), e-p jets are expected to be produced. A possible diagnostic of e-p jets is the presence of Doppler-shifted spectral lines, such as the H_α line as seen in SS433. The detection of such lines from jets having a Lorentz factor well in excess of unity (as is the case in the superluminal microquasars) may, however, be far more difficult than in SS433, as the lines are anticipated to be very broad ($\Delta\lambda/\lambda \gtrsim 0.1$). Furthermore, the conditions required to produce detectable flux in such sources may be far more extreme than in SS433 [18]. Here we propose another diagnostic of hadronic jets, namely, emission of TeV neutrinos. As shown below, for typical parameters the neutrinos are produced on scales much smaller than the IR and radio emission and, therefore, can provide a probe of the innermost structure of microquasar jets.

The picture envisaged is the following: We suppose that on sufficiently small scales ($\lesssim 10^{11}$ cm) a significant fraction of the energy liberated during ejection events dissipates, e.g., through the formation of internal collisionless shocks [19], leading to the acceleration of a nonthermal power-law distribution of protons (and electrons) up to the maximum energy achievable. The fraction of energy carried by the power-law component of electrons may depend sensitively on the injection process, but is generally thought to constitute a few percent. In the following, we denote by η_P the fraction of the total burst energy that is tapped for the acceleration of protons to nonthermal energies. We emphasize that η_p represents essentially the product of the fraction of total burst energy that dissipates behind the shocks (which depends primarily on the duty cycle of shock formation) and the efficiency at which protons are accelerated to nonthermal energies by a single shock (that may approach 100%). Observations of γ-ray bursts, which are believed to arise from a similar process of jet kinetic energy dissipation (albeit with different Lorentz factors and luminosities [20]), suggest that a significant fraction of the jet energy is dissipated and converted to a power-law distribution of accelerated electrons [21].

As shown below, synchrotron emission by electrons contributes large opacity to photomeson production over a range of radii that depends on the strength of the magnetic field. (Since the photomeson opacity is contributed predominantly by thermal electrons, the analysis outlined below is rendered independent of the details of electron injection to nonthermal energies.) Between 12% and 25% of the injected proton energy is converted, by virtue of the large optical depth, into muon neutrinos having a flat spectrum in the range between 1 and $\sim$100 TeV, which then freely escape the system.

Consider a jet of kinetic power $L_j = 10^{38} L_{j38}$ erg s^{-1}, and opening angle $\theta = 0.1\theta_{-1}$, with a corresponding proper energy density $U_j = L_j/(\pi\theta^2 r^2 \Gamma^2 c)$, propagating in the background of x-ray photons emitted by the accretion disk, with luminosity $L_x = 10^{38} L_{x38}$ erg s^{-1}. A significant fraction of the jet kinetic energy is assumed to dissipate via mildly relativistic internal shocks in the jet at radii $r \gtrsim \Gamma^2 c\delta t = 3 \times 10^7 \Gamma^2 \delta t_{-3}$ cm, where $\delta_t = 10^{-3}\delta t_{-3}$ s is the source dynamical time.

The Thomson optical depth across the jet is given by $\tau_T = \sigma_T n_p r\theta\Gamma \simeq 0.05 L_{j38}/(\Gamma r_8 \theta_{-1})$, where $r = 10^8 r_8$ cm and the proper proton number density in the jet is $n_p \simeq U_j/(m_p c^2)$. External photons can therefore penetrate the jet and interact with accelerated protons at radii larger than roughly 10^7 cm. The proton energy, in the jet frame, for which interaction with external x-ray photons is at the Δ resonance is $\epsilon_{p,\Delta} = 0.3$ GeV$^2/\Gamma\epsilon_x = 3 \times 10^{14}(\Gamma\epsilon_x/1\ \text{keV})^{-1}$ eV, where ϵ_x is the observed photon energy.

The time available, in the jet frame, for photopion interactions is $\approx r/\Gamma c$. Thus, the optical depth to photopion production at the Δ resonance, contributed by the external photons, is $\tau_{p\gamma} \simeq c\sigma_{\text{peak}} n_x r/\Gamma c = 1\Gamma L_{x38} r_8^{-1}(\epsilon_x/1\ \text{keV})^{-1}$, where $\sigma_{\text{peak}} = 5 \times 10^{-28}$ is the cross section at the Δ resonance and n_x is the comoving number density of background x-ray photons. Thus, photomeson interactions with external x-ray photons at the innermost dissipation radii will convert a large fraction of the accelerated proton energy to high-energy pions.

Additional contribution to the photomeson optical depth comes from synchrotron photons produced inside the jet by the thermal electrons. For a relativistic shock, the mean electron energy assuming rapid equilibration time is $\sim 0.5(\Gamma_s - 1)m_p c^2$, where Γ_s is the Lorentz factor of the shock. For the microquasars $\Gamma_s - 1 \sim 1$ typically. This then yields for the (jet frame) peak energy of the synchrotron spectrum,

$$\epsilon_{\gamma,\text{peak}} \simeq 50 \frac{\xi_{-1}^{1/2} L_{j38}^{1/2}}{\Gamma r_8 \theta_{-1}}\ \text{keV}, \tag{1}$$

where $\xi = 10^{-1}\xi_{-1}$ is the equipartition parameter, defined through the relation $B = \sqrt{8\pi\xi U_j}$. Note that $\epsilon_{\gamma,\text{peak}}$ reduces to infrared energies at a radius of $\sim 10^{13}$ cm, with a corresponding light crossing time of order minutes, compatible with the variability time of IR baby flares.

Assuming the total energy density of synchrotron photons to be on the order of U_j yields for the corresponding energy-dependent number density, defined as $n_{\text{syn}}(\epsilon) = \epsilon(dn_{\text{syn}}/d\epsilon)$,

$$\begin{aligned} n_{\text{syn}}(\epsilon) &\simeq \frac{U_j}{\epsilon_{\gamma,\text{peak}}}\left(\frac{\epsilon}{\epsilon_{\gamma,\text{peak}}}\right)^{-\alpha} \\ &\simeq 10^{20}\left(\frac{\epsilon}{\epsilon_{\gamma,\text{peak}}}\right)^{-\alpha} \frac{L_{j38}^{1/2}}{r_8 \Gamma \theta_{-1} \xi_{-1}^{1/2}}\ \text{cm}^{-3}, \end{aligned} \tag{2}$$

where $\alpha = 1/2$ for $\epsilon < \epsilon_{\text{peak}}$ and (assuming efficient electron injection) $\alpha = 1$ for $\epsilon > \epsilon_{\text{peak}}$. The spectrum at $\epsilon < \epsilon_{\text{peak}}$ results from the fast cooling of electrons. The ratio of electron cooling time, $t_{\text{syn}} \approx 6\pi m_e^2 c/\sigma_T B^2 m_p$, to the dynamical time, $r/\Gamma c$, is $\Gamma c t_{\text{syn}}/r \simeq 10^{-5}\Gamma^3 r_8 \theta_{-1}^2/(L_{j38}\xi_{-1})$. We note that at radii $r \lesssim 10^9$ cm, inverse-Compton scattering of electrons is suppressed since scattering is in the Klein-Nishsina regime for photons carrying significant fraction of the synchrotron energy density. At these radii, the fast cooling of electrons indeed implies that the energy density of synchrotron photons is of order U_j. At larger radii, the energy density of synchrotron photons will be somewhat suppressed due to inverse-Compton emission.

A rough estimate of the energy loss rate of protons due to photomeson interactions gives $t_{p\gamma}^{-1} \sim \sigma_{\rm peak} n_{\rm syn} c(\Delta\epsilon_p/\epsilon_p)$, where $\Delta\epsilon_p/\epsilon_p \simeq 0.2$ is the average fractional energy loss in a single collision. Equating the latter with the acceleration rate $\sim eBc/\epsilon_p$, we obtain for the maximum proton energy

$$\epsilon_{p,\rm max} \simeq 5 \times 10^{15} \xi_{-1}^{1/2} (\Gamma r_8 \theta_{-1})^{1/3} L_{j38}^{-1/6} \text{ eV}. \quad (3)$$

We emphasize that the proton energy cannot in any case exceed the upper limit imposed by the requirement that the protons be confined to the system; viz., $\epsilon_p < \theta eBr \simeq 2 \times 10^{16} \xi_{-1}^{1/2} L_{j38}^{1/2} \Gamma^{-1}$ eV. Comparing the latter with Eq. (8) then implies that the maximum proton energy is limited by photomeson losses at radii $r_8 < 10^2 L_{j38}^2 \Gamma^{-4} \theta_{-1}^{-1}$ and by confinement at larger radii.

The comoving proton energy for which interaction with synchrotron peak photons is at the Δ resonance is

$$\epsilon_{p,\rm peak} = 10^{13} \frac{\Gamma r_8 \theta_{-1}}{(\xi_{-1} L_{38})^{1/2}} \text{ eV}. \quad (4)$$

Consequently, protons may be accelerated to energy exceeding that required for photomeson interaction with synchrotron peak photons for radii

$$r \lesssim r_c = 5 \times 10^{11} \frac{L_{j38}^{1/2} \xi_{-1}}{\Gamma \theta_{-1}} \min(L_{j38}/\Gamma, \xi_{-1}^{1/2}) \text{ cm}. \quad (5)$$

From Eqs. (3) and (5) it is seen that for our choice of parameters protons can be accelerated to energies $\sim 10^{16} \xi_{-1}$ eV in the region where photomeson interactions take place. The contribution of synchrotron photons to the photomeson optical depth is

$$\tau_{p\gamma} \simeq c\sigma_{\rm peak} n_{\rm syn} \frac{r}{\Gamma c} \simeq \frac{10}{\Gamma^2 \theta_{-1}} \left(\frac{L_{j38}}{\xi_{-1}}\right)^{1/2} \left(\frac{\epsilon_p}{\epsilon_{p,\rm peak}}\right)^{\beta}, \quad (6)$$

where $\beta = 1/2$ for $\epsilon_p > \epsilon_{p,\rm peak}$ and $\beta = 1$ for $\epsilon_p < \epsilon_{p,\rm peak}$.

Equation (6) implies that at all radii the photomeson optical depth is large for protons of energy exceeding the threshold energy (4) for interaction with synchrotron peak photons. We therefore expect most of the energy of these protons to be lost to pion production. Neutral pions and charged pions are produced in the photomeson interaction with roughly equal probabilities, and the decay of a charged pion produces three neutrinos, $\pi^+ \rightarrow \mu^+ + \nu_\mu \rightarrow e^+ + \nu_e + \overline{\nu}_\mu + \nu_\mu$, each carrying $\sim 5\%$ of the initial proton energy. Because of the high photomeson optical depth, neutrons produced in photopion interactions will lose their energy by photomeson interactions on a time scale shorter than the neutron lifetime. This implies similar production rates of both negatively and positively charged pions and, hence, similar production rates of ν_μ and $\bar{\nu}_\mu$.

High-energy pions and muons may loose a significant fraction of their energy by inverse-Compton scattering prior to decaying. Inverse-Compton scattering of synchrotron photons by high-energy pions is in the Klein-Nishina regime for

$$\epsilon_\pi \gtrsim \epsilon_{\pi,c} = \frac{(m_\pi c^2)^2}{\epsilon_{\gamma,\rm peak}} \simeq 0.4 \frac{\Gamma r_8 \theta_{-1}}{(L_{j38} \xi_{-1})^{1/2}} \text{ TeV}. \quad (7)$$

Comparing this equation with Eq. (4), and noting that the energy of a pion produced by photopion interaction is $\approx 20\%$ of the initial proton energy, we find that pions produced by protons of energy exceeding the threshold energy satisfy $\epsilon_\pi \gtrsim \epsilon_{\pi,c}$. Thus, the pion inverse-Compton scattering energy loss time is

$$\tau_{\rm IC} \approx \frac{3}{8c\sigma_T n_{\rm syn}} \frac{\epsilon_\pi \epsilon_{\gamma,\rm peak}}{(m_e c^2)^2} = 3 \times 10^{-2} \xi_{-1} \left(\frac{\epsilon_\pi}{1 \text{ TeV}}\right) \text{s}, \quad (8)$$

and the ratio of the pion decay time to energy loss time is $\tau_{\rm IC}/\tau_{\rm decay} \approx 10^2 \xi_{-1}$. We therefore expect pions to decay prior to significant energy loss. However, muons, the lifetime of which is ≈ 100 times longer, may lose a significant fraction of their energy before decaying. In the following we therefore conservatively assume that a single high-energy ν_μ (or $\bar{\nu}_\mu$) is produced in a single photopion interaction of a proton (or neutron), corresponding to conversion of $1/8$ of the energy lost to pion production to muon neutrinos. In terms of η_p, the fraction of jet energy injected as a power-law distribution of protons, the flux at Earth of ν_μ and $\bar{\nu}_\mu$ can be expressed as

$$\mathcal{F}_{\nu_\mu} \simeq \frac{1}{2} \eta_p \Gamma^{-1} \delta^3 \frac{L_j/8}{4\pi D^2} = 0.5 \times 10^{-9} \eta_{p,-1} \Gamma^{-1} \delta^3 D_{22}^{-2} L_{j38} \text{ erg s}^{-1} \text{ cm}^{-2}, \quad (9)$$

where δ is the Doppler factor, $D = 10^{22} D_{22}$ cm is the distance to the source, and $\eta_p = 0.1\eta_{p,-1}$. The factor $1/2$ is due to the fact that for a flat power-law proton distribution, and for the characteristic parameters invoked, roughly half the energy of the power-law component is carried by protons which lose most of their energy to pion production [see Eqs. (4) and (6)].

Because of the high photomeson optical depth at energies $\epsilon_p > \epsilon_{p,\rm peak}$, we expect a neutrino spectrum similar to the proton spectrum, $dn_\nu/d\epsilon_\nu \propto \epsilon_\nu^{-2}$, above ~ 1 TeV. The probability that a muon neutrino will produce a high-energy muon in a terrestrial detector is [22] $P_{\nu\mu} \approx 1.3 \times 10^{-6} E_{\nu,\rm TeV}^{\beta}$, with $\beta = 2$ for $E_{\nu,\rm TeV} < 1$ and $\beta = 1$ for $E_{\nu,\rm TeV} > 1$. Thus, for a flat neutrino spectrum above 1 TeV, the muon flux at the detector is $\approx (P_0/E_0)\mathcal{F}_{\nu_\mu}$, where $P_0/E_0 = 1.3 \times 10^{-6}$ TeV^{-1}. The number of events detected in a burst of energy $E = 10^{43} E_{43}$ erg is therefore

$$N_\mu \simeq 0.2\eta_{p,-1}\Gamma^{-1}\delta^3 D_{22}^{-2} E_{43}(A/1 \text{ km}^2), \tag{10}$$

where A is the effective detector area.

For sources directed along our sight line $\Gamma^{-1}\delta^3 \sim 8\Gamma^2$ may exceed 100. Thus, if the fraction η_p exceeds a few percent, several neutrinos can be detected during a typical outburst even from a source at 10 kpc. The typical angular resolution of the planned neutrino telescopes at TeV energies (e.g., [23]) should be $\theta \sim 1$ deg. The atmospheric neutrino background flux is $\Phi_{\nu,\text{bkg}} \sim 10^{-7}\epsilon_{nu,\text{TeV}}^{-2.5}/\text{cm}^2\,\text{s}\,\text{sr}$, implying a number of detected background events $N_{\text{bkg}} \sim 3 \times 10^{-2}(\theta/\text{deg})^2 t_{\text{day}}\,\text{km}^{-2}$ per angular resolution element over a burst duration $1 t_{\text{day}}$ day. The neutrino signals above $\sim$10 TeV from a typical microquasar outburst should therefore be easily detected above the background.

The duration of the neutrino burst should be of the order of the blob's ejection time. It should precede the associated radio outburst, or the emergence of a new superluminal component, that originates from larger scales ($\sim 10^{15}$ cm), by several hours.

As an example, let us consider the 19 March 1994 outburst observed in GRS 1915 + 105 [6]. This source is at a distance of about 12 kpc, or $D_{22} = 3.5$. From the proper motions measured with the Very Large Array, a speed of 0.92 c and an angle to the line of sight of $\theta \simeq 70°$ are inferred for the ejecta, corresponding to a Doppler factor $\delta \simeq 0.58$. A conservative estimate of the total energy released during this event yields $E_{43} = 20$ [1,4]. From Eq. (10) we then obtain $N_\mu \simeq 0.03\eta_{p,-1}$ for a 1 km^2 detector. At $\gtrsim$10 TeV this is above the background but would require an average of $\sim$30 outbursts for detection. We note that if a similar event were to occur in a source that is located closer to Earth, at a distance of say 3 kpc as in the case of the superluminal source GRO J1655-40 (which has a similar Doppler factor), or from a jet oriented at a smaller angle, then a few or even a single outburst may produce several muon events in such a detector. This suggests that even sources that are beamed away from us might be detectable during particularly strong outbursts.

As a second example consider the source SS433. This source, at a distance of 3 kpc, exhibits a steady jet that moves with a speed of $\sim$0.3 c. The presence of H_α lines indicates a baryonic content. A conservative estimate of the jet's kinetic power yields $L_j \gtrsim 10^{39}$ ergs s^{-1} [24]. Thus of order $10^3\eta_{p,-1}$ events per year per km^2 are anticipated. This rate is consistent with the MACRO upper limit [25]. Emission of TeV neutrinos from this source has been proposed earlier [26], although the mechanism invoked in this reference is pp collisions.

Positive detection of TeV neutrinos from microquasar jets would imply a baryonic content (although nondetection does necessarily imply a pair dominated jet) and would have important consequences for the mechanisms responsible for the ejection and confinement of jets. It would also provide important information concerning dissipation and particle acceleration in shocks. Contemporaneous detections of gamma rays may even enable us to infer the relative efficiencies at which electrons and protons are accelerated in shocks by comparing the total energies emitted as neutrinos and gamma rays during outbursts, thereby resolving a long-standing issue.

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