Astron. Astrophys. 344, 199–204 (1999)

ASTRONOMY
AND
ASTROPHYSICS

Optically thick super-Eddington winds in Galactic superluminal sources

A.A. Belyanin

Institute of Applied Physics, Russian Academy of Science 46 Ulyanov st., 603600 Nizhny Novgorod, Russia (belyanin@appl.sci-nnov.ru)

Received 28 September 1998 / Accepted 7 January 1999

Abstract. A model of major ejection events in Galactic superluminal sources GRO J 1655-40 and GRS 1915-105 is proposed, in which episodes of e^-e^+ pair creation give rise to the formation of optically thick outflows accelerating protons to the limiting Lorentz factor of order 2-3. Probable source of e^-e^+ pairs is pion production in nucleon collisions. Subsequent deceleration of ejecta and formation of radio jets is analyzed in detail within a shock deceleration model. The detailed data available for GRO J 1655-40 are shown to be in good agreement with proposed scenario.

Key words: acceleration of particles – accretion, accretion disks – radiation mechanisms: non-thermal – stars: individual: GRO J 1655-40 – stars: individual: GRS 1915+105

1. Introduction

Galactic microquasars are now among the most intriguing astrophysical sources. Even the qualitative aspects of their highly unusual behaviour still remain unexplained. The sources are believed to be powered by accretion in a binary system containing a massive stellar remnant, presumably a black hole. The radio counterpart of the sources during periods of their activity consists of a compact core and two subparsec-scale radio jets. Four objects of this morphology have been recently identified: 1E 1740.7–2945, GRS 1758–258, GRS 1915+105, GRO J 1655–40, and the last two of them demonstrate episodic ejection of radio blobs moving with an apparent superluminal speed (Hjellming & Rupen 1995, Mirabel & Rodríguez 1994). In this paper we put forward a scenario for the formation and acceleration of relativistic outflows in accreting black holes of stellar mass, apply it to the above two sources and formulate the observational predictions of the model.

To explain the acceleration to bulk relativistic velocities, the model of e^-e^+ jets has been proposed (Li & Liang 1996), in which pairs are accelerated by radiation pressure in the optically thin regime, fly away from the black hole to the distance of $10^{16} - 10^{17}$ cm and produce the observed radio emission. However, models of this type face serious problems (see also remarks in Levinson & Blandford (1996)). The resulting e^-e^+ outflow is expected to be very opaque in the inner part, and its dynamics will be different from that described in Li & Liang (1996). In particular, an efficient annihilation prevails over the expansion and leads to very rapid decrease in the pair density (see Eq. (6)). As a result, the number of pairs survived after the wind becomes optically thin is practically independent on the initial pair density (see Eqs. (6), (7) and the subsequent discussion, and also Guilbert & Stepney 1985, Ghisselini et al. 1992) and seems to be too small to account for the observed radio emission.

In this paper, we propose a new model of the sources, in which episodes of enhanced pair production result in optically thick super-Eddington winds. As a result of expansion, the energy of trapped radiation is converted into bulk kinetic energy of protons, until the limiting Lorentz factor of order 2–3 is reached. After the e^-e^+ pairs annihilate, the flow becomes optically thin and the energy is transported further away in the form of bulk kinetic energy of protons. The observed radio emission is produced by electrons accelerated in the relativistic shock moving ahead of the outflow. The peak flux is reached when ejecta begin decelerating in the interstellar medium. Such a scenario is similar in some aspects to the scaled-down fireball models of cosmological gamma-ray bursts. It has relatively low energy demands and provides a natural explanation to the observed lightcurve and spectrum of the radio emission.

2. Formation of relativistic outflows

First of all, let us consider whether a wind with required terminal velocity and power may be formed in the innermost part of an accretion disk during enhanced X-ray activity. Let L and $\dot{M}$ be the total energy input rate to the wind (excluding the rest energy) and the rest mass ejection rate, respectively. We assume that the outflow is stationary, spherically symmetric (or, more generally, is uniform over its opening angle $\Delta\Omega < 4\pi$), and adiabatic in its inner part (below the photosphere). In this case the dynamics of a single-fluid flow is governed by conservation laws of the energy-momentum and particle number fluxes (Flammang 1982, Paczyński 1990), conveniently written in Derishev et al. (1998):

$$\frac{p + \rho c^2}{n}\Gamma = \mathrm{const}, \quad (1)$$

Send offprint requests to: A. Belyanin

$$n\beta c\Gamma r^2 = \mathrm{const}, \tag{2}$$

where n is the number density of protons, p and ρc^2 are the pressure and the total energy density (including the rest energy nmc^2 and radiation energy); $\beta = v/c$ and $\Gamma = (1-\beta^2)^{-1/2}$ are the hydrodynamic velocity and Lorentz factor of the flow. The pressure satisfies adiabatic equation of state. Here and below all quantities characterizing the fluid are measured in the rest frame of a fluid.

The constant in Eq. (1) can be calculated at $r \to \infty$, where the stored thermal energy (the energy of trapped radiation and e^-e^+ pairs) has been transferred to the kinetic energy of protons, and the Lorentz factor reaches the limiting value $\Gamma_m = 1 + L/\dot{M}c^2$:

$$\frac{p + \rho c^2}{n}\Gamma = mc^2\Gamma_m. \tag{3}$$

If the flow becomes optically thin before the above limiting value is reached, the maximum Lorentz factor is $\Gamma_0 r_{\rm ph}/r_0$, where $r_{\rm ph}$ is the radius of photosphere, and we assume that the wind originates at the surface $r = r_0$ with initial Lorentz factor $\Gamma(r_0) = \Gamma_0 \simeq (3/2)^{1/2}$ corresponding to the velocity of sound.

The observed value of Γ_m for both sources is 2.5. This should be expected for winds with pair-dominated opacity. Indeed, the optical depth of the wind, calculated as an integral

$$\tau = \int_{r_0}^{\infty} \frac{n_e(r)\sigma_T}{\Gamma(r)}\, dr \tag{4}$$

should be much greater than unity. Here $n_e = n + 2n_+$ is comoving lepton density in the presence of e^-e^+ pairs of density $2n_+$. For $\dot{M} \lesssim 10^{18}$ g/s the electrons associated with protons cannot provide the necessary opacity unless the wind is highly collimated from the very beginning. Therefore, the condition $n_e \simeq 2n_+ \gg n$ is required at the base of the wind. If we assume that pair-production region is confined to the region $r < r_0$ and pairs can only annihilate at $r > r_0$, the comoving positron density n_+ at $r \gtrsim r_0$ obeys the continuity equation

$$\frac{1}{r^2}\frac{d}{dr}\left(n_+ r^2 c\beta\Gamma\right) \simeq -\frac{3}{8}\sigma_T n_+^2 c, \tag{5}$$

where the expression on the right-hand side is the annihilation rate of positrons in the nonrelativistic limit (Berestetskii et al. 1989).

Eq. (5) should be integrated with boundary conditions $n_+(r_0) = n_0$, $\Gamma(r_0) = \Gamma_0$. Since $\Gamma \simeq \Gamma_0 r/r_0$ at the acceleration stage (Paczyński 1990), we have (see also Guilbert & Stepney 1985)

$$n_+ = \frac{n_0}{\left(1 + \frac{\tau_0}{16\Gamma_0}\right)\left(\frac{r}{r_0}\right)^3 - \frac{\tau_0}{16\Gamma_0}}, \tag{6}$$

where $\tau_0 = 2\sigma_T n_0 r_0$.

Substituting Eq. (6) into Eq. (4) gives

$$\tau = \begin{cases} \tau_0/3\Gamma_0, & \tau_0 \ll 16\Gamma_0, \\ 16/3, & \tau_0 \gg 16\Gamma_0. \end{cases} \tag{7}$$

The maximum Lorentz factor achieved in a wind with opacity, dominated by e^-e^+ pairs, is therefore of order $(16/3)^{1/3} \sim 2$, which is a known result (Guilbert & Stepney 1985). This estimate is, however, *very* approximate. One source of errors is an approximation for the dependence $\Gamma(r)$ which is too crude for small Γ. It is relatively easy to find an exact dependence of the hydrodynamic velocity β from radius by solving Eqs. (1), (2) supplemented with adiabatic equation of state. This gives algebraic equation for β (Derishev et al. 1998) which should be solved together with (5) to give more accurate dependence $n_+(r)$.

However, the main source of uncertainty is related to the pair production. In a more realistic situation, pairs are also produced in the wind. Consider an inner accretion disk with proton density decreasing with height above the disk surface, surrounded by an extended region of pair production. The outflow starts from the height where the ratio n_+/n becomes sufficiently high, but the pairs can be also created in the wind in significant amount. In this case, pair density decreases with radius more slowly than in Eq. (6), and the limiting Lorentz factor increases.

Anyway, the above estimates illustrate that for large initial optical depth the limiting Lorentz factor is rather weakly dependent on the initial pair density. It is difficult to achieve $\Gamma_m \sim 10$, but is quite possible to obtain $\Gamma_m \sim$ 2–3.

To achieve large values of the initial optical depth τ_0 requires very high rate of energy injection, L. Indeed, the proton density at the base of the wind is given by

$$n \simeq \frac{L}{4\pi r_0^2 mc^3\beta_s\Gamma_m}. \tag{8}$$

Here $\beta_s \simeq 1/\sqrt{3}$ is the velocity of sound in a radiation-dominated fluid. We define the starting point of the wind, $r = r_0$, as the radius at which the hydrodynamic velocity is equal to the velocity of sound. Substituting (8) into the expression for the optical depth gives

$$\begin{aligned}\tau_0 = \left(\frac{2n_0}{n} + 1\right)\sigma_T n r_0 = \\ \left(1 + \frac{m_e}{m}\left(\frac{2n_0}{n} + 1\right)\right)\frac{\sqrt{3}}{2}\frac{r_g}{r_0}\frac{L}{\Gamma_m\tilde{L}_E},\end{aligned} \tag{9}$$

where r_g is Schwarzschild radius of the black hole, and we introduced the Eddington luminosity $\tilde{L}_E \simeq 1.2\cdot 10^{38}(m_e/m + n/(2n_0+n))M/M_\odot$ erg/s in the presence of e^-e^+ pairs.

It is reasonable to expect $r_0 \sim 3r_g$. Then, one can see from Eq. (9) that relativistic wind requires highly super-Eddington input luminosity $L \gtrsim 10\tilde{L}_E$.

For GRO J 1655-40 the black-hole mass $M \simeq 7M_\odot$ (Orosz & Bailyn 1997), while the X-ray luminosity $L_x \sim 10^{38}$ erg/s during outbursts (Harmon et al. 1995). The total input luminosity is unlikely to be much higher than L_x. This means that $n_+ > 10^2 n$ should be satisfied at the base of the wind to reach the required Lorentz factor. As follows from shock deceleration model presented below, the kinetic power of the wind $\Gamma_m\dot{M}c^2 \sim 10^{38}$ erg/s is required by radio observations. This gives the proton density $n(r_0) \simeq 3\cdot 10^{15}$ cm^{-3} and the corresponding positron density $n_0 > 3 \cdot 10^{17}$ cm^{-3} at the base of the wind.

For GRS 1915+105 the X-ray luminosity is one order of magnitude higher, which may result in a lower ratio of n_+/n and more powerful wind.

2.1. Pair production

The above estimates clarify to some extent the physical conditions in the source necessary to launch a relativistic wind. It remains "only" to explain how the e^-e^+ pairs are created.

First of all, it is clear that at the base of the wind the pair production is balanced by annihilation, and not by the expansion. Therefore, the required pair production rate is given by

$$q \simeq \frac{3}{8}\sigma_T c n_+^2 \simeq 8 \cdot 10^{21} \left(\frac{n_+}{10^{18}\,\mathrm{cm}^{-3}}\right)^2 \ \mathrm{cm}^{-3}\mathrm{s}^{-1}, \qquad (10)$$

which corresponds to the injected power in pairs of order

$$L_{\mathrm{pairs}} \sim q 2 m_e c^2 (3r_g)^3 \sim 10^{37} \left(\frac{n_+}{10^{18}\,\mathrm{cm}^{-3}}\right)^2 \ \mathrm{erg/s}, \qquad (11)$$

assuming $M = 10M_\odot$.

Now, let us specify the pair production mechanism, which is necessarily nonthermal for such low luminosities; most likely variants include electromagnetic cascade initiated by hard γ-quanta and relativistic electrons (positrons) in copious soft X-ray photon field produced by disk. The question is where the primary particles come from. A natural source of energetic leptons and γ-quanta in the inner disk, which we discuss here in some detail, is the decay of pions created in nucleon collisions, which is capable of producing γ-ray and positron luminosities reaching $10^{37} - 10^{38}$ erg/s (Kolykhalov & Sunyaev 1979, Jordain & Roques 1994). An attractive feature of this mechanism is the direct conversion of the thermal energy of protons into relativistic particles, without any ad hoc assumptions on particle acceleration etc.

There seem to be at least two possibilities of such a conversion. The first one is realized if the inner accretion flow is advection-dominated, and the proton temperature reaches nearly virial values. A reasonable estimate of the proton temperature which is valid both for adiabatic spherical accretion (Kolykhalov & Sunyaev 1979) and for models of advection-dominated disk accretion (Narayan & Yi 1995, Narayan et al. 1997) is

$$T \simeq T_0 r_g/r, \qquad (12)$$

where $T_0 \simeq 200$ MeV. As follows from results of Monte-Carlo simulations presented in Fig. 2 of Kolykhalov & Sunyaev (1979), for Maxwellian plasma with temperatures in the interval 100–200 MeV the total probability of pion creation can be approximated as $\eta \simeq \eta_0 (T/T_0)^{1/2}$, where $\eta_0 = 5 \cdot 10^{-16}$ cm^3/s.

Integrating the pion production rate, $q_\pi \simeq (1/2)\eta n_p^2$, where n_p is the proton density in an accretion flow, over the volume of a pion production region, $V \sim (3r_g)^3$, one can obtain the total luminosity in pions

$$L_\pi \simeq 3 \cdot 10^{37} \dot{M}_{18}^2 M_{10}^{-1} \ \mathrm{erg/s}, \qquad (13)$$

where $\dot{M}_{18}$ is mass accretion rate in units of 10^{18} g/s, $M_{10} = M/(10M_\odot)$. Here we assumed advection-dominated disk accretion model, applying the self-similar solution by Narayan & Yi (1995) which gives for the proton density

$$n_p \simeq 4 \cdot 10^{17} \alpha^{-1} \dot{M}_{18} M_{10}^{-2} (r_g/r)^{3/2} \ \mathrm{cm}^{-3}, \qquad (14)$$

where α is viscosity parameter (Shakura & Sunyaev 1973).

For spherical accretion we obtain nearly the same value of L_π. Approximately $1/7$ of this value goes to gamma-quanta of mean energy 70 MeV, and 1/4 goes to positrons of energy 35 MeV.

The second possibility, suggested by Derishev (1998), is related to the existence of neutron halo within several r_g from the black hole. Neutrons are efficiently created in the inner accretion flow due to collisions of protons with He nuclei when the proton temperature is only about several MeV. Because of rather weak coupling with protons, neutrons occupy elliptic orbits, forming a halo around the central object. Let us find their equilibrium density. We will assume that He density in the accretion flow is $n_{\mathrm{He}} \simeq 0.1 n_p$, and take here and below in this section the expression (14) for proton density, and expression (12) for the ion temperature in the advection-dominated accretion flow.

The energy of He dissociation is about 28 MeV, which becomes comparable to the thermal energy $m v_t^2/2 = 1.5T$ of protons at $r \sim 10r_g$. Starting from this radius, the lifetime of He nuclei with respect to dissociation is $t_{\mathrm{dis}} \simeq (\sigma_{\mathrm{dis}} v_t n_p)^{-1}$. This is to be compared with accretion time $t_a \simeq r/v_r$, where the radial velocity $v_r \sim \alpha v_t$ (the neutron decay is unimportant due to very long decay time, 10 min). I have not found any data on the cross section of He dissociation by protons at relevant energies, but from purely geometrical reasons (calculating the de Broglie wavelength of nucleons with energy 7 MeV) it should be within 40–60 mb. In addition, partial dissociation of He also contributes to the neutron release, increasing the total cross section. The ratio

$$t_{\mathrm{dis}}/t_a \simeq 20\alpha^2 M_{10} \dot{M}_{18}^{-1} (40\,\mathrm{mb}/\sigma_{\mathrm{dis}}) (r/r_g)^{1/2} \qquad (15)$$

could become smaller than unity for $\alpha \lesssim 0.1$ in the innermost disk. Therefore, *if* the proton temperature in the innermost disk is sufficiently high, He nuclei can be completely destroyed, and the neutron density can be as high as $n_n \simeq 2n_{\mathrm{He}} \simeq 0.2n_p$.

In fact, it can be even higher if we take into account that neutrons accrete slower than ions. In this case the presence of neutrons may affect both collision processes and the dynamics of accretion, and more careful consideration is needed; the detailed discussion will be published elsewhere (Derishev & Belyanin 1999).

The collisions of neutrons on elliptical orbits with nucleons can occur at high relative velocities $v/c \sim (r_g/r)^{1/2}$ and lead to the pion production. Assuming $n_n \gtrsim 0.2n_p$ in the volume $(3r_g)^3$, typical luminosity in pions turns out to be of the same order or even higher than the above estimate for *pp* collisions. The advantages of this mechanism of pair production are lower proton temperatures and possibility to inject e^-e^+ pairs directly to the wind, high above the disk. The latter is especially important for producing a wind with small proton "pollution".

However, the above mechanism operates only if the transverse size of the disk is not much greater than the mean free path of neutrons with respect to np collisions. Otherwise, neutrons will be captured by the accretion flow and advected to the black hole. On the other hand, too low proton density can lead to low neutron densities and low value of pion production rate. The optimal case is when

$$\sigma_{\rm np} n_p h \lesssim 1, \tag{16}$$

which gives

$$n_p \lesssim 3 \cdot 10^{18} \frac{h}{r_g} M_{10} \ {\rm cm}^{-3}. \tag{17}$$

Here $\sigma_{\rm np} \simeq 100$ mb, h is the disk thickness. The value (17) is remarkably close to typical proton densities in the innermost accretion disk, predicted by disk models for mass accretion rates $\dot{M}_a \sim 10^{18}$ g/s; see Eq. (14).

2.2. Gamma-ray signatures

Due to efficient pair production on X-ray photons from disk and wind photosphere, pion-produced γ-quanta are completely absorbed in the source, initiating an electromagnetic cascade. Simulations by Jourdain & Roques (1994) predict power-law spectral tail extending to tens of MeV. Such a tail was observed by COMPTEL in both sources during outbursts; it disappeared in the quiescence periods (Iyudin et al. 1998). Of course, this is not a unique feature of pion production, but rather a signature of ultrarelativistic particles.

The spectral feature that may have the highest probability to be observed is a broad peak around 0.5 MeV resulting from positron annihilation in the accretion disk. Since we observe the disk nearly edge-on for both sources, the peak will be blueshifted and, in principle, can be spread over the whole range of frequencies between 0.43–0.7 MeV, where we took into account gravitational redshift and Doppler factor at the last stable orbit around the nonrotating black hole. Its intensity is, however, dependent on many uncertain factors. One of them is the fraction of positrons which is slowed down and annihilates in the disk. In the most optimistic case, when this fraction is of order 0.5, the resulting annihilation luminosity can be of the order of the total luminosity in pairs, which should be $\sim 10^{37}$ erg/s in our model. However, there always exists a possibility that annihilation region in the inner disk is completely hidden from view by outer parts of the disk or by optically thick outflow.

Positrons that survived in the wind after it becomes optically thin with respect to annihilation are captured by the magnetic field in the interstellar medium and decelerated, producing narrow and unshifted 0.511 MeV line. The total flux of positrons is practically independent on their initial density and is equal by order of magnitude to $4\pi r_{\rm ph}\Gamma_m^2 c/\sigma_T \sim 3\cdot 10^{42} M_{10}(r_{\rm ph}/r_g)\,{\rm s}^{-1}$. However, the deceleration time of positrons is 10^5 years if the density of an ambient medium is $1\,{\rm cm}^{-3}$. Therefore, the luminosity in the narrow 0.511 MeV line depends on the duty cycle of the source during previous epochs. If the duty cycle of major positron ejections exceeded 10^{-4}, the annihilation line can be detected by INTEGRAL.

3. Deceleration of ejecta and formation of radio jets

Suppose that the ejection starts at the moment $t = 0$ and continues during the time T. As follows from radio observations, the outflow eventually becomes highly collimated to an opening angle $\phi = 10^{-1}\phi_{-1}$ rad. The mechanism of collimation is beyond the scope of this paper; see Levinson & Blandford (1996) for the recent discussion.

As interstellar plasma is swept up by the expanding wind, an external shock is formed, which propagates into an ambient medium with Lorentz factor $\Gamma_s = \sqrt{2}\Gamma_m$ (Blandford & McKee 1976). When the pressure of swept-up particles becomes comparable to the wind pressure, ejecta begin decelerating, and Lorentz factor of the shock decreases according to the energy conservation law (Blandford & McKee 1976)

$$P_{\rm kin} \sim \phi^2 r^2 n_i m c^3 \beta^2 \Gamma^4. \tag{18}$$

Here $P_{\rm kin} \lesssim L$ is kinetic power of bulk motion of ejecta, $n_i = 10^{-1} n_{-1}\,{\rm cm}^{-3}$ is the density of the ionized fraction of an ambient medium, assumed to be uniform. The conservation law (18) follows from equating in the shock frame the injected momentum flux to the momentum flux of the unshocked gas.

According to Eq. (18), deceleration begins at the lab-frame radius r_d given by

$$\begin{aligned} r_d \ &\sim \ \left(\frac{P_{\rm kin}}{n_i m c^3 \beta_m^2 \Gamma_m^4 \phi^2}\right)^{1/2} \\ &\sim \ 8\times 10^{15} (P_{38}/n_{-1})^{1/2} \phi_{-1}^{-1} (\Gamma_m/2.5)^{-2} \ {\rm cm}. \end{aligned} \tag{19}$$

Using (19), (18), we obtain

$$\frac{\beta\Gamma^2}{\beta_m \Gamma_m^2} \sim \frac{r_d}{r}. \tag{20}$$

After the radius $r_{\rm nr} \sim r_d \Gamma_m^2$ the shock becomes nonrelativistic and propagates according to the law

$$r \sim (P_{\rm kin}/n\phi^2)^{1/5} t^{3/5}. \tag{21}$$

Eqs. (19)-(21) are valid in the case of a constant energy supply to the shock, $P_{\rm kin} =$ const. However, the ejection events last for the finite time T, and for the sources in question the inequality $cT \ll r_d \Gamma_m^2$ seems to be fulfilled. This means that the last portion of ejecta catches up the shock before it decelerates to nonrelativistic velocities. The radius r_c at which this occurs can be determined from equation of motion for the decelerating shock. For $\Gamma_m = 2.5$ and $cT \ll 2 r_d \Gamma_m^2$ we obtain $r_c \simeq 2.4 r_d$. After this radius the deceleration proceeds much faster, according to the energy conservation law for the case of impulsive energy release:

$$\frac{\beta\Gamma}{\beta_m \Gamma_m} \sim \left(\frac{\tilde{r}_d}{r}\right)^{3/2}, \tag{22}$$

where

$$\frac{\tilde{r}_d}{\mathrm{cm}} \sim \left(\frac{P_{\mathrm{kin}} T}{n_i m c^2 \beta_m^2 \Gamma_m^2 \phi^2}\right)^{1/3} \sim 2\times 10^{16}\left(\frac{P_{38} T_w}{n_{-1}}\right)^{1/3} (\beta_m \phi_{-1})^{-2/3}\left(\frac{\Gamma_m}{2.5}\right)^{-2/3}, \tag{23}$$

$T_w = T/1$ week. In the nonrelativistic limit the motion of ejecta approaches the Sedov-Taylor solution $v \propto r^{-3/2}$ and $r \propto t^{2/5}$.

An expanding shock accelerates relativistic electrons that produce synchrotron radiation in a turbulent magnetic field behind the shock. Assuming that fractions η_e and η_B of the thermal kinetic energy of the shocked gas are carried by electrons and magnetic field, we arrive at the expression for the characteristic Lorentz factor of relativistic electrons and the magnetic field behind the shock:

$$\gamma_e \sim \eta_e \beta^2 \Gamma m/m_e, \tag{24}$$

$$B \sim 0.1(\eta_B n_{-1})^{1/2} \beta\Gamma\ G. \tag{25}$$

For electrons with Lorentz factor (24) in a magnetic field (25) the gyroradius is much smaller than the thickness of a shocked gas shell, and the characteristic synchrotron frequency is

$$\nu_{\mathrm{syn}} \simeq 10^6 B\gamma_e^2 \simeq 5\cdot 10^{11} \eta_B^{1/2} \eta_e^2 \beta^5 \Gamma^3 n_{-1}^{1/2}\ \mathrm{Hz}, \tag{26}$$

which falls into GHz range for $\eta_e \sim \eta_B \sim 0.1$. Note the rapid decrease of ν_{syn} after the beginning of deceleration.

It is easy to see that the time of synchrotron radiation losses, $t_{\mathrm{syn}}(\mathrm{sec}) \sim 5\cdot 10^8/\gamma_e B^2 \sim 2.5\cdot 10^7/(\eta_e \eta_B \beta^4 \Gamma^3)$, is much greater than the characteristic expansion timescales $t_d = r_d/c$ and $t_{\mathrm{nr}} = r_{\mathrm{nr}}/c$, which justifies adiabatic approximation.

The total amount of relativistic electrons behind the shock in the comoving frame scales with expansion as

$$N \sim 10^{44} \phi_{-1}^2 r_{16}^3 n_{-1}, \tag{27}$$

where $r_{16} = r/10^{16}$ cm.

The total comoving synchrotron luminosity of a shocked gas is

$$L_{\mathrm{syn}} \sim 4\cdot 10^{32} \phi_{-1}^2 r_{16}^3 \eta_e^2 \eta_B n_{-1} \beta^6 \Gamma^4\ \mathrm{erg/s}, \tag{28}$$

where it is assumed that all electrons have Lorentz factor γ_e. For $r < r_d$ the luminosity increases with time as t^3; after reaching r_d it grows more slowly due to decrease in the magnetic field and electron energy; after reaching $r = r_c$ it falls down rapidly. Of course, in reality the spectrum of electrons is not a delta-function peaked at equipartition value, but is determined by a competition between acceleration rate, synchrotron losses and escape from the acceleration region. However, the spectral flux at a given frequency has a similar behaviour which can be easily found for given spectral distribution of electrons derived from radio spectrum.

4. Application to GRO J 1655-40

Let us apply the above model to correlated X-ray and radio outbursts in August-September 1994, for which the detailed data are available (Hjellming & Rupen 1995, Harmon et al. 1995). The jets were oriented nearly perpendicular to the line of sight, with inclination angle $i \simeq 85°$ and the Lorentz factor of bulk motion $\Gamma_m \simeq 2.5$. The first X-ray outburst occurred between the days 9561-9581 (the truncated Julian day 9561 corresponds to 28 July 1994), with total X-ray luminosity reaching 10^{38} erg/s, as inferred from BATSE observations (Harmon et al. 1995).

VLA observations began on day 9576, revealing an unresolved source with a flat spectrum. The peak flux of 5.5 Jy at 1.49 GHz was reached six days later. The spectral index was equal to $\alpha \simeq 0.65$ (assuming the flux $S_\nu \propto \nu^{-\alpha}$). After the day 9582 the flux falls down in all frequencies, from 1.49 to 22.5 GHz, with timescale of the order of one week and the same spectral index.

According to Fig. 4 of Hjellming & Rupen (1995), the corresponding ejection event started on day 9574 ± 1 and apparently lasted for several days. We will assume that ejection ended on day 9581, simultaneously with the end of an X-ray outburst, so that $T \sim 1$ week. Let us identify the moment of peak flux (day 9582) with the time when ejecta reached the radius r_c which marked the onset of a rapid deceleration stage described by Eqs. (22), (23). For observed angular velocity of ejecta this corresponds to a separation of 300 mas, or $1.5\cdot 10^{16}$ cm from the core for a distance to the source equal to 3.2 kpc. Then the radius r_d is equal to $r_c/2.4 \simeq 6.3\cdot 10^{15}$ cm, and the radius where the shock is decelerated to nonrelativistic velocities is given by $r_{\mathrm{nr}} \sim r_d \Gamma_m^2 \sim 4.1\cdot 10^{16}$ cm. This well corresponds to the observed distance from the core at which ejecta became noticeably decelerated (Hjellming & Rupen 1995).

Thus, the observed motion of ejecta appears to be in agreement with shock deceleration model. The same is true for radio flux magnitude. Indeed, let us assume that electrons in the shocked gas have power-law spectrum $n_e(\gamma) = K\gamma^{-p}$ with $p = 2\alpha + 1 \simeq 2.3$ in the interval $\gamma > \gamma_{\mathrm{min}}$. Then, using standard synchrotron formulas and allowing for relativistic corrections, we obtain the total number of electrons required to produce the peak value of the radio flux:

$$N(\gamma > \gamma_{\mathrm{min}}) \simeq 2\cdot 10^{46} \gamma_{\mathrm{min}}^{-1.3}\left(\frac{\Gamma}{B}\right)^{1.65} \times \frac{S_\nu(1.49\,\mathrm{GHz})}{5.5\,\mathrm{Jy}}. \tag{29}$$

Choosing $\gamma_{\mathrm{min}} \sim \gamma_e$ from Eq. (24) and taking the value of magnetic field from Eq. (25), we obtain a good agreement with an estimation (27) of the total number of electrons behind the shock, calculated for the shock radius $r = r_c \simeq 1.5\cdot 10^{16}$ cm.

To calculate the evolution of the spectral flux, we assume that the total number of relativistic electrons increases as in Eq. (27) during all stages of expansion, the electrons have power-law spectrum with the constant index p, and the lower cutoff energy of the spectrum, γ_{min}, decreases as γ_e in Eq. (24). The resulting behaviour is qualitatively similar to that of the total luminosity (28): the spectral flux at all observation frequencies increases with time before reaching r_c, and drops with characteristic timescale of order $\tilde{r}_d/c \sim 1$ week after this radius.

5. Conclusions

In conclusion, we have shown that relativistic jets observed in Galactic superluminal sources may result from the formation of optically thick super-Eddington outflows accelerating protons to relativistic velocities. Radio emission is produced in a strong relativistic shock, which is formed when protons begin decelerating in an interstellar medium. Kinematics of ejecta and properties of the radio emission, observed in GRO J 1655-40, are shown to be in a good agreement with a shock deceleration model.

Acknowledgements. The author is grateful to E. V. Derishev and Vl. V. Kocharovsky for fruitful discussions, and to the referee, Igor Moskalenko, for many helpful comments. This work has been supported by the Russian Foundation for Basic Research.

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