The bulk kinetic power of the jets of GRS 1915+105

M. Gliozzi,1 G. Bodo2 and G. Ghisellini3
1Dipartimento di Fisica Generale di Torino, Via P.Giuria, I-10125 Torino, Italy
2Osservatorio Astronomico di Torino, Strada Osservatorio 20, I-10025 Pino Torinese, Italy
3Osservatorio Astronomico di Brera, Via Bianchi, 46, I-23807 Merate, Italy

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ABSTRACT
We calculate the minimum value of the power in kinetic bulk motion of the galactic superluminal source GRS 1915+105. This value far exceeds the Eddington luminosity for accretion onto a black hole of mass 10 M\(_\odot\). This large value severely limits the possible carriers of the kinetic luminosity at the base of the jet, and favours a jet production and acceleration controlled by a magnetic field, the value of which, at the base of the jet, exceeds 10\(^8\) G. The Blandford & Znajek process can be responsible for the extraction of the rotational energy of a Kerr black hole, if it lasts long enough to provide the required kinetic energy. The time required, of the order of a day, implies that the process must operate in a stationary, not impulsive, mode.

Key words: radiative mechanisms: non-thermal – stars: individual: GRS 1915+105.

1 INTRODUCTION
GRS 1915+105 was discovered in 1992 with the WATCH telescope on board the GRANAT satellite (Castro-Tirado, Brandt & Lund 1992). A radio counterpart was subsequently identified and bipolar outflows with apparent superluminal motions were observed; the standard interpretation of this phenomenon in terms of relativistic jets (Rees 1966) places the source at a distance \(D = 12.5\) kpc at an angle \(i = 70^\circ\) to the line of sight (Mirabel & Rodriguez 1994). The source cannot be observed in the optical band, because of the heavy extinction in the Galactic plane, but the X-ray luminosity, often well above the Eddington luminosity for a neutron star, suggests the presence of a \(\sim 10\)-M\(_\odot\) black hole. Other circumstantial evidence comes with the similarities with the other galactic superluminal source GRO J1655−40, which has been shown unambiguously to harbour a compact object of \(\sim 7\) M\(_\odot\) (Orosz & Bailyn 1997). Since its discovery, GRS 1915+105 has displayed an extraordinary richness in variability in the X-ray (Greiner, Morgan & Remillard 1996; Belloni et al. 1997; Chen, Swank & Taam 1997), infrared and radio bands (Fender et al. 1997; Pooley & Fender 1997). Recent (Eikenberry et al. 1997; Mirabel et al. 1998) simultaneous multiwavelength observations of this superluminal source show a strict link between activity in the inner accretion disc and plasma ejections.

In this paper we estimate the minimum kinetic power associated with the ejections and investigate the possible energy transport mechanism from the compact source to distances where the radio blobs are resolved, showing that the more suitable solution implies the presence of a strong magnetic field at the footpoint of the jet.

In Section 2 we summarize the existing radio and infrared data for the ejection events. In Section 3 we directly calculate the minimum kinetic power during the ejection. In Section 4 we investigate the possible origin of the kinetic power. In Section 5 we summarize the main results of the paper and discuss their consequences.

2 OBSERVATIONAL CONSTRAINTS DURING THE OUTFLOWS
Radio observations show that GRS 1915+105 in its active state is characterized by the increase of the flux level from \(S_v \leq 10\) mJy to \(S_v \sim 100\) mJy, the so called plateau state, on top of which radio flares are superimposed, with durations from days to a month and fluxes up to a few Jy (Foster et al. 1996). The spectra during the plateau state are generally flat or inverted, \(\alpha \leq 0\) (with \(S_v \propto \nu^{-\alpha}\)), probably indicating synchrotron self-absorption. On the other hand, the spectra of radio flares reveal a transition from optically thick to thin (\(\alpha > 0\)) emission during the rise stage, with spectra harder (\(\alpha \sim 0.5\)) when the fluxes reach the maximum and softer (\(\alpha \sim 1\)) during the subsequent decay, generally interpreted as an indicator of synchrotron radiation from an expanding radio cloud.

Observations with the Very Large Array and multiwavelength campaigns (with VLA and UKIRT simultaneously) pointed out that some strong radio flares are related to the ejection of radio clouds. There are two different kind of ejections:

(i) major ejections, such as the prominent radio outburst of 1994 March (Mirabel & Rodriguez 1994), where the ejecta actually moved for several weeks along a direction forming an angle \(\theta = 70^\circ\) to the line of sight, with bulk velocity 0.92c and expansion at \(\sim 0.2c\); with a spectral index change from \(\alpha = 0.49\) (when the blobs could not be resolved) to \(\alpha = 0.84\) (when the two condensations had moved apart), and

(ii) so-called ‘baby jets’ (Eikenberry et al. 1997; Mirabel et al. 1998).
3 MINIMUM KINETIC POWER CONDITION

In order to deduce the minimum kinetic luminosity $L_k$ related to a major ejection we consider the 1994 March event, for which the observational data are the most exhaustive. As we are interested more in power than in energy estimates, we follow an alternative method rather than the standard one, which consists of determining first the internal energy of the blob via the minimum energy criterion and then obtaining an estimate of the kinetic power, by making some (highly uncertain) assumptions about the energization time.

We instead directly calculate the kinetic power, equal to the energy flux through a cross-section of the jet (see also Ghisellini & Celotti 1999). This energy flux can be carried by particles and by the toroidal magnetic field, and correspondingly we have

$$L_{k,p} = \pi R^2 \Gamma^2 \beta c n_e (\gamma) m_e c^2 \left( 1 + \frac{m_e}{m_i (\gamma)} \right),$$

$$L_{k,B} = \pi R^2 \Gamma^2 \beta c B_u,$$

where $m_i$ is either $m_e$ in the case of a ‘normal’ e–p plasma or $(\gamma)m_i$ for $e^\pm$ pairs, $\Gamma$ is the Lorentz factor of the bulk motion, $R$ is the cross-section radius of the jet, $n'$ is the comoving particle density, $B_u$ is the magnetic energy density measured in the comoving frame and $(\gamma)$ is the mean Lorentz factor of the electrons, as measured in the comoving frame. A lower limit on $n'$ can be estimated from the observed synchrotron emission $L_{\text{syn}}$ of the blob. Assuming a spherical emitting volume of radius $R$, the number density of leptons producing the observed radiation is

$$n' = \frac{9L_{\text{syn}}}{2(\gamma^2 \beta^2 c^5 B^2 R^2)},$$

where $\langle \gamma \rangle$ is averaged over the relativistic electron distribution, and $\delta = (1 - \beta \cos \theta)^{-1}$ is the Doppler factor. The $R^{-2}$ dependence of the estimated particle density allows us to minimize the total power $L_{k,p} + L_{k,B}$ with respect to the magnetic field, because both the bulk Lorentz factor $\Gamma$ (\approx 2.55) and the viewing angle $\theta$ (\approx 72\degree) are known:

$$\frac{\partial}{\partial B} (L_{k,p} + L_{k,B}) = 0,$$

yielding the value of the magnetic field $B_{\text{min}}$ corresponding to the minimum power,

$$B_{\text{min}} = \frac{[36\pi L_{\text{syn}} m_e c^2]}{\Gamma^4 \delta^5 c^3},$$

where $f$ is a parameter directly related to the shape of the electron energy distribution, and therefore to the shape of the synchrotron radiation they emit:

$$f = \frac{\langle \gamma \rangle}{(\langle \gamma \rangle)^2} \left( 1 + \frac{m_e}{m_i (\gamma)} \right).$$

As the radio emission is a power law with spectral index $\alpha \sim 0.5$, the particle distribution $N(\gamma) \propto \gamma^{-\alpha}$ between $\gamma_1$ and $\gamma_2$, yielding $f = \ln(\gamma_2/\gamma_1)\gamma_2/2 [1 + m_e/m_i (\gamma_1 \ln(\gamma_2/\gamma_1)]$. With $R \approx 7 \times 10^{15}$ cm (geometric average of the size of blob of GRS 1915+105, assuming a distance of 12.5 kpc), $L_{\text{syn}} = 10^{35}$ erg s$^{-1}$ and $\beta = 0.92$ (corresponding to $\Gamma = 2.55$), we obtain $B_{\text{min}, \text{e-p}} = 0.12$ G and $B_{\text{min}, e^\pm} = 0.036$ G if $\gamma_1 = 1$ and $\gamma_2 = 10^3$, as required for production of the observed synchrotron photons at $\nu \approx 300$ GHz. A low-energy cut-off in the particle distribution would decrease $B_{\text{min}}$ somewhat ($B_{\text{min}, \text{e-p}} = 0.054$ G and $B_{\text{min}, e^\pm} = 0.031$ G for $\gamma_1 = 30$), while the high-energy cut-off is consistent with the production of the observed radiation at 300 GHz (with $B = B_{\text{min}}$).

However, the dependence of $B_{\text{min}}$ on the extremes of the electron distribution is rather weak ($B_{\text{min}} \propto f^{1/4}$).

Note that the value of $B_{\text{min}, \text{e-p}} = 0.12$ is almost a factor of 3 greater than the one obtained by Mirabel & Rodriguez (1995) and Liang & Li (1995), while it more-or-less agrees with the value estimated from the requirement of optical transparency of the plasmoid with respect to synchrotron self-absorption (Atoyan & Aharonian 1999).

With a magnetic field value equal to $B_{\text{min}}$, the particle kinetic power and the Poynting flux are nearly equal, and the total power,

$$L_{k,\text{tot}} = 3 \Gamma^2 \beta c \left[ \frac{\pi m_e^2 c^2 R_{\text{syn}} f}{\sigma_T c} \right]^{1/2},$$

is, respectively for e–p plasma and $e^\pm$ pairs, $\sim 3.3 \times 10^{40}$ erg s$^{-1}$ and $\sim 2.9 \times 10^{39}$ erg s$^{-1}$.

This is the minimum kinetic luminosity involved in major ejection events, calculated assuming $\gamma_1 = 1$ and $\gamma_2 = 10^3$ ($L_{k,\text{tot}, \text{e-p}} \sim 6.4 \times 10^{39}$ erg s$^{-1}$ and $L_{k,\text{tot}, e^\pm} \sim 2.1 \times 10^{39}$ erg s$^{-1}$ if $\gamma_1 = 30$).

Setting $B = B_{\text{min}}$ in equation (2), the total numbers of emitting particles in the magnetic cloud are respectively $N_e' = 5.9 \times 10^{47}$ and $N_{e^\pm} = 6.9 \times 10^{47}$, corresponding to internal energies

$$E_i = m_e c^2 \left[ \pi N(\gamma) (\gamma) d\gamma = N(\gamma) m_e c^2 \right.$$

$$\sim 3.3 \times 10^{42} \text{ erg} \quad \sim 3.9 \times 10^{43} \text{ erg},$$

These internal energies can be compared with the corresponding kinetic energies in bulk motion $E_k$ given by

$$E_k = L_k t = \frac{L_k}{R/c},$$

$$\sim 8.5 \times 10^{45} \text{ erg} \quad \text{and} \quad \sim 7.4 \times 10^{46} \text{ erg},$$

where $t$ is the time needed for the blob to cross the jet section.

4 ORIGIN OF THE KINETIC LUMINOSITY

The kinetic power calculated above refers to the radio-emitting blob, at a huge distance from the putative black hole and accretion disc. The most economic assumption is that this power is approximately conserved along the jet, with very small dissipation giving rise, e.g., to the random particle energy responsible for the emission. If not, we would have to assume that at the base of the jet an even larger value of $L_k$ is produced. In principle there are several possible energy carriers: $e^\pm$ pairs, $\text{p–e}^\pm$, pure magnetic field or a mixture of these components. Consider also that the particles can in principle be ‘hot’ or ‘cold’: at large distances at least some of the electrons are ‘hot’ (indeed, in our estimates we assumed that all the
electronics are ‘hot’, because this is the most economic assumption), however these ‘hot’ particles cannot come directly from the inner region (see below), but need to be accelerated or reaccelerated, therefore for the inner jet energy carriers one has to consider both cases. Furthermore, we note that in the ‘hot’ case, the particle random energy, increasing the mass, can contribute to the kinetic power. Let us examine the possible cases.

4.1 ‘Cold’ $e^\pm$ pairs

If the kinetic power is carried by a pure $e^\pm$ cold pair plasma, we can calculate the corresponding pair density and scattering optical depth at some jet radius $R$ close to the base of the jet:

$$\tau_\pm = \frac{\sigma_T L_k}{\pi R^2 \beta m c^2} \approx 4.2 \times 10^2 \frac{L_k}{2.9 \times 10^{39} \text{ erg s}^{-1}} \frac{10^7 \text{ cm}}{R}.$$  \hspace{0.5cm} (9)

As the annihilation time-scale is of the order of $R/(c\tau_\pm)$, cold pairs cannot survive annihilation.

4.2 ‘Hot’ $e^\pm$ pairs

If the pairs are relativistic, with average random Lorentz factor $\langle \gamma \rangle$, the above estimate of $\tau_\pm$ decreases by a factor $\langle \gamma \rangle$. In addition, the annihilation cross-section decreases. However these relativistic pairs are embedded in a dense radiation field, produced by the accretion disc, and they cool on time-scales shorter than the dynamical time-scale. In a region of size $R_{d}$, the radiation energy density resulting from the accretion disc luminosity $L_d$ produced within $R_{d}$ is of the order of $U_d = L_d/(4\pi R_d^2 c)$. The ratio between the inverse Compton cooling time and the dynamical time $R_d/c$ for a particle of Lorentz factor $\gamma$ is

$$t_{\text{IC}} \approx \frac{4\pi R_d m c^3}{\sigma_T L_d \gamma} \ll 1.$$  \hspace{0.5cm} (10)

A large fraction of $L_k$ would be lost and converted into radiation, contrary to what is observed. We therefore conclude that a pure $e^\pm$ pair plasma cannot carry the derived jet kinetic power, irrespective of whether it is hot or cold.

4.3 ‘Normal’ $e^-p$ plasma

The case of electrons with an average random Lorentz factor $\langle \gamma \rangle$ greater than $m_e/m_p$ and cold protons is analogous to the above case: inverse Compton scattering of seed accretion disc photons cools the electrons in a time shorter than $R_d/c$.

The case of ‘hot’, relativistic protons is instead immune to radiative losses, and our main concern is with the required confining pressure. If the kinetic power is carried by protons with an average Lorentz factor $\langle \gamma_p \rangle >$, the corresponding (comoving) pressure is of the order of their energy density. If the confinement is magnetic, the required value of the magnetic field is

$$B = \left( \frac{8L_k}{R^2 \Gamma^2 \beta c} \right)^{1/2},$$  \hspace{0.5cm} (11)

which is equal to the value that the magnetic field must have if it is the main carrier of the kinetic power.

If both electrons and protons are cold, the electron-scattering optical depth is a factor $m_e/m_p$ smaller than that of equation (9), i.e. of the order of a few. There are no severe limits in this case: the Compton drag is not sufficient to slow down the plasma in the jet. It is, however, useful to estimate the predicted power emitted by the Compton drag process, because it can turn out to be important in sources where the jet is more aligned with the line of sight. For this simple estimate we will assume that the jet is Compton-thick (i.e. $\tau_\gamma > 1$), so that the effective cross-section of the process is the geometrical one, i.e. $\pi R^2$. As before, we assume that the accretion disc produces the luminosity $L_d$ within the region $R_d > R$, at the typical frequency $\nu_d$. With these hypotheses, the observed bulk Compton luminosity $L_{\text{dc}}$ is

$$L_{\text{dc}} \approx \left( \frac{R}{R_d} \right)^2 \frac{L_d}{4} \nu^2 \delta^4.$$  \hspace{0.5cm} (12)

Most of this luminosity is observed at the frequency $\nu \approx \delta \nu_d$. With $R/R_d \approx 1/10$, $L_d \approx 10^{39}$ erg s$^{-1}$ and $\delta = 0.57$ we obtain $L_{\text{dc}} \approx 1.7 \times 10^{36}$ erg s$^{-1}$ and $\nu \approx \nu_d$. This radiation is therefore unobservable in the case of the known Galactic superluminals, which are both observed with large viewing angles, but may be very important for sources with $\delta \approx \Gamma > 1$ (see also Sikora et al. 1997).

4.4 Poynting vector

The value of the magnetic field needed to carry $L_k$ in the vicinity of the black hole is equal to that given in equation (11), if the bulk of the magnetic field is moving with $\Gamma$, while the $B$ value before acceleration is a factor of $\Gamma$ greater, corresponding to $B \approx 3 \times 10^8$ G at $R = 10^7$ cm.

It must be remarked that the standard theory of accretion discs (Shakura & Sunyaev 1973) predicts that the maximum possible magnetic field at a given $R/R_{d}$ (where $R_d$ is the Schwarzschild radius) in a radiation-dominated disc is $B \propto (M/M_{\odot})^{-1/2}$, so that microquasars will have magnetic fields $\sim 10^7$ times stronger than those in quasars. Therefore, the theoretical estimate $B \approx 10^{-10} - 10^9$ G obtained with the Blandford & Payne model (Blandford & Payne 1982) near the massive black holes in active galactic nuclei (AGN) clearly shows that in microquasars $B \sim 10^9$ G can easily be attained.

Note that this value of the magnetic field would be of the same magnitude as the magnetic field required by the Blandford–Znajek (Blandford & Znajek 1977) process to produce $L_k$ by extracting the rotational energy of a 10$M_{\odot}$ Kerr black hole:

$$L_{\text{rot}} = 10^{41} B_0^2 M_1^2 \text{ erg s}^{-1},$$  \hspace{0.5cm} (13)

where $B = 10^9 B_0$ G and $M = 10 M_1 M_{\odot}$.

From the above arguments, we conclude that $e^\pm$ or relativistic electrons with $\langle \gamma \rangle$ greater than $m_e/m_p$ cannot carry $L_k$ in the inner region of the jet, while protons and the Poynting vector can.

An alternative solution could be that the energy that is carried in the inner region by a ‘normal’ cold plasma or toroidal magnetic field is converted into $e^\pm$ pairs at a large distance from the black hole and accretion disc. Pairs are produced most efficiently through the inverse Compton drag until the energy density resulting from the accretion disc luminosity $L_d$ produced within $R_d$ is of the order of $U_d = L_d/(4\pi R_d^2 c)$. The ratio between the pair plasma cannot carry the derived jet kinetic power, irrespective of whether it is hot or cold.

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initially carried by a large magnetic field there is the possibility of tapping a great reservoir of energy (the rotational energy of the black hole) and accelerating the plasma to relativistic speeds.

We therefore conclude that the kinetic power carried initially by the magnetic field is the most economic way to explain the observed energetics. Differential accretion disc rotation or rotating black holes can amplify magnetic fields to the required values, and the magnetic field will then accelerate particles to relativistic speeds.

The amount of matter present at the base of the jet may not be negligible, especially if magnetic field lines help to channel particles from the disc to the jet. We can compare the inflow rate \( M_{\text{in}} \) of the accretion process with the outflow rate \( M_{\text{out}} \) necessary to account for the kinetic power, in particles, of the radio blob. If we write the accretion luminosity as \( L_{\text{acc}} = \eta M_{\text{in}} c^2 \), as usual, and the kinetic power as \( L_k = (\Gamma - 1) M_{\text{out}} c^2 \), we derive

\[
M_{\text{out}} = \frac{\eta}{\Gamma - 1} L_k M_{\text{in}}.
\]

As \( \eta \sim 0.1 \) and \( L_k \sim 10 L_{\text{acc}} \), we have that the inflowing and outflowing mass rates are comparable. This in turn suggests that most of the matter in the jet may come from the accretion disc.

5 THE EJECTION TIME

Another important question concerns the duration of the ejection events. When the blob becomes visible in the radio, it has a size of a few light-days. It seems unlikely that it corresponds to an ejection duration lasting longer than that, while a shorter ejection time may be possible. In the latter case, the requirement on the initial kinetic power correspondingly increases. Note that direct radio observations of flare events do not directly constrain the ejection time, because the radio flux eventually produced in the first parts of the jet is heavily self-absorbed. Again, the minimum power requirement criterion favours ejection events that last for a few light-days. It seems unlikely that it corresponds to an ejection duration lasting longer than that, while a shorter ejection time may be possible. In the latter case, the requirement on the initial kinetic power correspondingly increases. Note that direct radio observations of flare events do not directly constrain the ejection time, because the radio flux eventually produced in the first parts of the jet is heavily self-absorbed. Again, the minimum power requirement criterion favours ejection events that last for \( t_{\text{out}} \sim \) a few d. This corresponds to \( \sim 10^{-5} R_{S,10} c / \)c, where \( R_{S,10} \) is the Schwarzschild radius for a 10-\( M_\odot \) black hole. The immediate consequence is that the ejection event must be considered as a continuous and stationary process. Scaling for a superluminal AGN, we would have, for a 10\(^9\) - \( M_\odot \) black hole, an ejection phase lasting for \( 2 \times 10^5 \) yr.

6 DISCUSSION

We have obtained a reliable lower limit to the kinetic power corresponding to major ejection events in superluminal galactic sources, in particular for GRS 1915+105. This limit is of the order of \( 3 \times 10^{46} \) erg \( \text{s}^{-1} \), much greater than the observed radiative luminosity, which is probably Eddington-limited to values of the order of \( 10^{39} \) erg \( \text{s}^{-1} \), corresponding to a black hole of 10-\( M_\odot \). This by itself suggests that the jet acceleration mechanism cannot be radiative.

We have investigated the role of e\(^{-}\) pairs, normal plasma and the magnetic field as energy carriers of the kinetic power in the inner jet regions, excluding an important role for the e\(^{-}\) pairs, and favouring a scenario in which the initial acceleration phase is controlled by a magnetic field of the order of \( 10^9 \) G or more, able to channel and accelerate accretion disc matter in the jet.

Conservation of kinetic power dictates that the injection timescale is of the order of a day, a very long time-scale if measured in units of the light-crossing time of a 10-\( M_\odot \) Schwarzschild radius, indicating a stationary, not impulsive, process. Shorter injection times, although possible, correspond to larger kinetic powers, exacerbating the problem of how to obtain them. The rough equality between the value of the magnetic field needed to carry the kinetic power in the inner jet and the value necessary to tap the rotational energy of a Kerr black hole via the Blandford & Zhakajev process can be regarded as circumstantial evidence that this process is indeed the one responsible for jet formation and acceleration. This process is also a candidate to power the jets in radio-loud quasars, and we can compare the results obtained here with the corresponding estimates of the kinetic power of AGN derived by Celotti, Padovani & Ghisellini (1997) for a sample of radio-loud sources. For those AGN, the above authors find a kinetic power between \( 10^{45} \) and \( 10^{49} \) erg \( \text{s}^{-1} \), of the same order as the luminosity needed to ionize the broad-line region of the same objects, and in agreement with the power required by the existence of the outer radio lobes. In the AGN case the kinetic and accretion luminosities are roughly equal, while in Galactic superluminal objects the kinetic power dominates. Another obvious difference is the Lorentz factor of the bulk motion, which is much greater in the AGN case. This results in a ratio of the outflowing to infalling mass rate \( M_{\text{out}} / M_{\text{in}} \) of order unity for Galactic superluminal objects and two orders of magnitude smaller for AGN. These estimates make Galactic superluminal sources the most efficient engines to produce collimated relativistic bulk motion, with the possible exception of gamma-ray bursts.

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