Core shifts in blazar jets

Andrzej A. Zdziarski,1 Marek Sikora,1 Patryk Pjanka2 and Alexander Tchekhovskoy3,4

1Centrum Astronomiczne im. M. Kopernika, Bartycka 18, PL-00-716 Warszawa, Poland
2Observatorium Astronomiczne Uniwersytetu Warszawskiego, Al. Ujazdowskie 4, PL-00-478 Warszawa, Poland
3Department of Physics and Department of Astronomy, University of California, Berkeley, CA 94720–3411, USA
4Lawrence Berkeley National Laboratory, 1 Cyclotron Rd, Berkeley, CA 94720, USA

Submitted 2014 October 16

ABSTRACT

We study the effect of core shift in jets, which is the dependence of the position of the jet radio core on the frequency. We derive a new method to measure the jet magnetic field based on both the value of the shift and the observed flux, which compliments the standard method assuming equipartition. Using both methods, we re-analyse the blazar sample of Zamaninasab et al. We find that equipartition is satisfied only if the jet opening angle in the radio core region is close to the values found observationally, ≃0.1–0.2 divided by the bulk Lorentz factor, Γ. Larger values, e.g., 1/Γ, would imply very strong departures from equipartition. A small jet opening angle implies in turn the magnetization parameter of ≲1. We determine the jet magnetic flux taking this effect into account. We find that the average jet magnetic flux is compatible with the model of jet formation due to black-hole spin energy extraction and accretion being magnetically arrested. We calculate the jet average mass-flow rate corresponding to this model, and find it consists of a substantial fraction of the mass accretion rate. We also calculate the average jet power, and find it moderately exceeds Ṁc2, reflecting black-hole spin energy extraction.

Key words: acceleration of particles–ISM: jets and outflows–magnetic fields–radiation mechanisms: non-thermal.

1 INTRODUCTION

We study here extended jets, which low-frequency emission originates in both a part of the jet which is optically thick to synchrotron self-absorption, and one which is optically thin (Blandford & Königl 1979). Then, the partially self-absorbed emission peaks at a distance ≃ν−1 along the jet (where ν is the frequency). The dependence of core position on frequency of observation is called the core shift. In this work, we study this effect theoretically.

Then, we study the samples of blazars and radio galaxies of Zamaninasab et al. (2014), hereafter Z14. We apply to them our theoretical results, and study equipartition and the jet opening angles. We then re-consider the application by Z14 of the model of jet formation from black-hole spin-energy extraction (Blandford & Znajek 1977), with the accretion being magnetically arrested. The powers of (1 + z) in remaining formulae follow from that, and they are (p + 4)/2 in equation (23), (3 − p)/2 in equation (24), and (p − 3)/2 in equation (26). Also, the power of the Doppler factor in equation (26) has been misprinted as 1 instead of 1/2.

2 CORE SHIFT IN THE MODEL OF BLANDFORD & KÖNIGL

Lobanov (1998) and Hirotani (2005). We use the formulation of the model of Blandford & Königl (1979) of Zdziarski et al. (2012) (ZLS12), adding to their formulae the dependences on the cosmological redshift, which we denote as z∗ to distinguish from the jet height z. The jet is emitting above some minimum height, z0, over its length. Its
Figure 1. An example of the dependences of the flux per ln $z$ at three different frequencies, 2, 8 and 32 GHz shown by the blue dotted, red dashed and black solid curves, respectively. The peaks correspond to the position of the radio core. In this example, $p = 2$, the jet begins at $z_{0} \approx 3 \times 10^{15}$ cm, at which it is optically thin above $\nu_{t} \approx 3 \times 10^{14}$ Hz.

The synchrotron spectrum consists of a low-frequency part, in which a part of the jet up to some height is optically thick to synchrotron self-absorption and the rest is optically thin, and a high-frequency part, in which the entire emitting part of the jet is optically thin. The boundary is called the turnover frequency, $\nu_{t}$, and its value depends on $z_{0}$. The partially optically-thick regime has the energy spectral index of 0, and thus $F_{\nu} = \text{constant}$ at $\nu \lesssim \nu_{t}$. The model of Blandford & Königl (1979) assumes conservation of the relativistic-electron flux and the toroidal magnetic energy flux in a jet with constant both the bulk Lorentz factor, $\Gamma$, and the (half) opening angle, $\Theta$, implying

$$K(z) = K_{0}(z/z_{0})^{-2}, \quad B(z) = B_{0}(z/z_{0})^{-1}, \quad N(\gamma, z) = K(z)\gamma^{-p},$$

where $K$ is the normalization of the electron distribution, $N$, $\gamma$ is the Lorentz factor, $\gamma_{\text{min}} < \gamma \leq \gamma_{\text{max}}$, and $p$ is the electron index.

Often, we know neither $\nu_{t}$ nor $z_{0}$. However, in the partially optically-thick regime, which we consider here, the emission at a given frequency is mostly emitted by a narrow range of height, $\nu \propto z^{-1}$. Thus, the actual value of $z_{0}$ is of no importance for emission below the turnover. Therefore, we can parameterize the jet using the dependences (1) down to an arbitrary position, which we take at the gravitational radius, $r_{g} \equiv GM/c^{2}$ (where $M$ is the black-hole mass), defining the values of $B$ and $K$ there as $B_{k}$, $K_{k}$ (though, we stress, this does not imply a jet emission there but only provides a parameterization). Then, we have

$$K(z) = K_{k}(z/r_{g})^{-2}, \quad B(z) = B_{k}(z/r_{g})^{-1}, \quad \nu_{t}(z) = \nu_{t}(z_{0}/r_{g})^{-1},$$

where $\nu_{t}$ is the turnover frequency at $z_{0} = r_{g}$.

The peak flux per unit ln $z$ is emitted at $z \approx r_{g}\nu_{t}/\nu$. This follows, e.g., from equation (21) of ZLS12, which gives the spatial profile of the emission provided the outer integration is not carried out. We can calculate those profiles in the optically thick and thin regimes, which give $dF/d\ln z \propto z^{1/2}$ and $z^{1/p-2}$, respectively, see Fig. 1 and the intersection at $z = (\pi/4)^{(p-2)/(p+4)} r_{g} \nu_{t}/\nu$, where the numerical factor is indeed close to unity for the usual power-law electrons, in particular in ions (excluding the rest energy). This is because the magnetic pressure depends on the field configuration, being $B^{2}/8\pi$ and $B^{2}/24\pi$ and for a toroidal field and fully tangled field, respectively, e.g. Leahy (1991), and the particle pressure depends on its adiabatic index. Also, the energy density ratio is usually used for defining equipartition in astrophysics. From the above definition and $K/B^{2} = K_{k}/B_{k}^{2}$, we obtain the magnetic field strength (in the comoving frame) at the height $z_{\ast}$, we denote $B_{\beta}(z) = z^{-1} \times $

$$
\left[\frac{hD_{i}\Delta\theta}{(\nu_{1}^{2} - \nu_{2}^{2})(1 + z_{i})} \right]^{\frac{p+4}{2p}} \frac{B_{\gamma}}{B_{\ast}} \frac{\nu_{i}^{2}}{\nu_{\ast}^{2}} \frac{8\sigma_{T}(1 + k)f}{\beta C(z/p)\sigma_{T}\tan\Theta_{i}} \right]^{\frac{1}{p}}
$$

This expression is almost the same as the corresponding equation (43) of Hirota (2005) except for his dependence on $(1 + z_{i})$, which appears incorrect. Consequently, an incorrect power of $(1+z_{i})$ appears in the expressions of O’Sullivan & Gabuzda (2009), Pushkarev et al. (2012) and Z14, who used the result of Hirota (2005) for $p = 2$. The numerical coefficient at $p = 2, k = z_{i} = 0, \gamma_{\text{min}}/\gamma_{\text{max}} = 10^{-34}$ (corresponding to $f = 10$ assumed by Hirota (2005), $z$ and $D_{i}$ in pc, $\nu$ in GHz and $\Delta\theta$ in mas is $1.42 \times 10^{-8}$.

2 This discrepancy originates in the assumption of Hirota (2005) that the special-relativistic invariance includes $z_{i}$, see his equations (17–18), which leads to the transformation of $\sin i$ from the jet to the observer frame including $(1+z_{i})$. However, the transformation from the jet frame to the stationary frame at the jet redshift is $\sin i = \delta^{-1} \sin i’$. Then, the photon travels to the observer radially without changing its direction, as well as the jet axis is no more a distinct direction. Also, note that equation (10) of Lobanov (1998), giving $B$ as a function of the core shift, contains some incorrect exponents.
which small difference with respect to $1.45 \times 10^{-8}$ in the expression of O’Sullivan & Gabuzda (2009) (including their coefficient for $\Omega_{c,1}$) is due to rounding errors in the latter.

On the other hand, the above method ignores the information contained in the flux in the partially self-absorbed spectrum, $F_{\nu}$. In the model of Blandford & Königl (1979), it can be used to find the magnetic field without any assumptions about equipartition. We use here equation (22) of ZLS12 to relate $F_{\nu}$ [in erg/(cm$^2$ s Hz)] to $\nu$ and $B_{\nu}$. This way, we derive

$$B_{\nu}(z) = \frac{D_{\nu}h(m_{\nu})\sqrt{\epsilon_{\nu}}}{1 + z_{\nu}} \left( \frac{\Delta \theta}{\nu_{1}^{-1} - \nu_{2}^{-1}} \right)^{\frac{1}{2}} \frac{\sigma_{i}C_{i}(p)C_{i}(p) \sin \theta}{2\pi \nu^{9} C_{(p)F_{\nu}}},$$

(8)

where $C_{i}(2, 3) \simeq 1.14$, $C_{i}(2, 3) = 3.61$, $2.10$, respectively, are coefficients defined in ZLS12. The numerical coefficient at $p = 2$, $z$ and $D_{\nu}$ in pc, $\nu$ in GHz, $F_{\nu}$ in Jy and $\Delta \theta$ in mas is $3.35 \times 10^{-11}$, and it is only weakly dependent on $p$ through $C$. Note the relatively high power of $\Delta \theta$, requiring its accurate measurement. Both formulae for $B(\nu)$ are independent of $M$.

We can also calculate the equipartition coefficient using equations (8,9). We obtain

$$\beta = \frac{f(1 + k)m_{e}^{2+3}c^{3+5}p^{4+5} \nu_{1}^{4-1}}{D_{\nu}^{2}r_{h}^{2+3} \nu_{1}^{4-1} \nu^{6} F_{\nu}^{6} \sin \theta} \times \left( \frac{C_{i}}{c_{i}} \right) \left( \frac{\sigma_{i}F_{\nu}}{C_{(p)C_{i}}} \right).$$

(9)

The numerical coefficient at $p = 2$, $f = 1$, $k = 0$, $D_{\nu}$ in pc, $\nu$ in GHz, $F_{\nu}$ in Jy and $\Delta \theta$ in mas is $3.28 \times 10^{8}$. Given the high powers of most of the measured quantities, an application of this formula can yield a relatively large fractional error.

3 EQUIPARTITION IN THE SAMPLE OF Z14

We re-analyse the sample of Z14, which contains 76 blazars and radio galaxies. This allows us to study in detail their average properties. Hereafter, we assume $\Theta_{0} \equiv 1$, implying $\Gamma = \frac{\Theta_{0}}{\Theta_{0}}$.

In calculating the inferred magnetic field at $z = 1$ pc in radio galaxies, Z14 applied their equation (3) using estimates of $i$ from literature, from which $\Gamma(i, \Theta_{0})$ was calculated, where $\theta_{\text{app}}$ is the apparent velocity, which they used to calculate $\delta$. On the other hand, they made some approximate substitutions for blazars, namely $\Gamma_{i} = \Gamma_{\min} = (1 + \beta_{\text{app}})^{1/2}$, which is the minimum Lorentz factor for a given observed apparent velocity, $i = \Gamma_{\min}$ (in which case $\delta = \Gamma_{\min}$), and $\Theta_{0} = 0.13 \Gamma_{\min}^{-1}$ (from the fit of Pushkarev et al. 2009). These assumptions were used in obtaining their equation (4) from their equation (3). Adopting the same further assumptions as Z14, namely $p = 2$, $f = 2$ and $\beta = 1$, we obtain slightly higher values of $B_{\nu}$ than Z14, owing to our formula having the corrected power of $(1 + z_{\nu})$.

However, as Z14 note, the relationship between $\Theta_{0}$ and $\Gamma_{i}$ is observationally determined at larger distances than the radio cores, and it is not known at the radio cores. In fact, Z14 used $\Theta_{0} = \Gamma_{i}^{-1}$ in some of their following discussion. Thus, we allow here the coefficient of the opening angle, $\theta_{i}$, to be a free parameter, $\Theta_{0} / \Gamma_{i}$.

To be able to calculate the magnetic field, $B_{\nu}$, from equation (8), we need to know the flux. For the sample of blazars, we get $F_{\nu}$ measured at 15 GHz at the same time as the core shift from the MOJAVE (Lister et al. 2009). For radio galaxies, we have found the radio fluxes measured at the same time as the core shifts of 6 out of 8 sources in the sample of Z14. In a few cases, the core shift was based on two observations and/or two radio bands, for which we took the average flux. We do not consider the upper limits on the core shift (9 blazars in the sample of Z14).

We then define a reference value of $\theta_{0} \equiv \theta_{i}$ by $(\langle B_{\nu}/B_{0} \rangle = 1$ (hereafter the symbol $\langle \rangle$ denotes a geometric average), assuming all sources have the same $\theta_{i}$ and $\beta = k = 1$. Notably, we find $\theta_{0} \approx 0.12$, very close to the best-fit value of Pushkarev et al. 2009. Fig. 2 shows the histogram of $B_{\nu}/B_{0}$ at $\theta_{i}$. We consider the errors in determining the core shift to be a major cause of the relatively large scatter seen in Fig. 2 corresponding to the standard deviation of $\log B_{\nu}/B_{0}$ of $\approx 1.1$. Still, there is a pronounced peak at $B_{\nu}/B_{0} \approx 2-3$. This argues for a value of $\beta$ with a relatively small intrinsic dispersion in the sample.

On the other hand, we see that all the radio galaxies for which we have $F_{\nu}$ have $B_{\nu}/B_{0} \lesssim 0.2$, and $(\langle B_{\nu}/B_{0} \rangle = 0.04$. This is a systematic effect unlikely to be due to the measurement errors. This indicates that those objects have properties different from blazars, possibly the magnetic fields in those objects are much below equipartition, or, alternatively, they have substantially wider opening angles.

Hereafter we consider only blazars. Proceeding as above, we find $\theta_{0} \approx 0.11$ at $\beta = 1$, and $(\langle B_{\nu}/B_{0} \rangle = 1.3$ G at 1 pc. On the other hand, we have the following dependences from equations (2,9):

$$\langle \theta_{i} / \Theta_{0} \rangle = (\beta)^{-1/(p+7)}, \quad \langle B_{\nu}/B_{0} \rangle = (\beta)^{-2/(p+6)} = (\theta_{i} / \Theta_{0})^{2/(p+6)}(10).$$

Then, if the true value average of $\theta_{i}$ were different from $\theta_{0}$, the distribution in Fig. 2 would be shifted by $(\theta_{i} / \Theta_{0})^{2/(p+6)}$ preserving the shape of the histogram. Thus, we can have a larger value of $\theta_{i}$ if $\beta < 1$. Still, $(\langle B_{\nu}/B_{0} \rangle = 160$ for $\theta_{i} = 1$. This corresponds to a very large departure from equipartition, $(\beta) \approx 1.6 \times 10^{-9}$, with the standard deviation of $\log \beta$ of $\approx 4.4$. Since such large departures from equipartition are unlikely, this argues for $(\theta_{i}) \lesssim 1$ in this sample (in agreement with the observations, see above).

4 IMPLICATIONS FOR THE JET PHYSICS

The magnetization parameter is the ratio of the Poynting flux to the kinetic energy flux in the black-hole frame. In the jet rest frame, it is equal to the ratio of the proper magnetic enthalpy, $w_{\text{B}}$, to that for particles including the rest energy, $w_{p}$.

$$\sigma = \frac{w_{\text{B}}}{w_{p}} = \frac{B^{2}/4\pi}{\eta \rho_{p} + \rho c^{2}},$$

(11)

where $\rho$ is the rest-mass density, $4/3 < \eta < 5/3$ is the particle adiabatic index and the magnetic field is predominantly toroidal.
Using equation (6), $\beta < 2/(\eta c^3)$. If we know both $\sigma$ and $\beta$, we can constrain the plasma parameters,

$$
\frac{u_p}{\rho c^2} = \frac{\beta \sigma/2}{1 - \beta \eta/2}, \quad \frac{B^2/4\pi}{\rho c^2} = \frac{\sigma}{1 - \beta \eta/2}.
$$

(12)

Also, the opening angle is approximately related to $\sigma$,

$$
\Theta_j = \frac{a \sigma}{\Gamma_j}, \quad \Theta_j = \frac{a \sigma}{\Gamma_j}.
$$

(13)

where $a \leq 1$ (Tchekhovskoy et al. 2009; Komissarov et al. 2009). Since $(\Theta_j) \sim 0.1$–0.2 at least at large distances beyond the radio core (Pushkarev et al. 2009; Clausen-Brown et al. 2013; Forstad et al. 2005), the average value of $\sigma$ there is $\ll 1$. A similar situation occurs in black-hole binaries, whose jets have the opening angles $\ll 1/\Gamma_j$ (Miller-Jones et al. 2006). Thus, some mechanisms in the jet have to be able to decrease $\sigma$ below unity, see Komissarov (2011), Tchekhovskoy et al. (2009), Lyubarsky (2010).

The Bernoulli equation for negligible energy losses is,

$$
\Gamma_j \left[1 + \frac{\eta e B^2/4\pi}{\rho c^2}\right] = \Gamma_j \left[1 + \sigma - \beta \eta/2\right] = \Gamma_{max},
$$

(14)

where $\Gamma_{max}$ is the Lorentz factor corresponding to the conversion off all of the magnetic and particle internal energy into acceleration. This gives $\Gamma_j$ as a function of $\sigma$ and $\beta$. We can then use equations (12,14) together with the conservation of mass, equation (15), below, to determine the evolution of the toroidal magnetic field and particle energy density along the jet. We derive

$$
B^2(c) = B_0^2 \left[1 + \gamma_0 \sigma + \frac{(1 + \beta \eta/2)^2}{c^2} \right] \left[1 + \frac{\eta e B^2/4\pi}{\rho c^2}\right], \quad u_p = u_{0p} \left[1 + \frac{(1 + \beta \eta/2)^2}{c^2} \right],
$$

(15)

where $B(c)$ differs from the corresponding dependence of equation (14), which assumes constant $\Gamma_j$ and $\Theta_j$, whereas they vary in the present approach. Here $u_{0p}$ is the particle energy density at a reference point $z_0$ (with $\sigma = \gamma_0$). Note that equation (15) also conserves the enthalpy flux, equation (20) below. In deriving equation (15), we have defined $\Theta_j$ as the ratio of the jet radius to height, in which case the limit of $\sigma = 0$ cannot be applied.

With equations (13,14), we can rewrite equation (5) of Z14, giving the jet magnetic flux in the model with extraction of the rotational power of the black hole (Blandford & Znajek 1977), as

$$
\Phi_j = \frac{2\pi r_{hl} a B}{\ell a} \frac{1}{\sigma^{1/2}(1 + \beta \eta/2)},
$$

(16)

where $r_{hl} = [1 + (1 - a^2)^{1/2}]r_g$ is the black-hole horizon radius, $a$, is the dimensionless spin parameter, and $\ell = 0.5$ is the ratio of the angular frequency of the field lines to that of the black hole. See Z14 for the adopted assumptions and references. Hereafter $a_\ast = 1$ is assumed, following Z14. We determine the values of $B$ using equation (7), i.e., $B = B_0$, but impose the condition of $\langle B_r \rangle = (B_0)$, which implies $\beta = 0.11$, see Section 2. We do not directly use the values of $B_r$ due to their large errors associated with the high power of the core shift, see Section 2.

We can then equate $\Phi_j$ to the poloidal flux threading the black hole on one hemisphere, which is limited by the ram pressure of the accretion flow (Narayan et al. 2003). This can be written as

$$
\Phi_{mzj} = \Phi_{blz}(M_{\ast})^{1/2}r_{hl},
$$

(17)

where saturation values of $\Phi_{blz} \approx 50$–$100$ have been found in GRMHD simulations of magnetically arrested accretion (Tchekhovskoy et al. 2011; McKinney et al. 2012). Z14 estimated $M$ as $L/\epsilon$, where $L$ is the estimated bolometric luminosity, assuming the efficiency of $\epsilon = 0.4$. For the sample of Z14, $\ell = 0.5$, $a = k = \beta = 1$, and $\eta = 3/2$ for a mixture of relativistic electrons and nonrelativistic protons, we obtain $\langle \Phi_{blz} \rangle \approx 180(\epsilon/0.4)^{1/2}$. The cause of the difference with respect to the result of Z14 of $\approx 50$ is our use of a lower jet opening angle, motivated both observationally and by the equipartition argument. However, the average radiative efficiency in magnetically arrested discs is unknown and may be $< 0.4$; for $\epsilon = 0.2$, $\langle \Phi_{blz} \rangle = 130$, and the standard deviation of $\Phi_{blz}$ is 0.23. Within this uncertainty and systematic uncertainties of our adopted idealized model, this is in a satisfactory agreement with the simulation results. Thus, our results confirm that blazars can have jets originating from magnetically arrested discs and are powered by the Blandford-Znajek mechanism, but at the same time they can have $\sigma \ll 1$ at the radio core.

Note that both $\Phi_j$ and $\Phi_{bi}$ depend linearly on the black-hole mass. Thus, the actual strength of the correlation is completely independent of it. Therefore, we present the correlation (for blazars only) for both quantities divided by $M$ in Fig. 3. We see a relatively good correlation, though the visual scatter is much larger than in fig. 2 of Z14, where the common dependence on $M$ was included.

5 THE JET MASS-FLOW RATE AND POWER

The total jet+counterjet mass-flow rate is $M_j = 2\pi \rho c(\Theta_j c)^2 \Gamma_j$. Using equations (12,13), it can be written as

$$
M_j = \frac{(B_a c)^2}{2\epsilon \Gamma_j} \left(1 - \eta \beta / 2 \right).
$$

(18)

For the sample of Z14 at $a = k = \beta = 1$, we obtain $\langle M_j / M \rangle \approx 0.15(\epsilon/0.2)$, i.e., a relatively large fraction of the accretion flow is channelled into the jet. This further supports a modest value of the accretion radiative efficiency. Fig. 4(a) shows the dependence of this ratio on the Eddington ratio (for the H abundance of $X = 0.7$) for blazars. We see that most of the sources have $L/L_E$ in the 0.1–1 range, and there is no apparent trend seen.

The jet-counterjet power for $\Gamma_j \gg 1$ is $P = 2\pi \nu c (\Theta_j c \Gamma_j)^2$ (e.g., Levinson 2004). We find the power in the magnetic field and the total power, respectively, as

$$
P_B = \frac{(B_c a)^2 c}{8\pi} = \frac{(\Phi_{bi} c)^2 M_c^2}{8\pi^2} \frac{\sigma (1 + \eta \beta / 2)}{1 + \sigma},
$$

(19)

$$
P_j = \frac{(B_c a)^2 c}{8\pi} \frac{1}{1 + \sigma} = \frac{(\Phi_{bi} c)^2 M_c^2}{8\pi^2} \frac{\sigma (1 + \eta \beta / 2)}{1 + \sigma}.
$$

(20)

For the sample of Z14 at $a = k = \beta = 1$, we obtain $(P_j / M_c^2) \approx 1.7(\epsilon/0.2)$. Within the systematic uncertainties, this agrees well with the average from numerical simulations (Tchekhovskoy et al. 2011; McKinney et al. 2012). Fig. 4(b) shows the dependence of

\[\begin{figure}
\centering
\includegraphics[width=\textwidth]{figure3.png}
\caption{The correlation between $\Phi_j / M$ and $L_j^{1/2}$ (normalized as in equation (17) for blazars. The dashed line corresponds to $\langle \Phi_{bi} \rangle$.}
\end{figure}\]
with modelling of blazars for the strong magnetic jets measured by the core shift method pointed out by Nalewajo et al. (2014).

We have calculated the jet mass-flow corresponding to the above model, and found it consists of a substantial fraction of the mass accretion rate, $\dot{M}_j \approx 0.15 M_\odot$ for our chosen parameters. For those, the jet power is $P_j \approx 1.7 M^2 c^2$, consistent with the results of numerical simulations. Finally, we have found that the radio galaxies in the sample of Z14 appear to have substantially lower magnetic fields and fluxes.

7 ACKNOWLEDGEMENTS

We thank Alexander Pushkarev for help with using radio data, and Tuomas Savolainen for providing us with the inclinations used by Z14 for the sample of the radio galaxies. This research has made use of data from the MOJAVE database that is maintained by the MOJAVE team, and it has been supported in part by the Polish NCN grants 2012/04/M/ST9/00780, 2013/10/M/ST9/00729 and DEC-2011/01/B/ST9/04845. AT was supported by NASA through Einstein Postdoctoral Fellowship grant number PF3-140115 awarded by the Chandra X-ray Center, which is operated by the Smithsonian Astrophysical Observatory for NASA under contract NAS8-03060.

REFERENCES

Blandford R. D., Königl A., 1979, ApJ, 232, 34
Blandford R. D., Znajek R. L., 1977, MNRAS, 179, 433
Clausen-Brown E., Savolainen T., Pushkarev A. B., Kovalev Y. Y., Zensus J. A., 2013, A&A, 558, A144
Hirotani K., 2005, ApJ, 619, 73
Jorstad S. G., et al., 2005, AJ, 130, 1418
Komissarov S. S., 2011, Mem. Soc. Astron. Ital., 82, 95
Komissarov S. S., Vlahakis N., Königl A., Barkov M. V., 2009, MNRAS, 394, 1182
Leahy J. P., 1991, in P. A. Hughes, ed., Beams and Jets in Astrophysics. Cambridge Univ. Press, Cambridge, p. 100
Levinson A., 2006, Int. J. Mod. Phys., 21, 6015
Lister M. L., et al., 2009, AJ, 137, 3718
Lobanov A. P., 1998, A&A, 330, 79
Lyubarsky Y. E., 2010, MNRAS, 402, 353
McKinney J. C., Tchekhovskoy A., Blandford R. D., 2012, MNRAS, 423, 3083
Miller-Jones J. C. A., Fender R. P., Nakar E., 2006, MNRAS, 367, 1432
Nalewajo K., Sikora M., Begelman M. C., 2014, ApJ, in press, arXiv:1410.4571
Narayan R., Igumenshchev I. V., Abramowicz M. A., 2003, PASJ, 55, L69
O’Sullivan S. P., Gabuzda D. C., 2009, MNRAS, 400, 26
Pushkarev A. B., Kovalev Y. Y., Lister M. L., Savolainen T., 2009, A&A, 507, L33
Pushkarev A. B., Hovatta T., Kovalev Y. Y., Lister M. L., Lobanov A. P., Savolainen T., Zensus J. A., 2012, A&A, 545, A113
Tchekhovskoy A., 2015, in I. Contopoulos, et al., eds., The Formation and Disruption of Black Hole Jets. Springer, in press
Tchekhovskoy A., McKinney J. C., 2012, MNRAS, 423, L55
Tchekhovskoy A., McKinney J. C., Narayan R., 2009, ApJ, 699, 1789
Tchekhovskoy A., Narayan R., McKinney J. C., 2011, MNRAS, 418, L79

© 2014 RAS, MNRAS 000, 1-4
Zamaninasab M., Clausen-Brown E., Savolainen T., Tchekhovskoy A., 2014, Nature, 510, 126 (Z14)
Zdziarski A. A., 2014, MNRAS, 445, 1321
Zdziarski A. A., Lubiński P., Sikora M., 2012, MNRAS, 423, 663 (ZLS12)