Does the plasma composition affect the long term evolution of relativistic jets?

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ABSTRACT

We study the influence of the matter content of extragalactic jets on their morphology, dynamics and emission properties. For this purpose we consider jets of extremely different compositions including pure leptonic and baryonic plasmas. Our work is based on two-dimensional relativistic hydrodynamic simulations of the long-term evolution of powerful extragalactic jets propagating into a homogeneous environment. The equation of state used in the simulations accounts for an arbitrary mixture of electrons, protons and electron-positron pairs. Using the hydrodynamic models we have also computed synthetic radio maps and the thermal Bremsstrahlung X-ray emission from their cavities.

Although there is a difference of about three orders of magnitude in the temperatures of the cavities inflated by the simulated jets, we find that both the morphology and the dynamic behaviour are almost independent on the assumed composition of the jets. Their evolution proceeds in two distinct epochs. During the first one multidimensional effects are unimportant and the jets propagate ballistically. The second epoch starts when the first larger vortices are produced near the jet head causing the beam cross section to increase and the jet to decelerate. The evolution of the cocoon and cavity is in agreement with a simple theoretical model. The beam velocities are relativistic \( \Gamma \approx 4 \) at kiloparsec scales supporting the idea that the X-ray emission of several extragalactic jets may be due to relativistically boosted CMB photons. The radio emission of all models is dominated by the contribution of the hot spots. All models exhibit a depression in the X-rays surface brightness of the cavity interior in agreement with recent observations.
Key words: hydrodynamics – relativity – plasmas – ISM: jets and outflows – radiation mechanisms: thermal – galaxies: jets.

1 INTRODUCTION

The standard model for powerful jets associated with extragalactic radio sources (Blandford & Rees 1974) assumes that the energy radiated in the radio lobes of such sources is produced in the active nuclei of their host galaxies, the central engine being a supermassive black hole surrounded by an accretion disc. The accretion process fuels a couple of twin supersonic jets which transport away bulk kinetic energy from the neighbourhood of the central black hole to the lobes, i.e., from scales of the Schwarzschild radius of the supermassive black hole

\[ R_S = \frac{2GM_{bh}}{c^2} = 3 \cdot 10^{13} \cdot \left(\frac{M_{bh}}{10^8 M_\odot}\right) \text{cm}, \]





to kiloparsec scales. This bulk kinetic energy is dissipated by shocks within the beam and (mostly) at the jet terminal shocks where electrons are accelerated and radiate via synchrotron and inverse Compton mechanisms. Typical lifetimes and kinetic powers of powerful radio sources are \( \approx 10^7 \text{y} \) and \( 10^{44} - 10^{47} \text{erg s}^{-1} \), respectively (Rawlings & Saunders 1991; Daly 1995).

Among the problems that still remain open even after more than 30 years of research is the composition of extragalactic jets. Within the standard model (Blandford & Rees 1974) jets are made of a relativistic plasma that contains relativistic electrons and thermal protons (ep-plasma). On the other side, based on equipartition arguments Kundt & Gopal-Krishna (1980) claim that extragalactic jets consist of beams of extremely relativistic electrons and positrons (e\( ^\pm \)-plasma) with almost no ions. Ultimately, the composition of jets is tightly related to their formation mechanisms. As discussed by Celotti & Blandford (2001), electromagnetically dominated outflows, as those generated by the extraction of spin energy of the black hole, will become pair dominated jets with a low baryonic pollution. Jets generated from the accretion disk by hydromagnetic winds will be made of baryonic plasma. Probably, both processes are operating simultaneously in nature. Sol, Pelletier & Asseo (1989) propose a two-flow model where the jet consists of a beam of relativistic particles (pair plasma) surrounded by a Newtonian or mildly relativistic (ep-plasma) wind accelerated from the disk. It is also possible that ep-jets become pair-loaded later on, e.g., by interactions with high energy photons from the disk corona – e.g., Begelman, Blandford & Rees (1984) – or by proton-proton collisions at parsec scale – Anyakoha, Okeke & Okoye (1988) –.

Observations show that a number of jets are highly polarised. This is a further argument
in favour of jets made of e\textsuperscript{\pm} pairs as there is no internal Faraday rotation or depolarisation (the rotation produced by the electrons is compensated by that of the positrons), while an ep-plasma gives rise to these effects. If jets were dominated by ep-plasmas, observational limits on the degree of Faraday rotation and depolarisation imply thermal electron densities $< 10^{-3}$ cm$^{-3}$ – e.g., Walker, Benson & Unwin (1987) – or a minimum energy cut-off for the electrons $\approx 50$ MeV (Wardle 1977). This suggests a lack of thermal (cold) matter, at least, in the most polarised sources. However, as the upper bounds on the thermal electron density derived from the depolarisation argument depend on the magnetic field structure, they might be underestimated if there are field reversals within the VLBI beam. Finally, the detection of circular polarisation at parsec scales in several radio sources (Wardle et al.1998; Homan & Wardle 1999) suggests that, in general, extragalactic radio sources are mainly composed of e\textsuperscript{\pm}-plasma. The argument is based on the fact that circular polarisation is produced by Faraday conversion, which requires that the energy distribution of the jet emitting particles extends to very low energies. This in turn indicates that e\textsuperscript{\pm}-pairs are an important component of the jet plasma.

Sikora & Madejski (2000) conclude that X-ray observations of blazars associated with OVV quasars impose strong constraints on the e\textsuperscript{\pm} pair content of jets of radio-loud quasars. According to these authors, pure electron-positron pair jets can be excluded because they may produce too much soft X-rays, while pure electron-proton jets may emit too few non-thermal X-ray radiation. Therefore, only jets may be viable where the number density of electron-positron pairs is much larger than the proton number density, but small enough for protons still being dynamically more important. Hirota et al.(2000) find that parsec-scale jets cannot be dominated by baryons, if the electron density determined by the amount of synchrotron self-absorption necessary for an optically thick jet component, is of the same order as the one derived from the kinetic luminosity.

For jets composed of a pair plasma additional questions arise. How can they maintain their stability, as for a given particle density they are lighter and hence less stable than jets made of an electron-proton plasma? Can they transport enough energy to fuel the observed kiloparsec scale radio lobes? Concerning the stability of electron-positron plasma flows, Sol et al.(1989) demonstrate that the two-flow model is stable against excitation of electrostatic waves as long as the magnetic field is larger than a critical value $B_c$. In their two-flow model, the relativistic beam is responsible for the VLBI jet and the observed superluminal motion, while the slower wind gives rise to the kiloparsec scale jet. Achatz & Schlickeiser (1993)
also consider the stability of $e^\pm$-plasmas against the excitation of electromagnetic waves when the magnetic field exceeds $B_c$. They find that the plasma may be stable over large distances provided plasma waves are damped thermally. The problem of the energy supply to large scales with a limited number of particles is turned around by Celotti (1998). She argues that the jets of low power radio sources mainly consist of $e^\pm$-plasma, while those of powerful radio-loud quasars are made of $e^p$-plasma. However, measurements of the circular polarisation of VLBA-jets of several FR II sources by Wardle et al.(1998) favour $e^\pm$-jets also in powerful sources (at least, at parsec scales).

Another important and not yet solved question concerns the impact of the composition on the morphology and dynamics of jets from sub-parsec to kiloparsec scales. Observations indicate that jets are relativistic at parsec scales – e.g., Laing (1996) –, then decelerate until they become sub-relativistic or mildly relativistic – e.g., Bridle et al.(1994) – at kiloparsec scales with advance speeds of the terminal hotspots in the range $0.01c$ to $0.1c$ (Liu, Pooley & Riley 1992; Daly 1995), where $c$ is the speed of light. At large scales the observed morphologies and deceleration of powerful radio sources is governed by interaction with the external medium. Several theoretical models have considered the gross features of kiloparsec scale jets, i.e., their long term evolution. Begelman & Cioffi (1989) – BC89 hereafter – consider a simple model to describe the evolution of the cocoon. Self-similar expansion is suggested by Falle (1991). Komissarov & Falle (1998) explore the large-scale flow caused by classical and relativistic jets in a uniform external medium. They find that jets with finite initial opening angles are recollimated by the high pressure in the cocoon and that the flow becomes approximately self-similar at large times.

Numerical investigations have also addressed the kiloparsec scale regime. Two-dimensional Newtonian hydrodynamic simulations of axisymmetric light jets were performed by Cioffi & Blondin (1992) in order to understand the evolution of the cocoon, and were compared with the simple analytic theory of BC89. The grid resolution is quite high (15 zones per jet radius), but the simulations do only cover a relatively short period of the cocoon’s evolution. The Newtonian simulations of Hooda, Mangalam & Wiita (1994) cover the evolution of axisymmetric extragalactic jets up to $10^8$ y. Their simulations include isothermal atmospheres with a density stratification given by a power-law, surrounded by an even hotter, but less dense intra-cluster medium where the jets accelerate and collimate. Hooda & Wiita (1996) extended the results of Hooda et al.(1994) to three-dimensions, but their simulations cover...
only a distance of 35 (initial) jet radii, and thus are too short to shed light on the long term evolution of real sources.

Martí, Müller & Ibáñez (1998) – MMI98 hereafter – studied the long term evolution of powerful extragalactic jets on the basis of relativistic hydrodynamic simulations (up to an evolutionary time of $3 \cdot 10^6$ y with a relatively low numerical resolution). The results are compared and interpreted with a simple generalisation of the model of BC89. They find an evolution divided into two epochs. After a transient initial stage, the jet’s evolution is dominated by a strong deceleration. The jet advance speed becomes as small as $0.05c$ due to the degradation of the beam flow by means of internal shocks and the broadening of the beam cross section near the hotspot.

The present work extends the investigations of MMI98 to much later evolutionary times and also addresses the question whether the content of thermal matter in powerful kiloparsec jets does influence their morphology, dynamics and emission properties. Like MMI98 we assume that the dynamics of the jets is dominated by the thermal plasma – e.g., Sikora & Madejski (2000) – and, thus, a hydrodynamic approach is appropriate. For this purpose we have performed long-term simulations of axisymmetric, relativistic jets of extremely different composition, i.e., jets made of a pure leptonic or baryonic plasma, propagating into a uniform environment.

2 MODELS

2.1 Equation of state

In almost all previous jet simulations an ideal gas equation of state (EoS) with constant adiabatic index $\gamma$ has been used. This is a good approximation in both the nonrelativistic ($\gamma = 5/3$) and the ultra-relativistic limit ($\gamma = 4/3$). However, when there exist very large temperature gradients in the flow, it is more accurate to use a EoS including a temperature dependent $\gamma$. For the simulations that we present below temperatures range from about $10^7$ K in the ambient medium to $10^{13}$ K in the hotspots. Thus, both relativistic and non-relativistic particles will participate on an equal footing. Additionally, extragalactic jets are likely composed of a mixture of particles of different masses, i.e., the value of $\gamma$ depends on the composition as protons become relativistic at higher temperatures than electrons.

An equation of state that describes a mixture of ideal, relativistic Boltzmann gases has been derived by Synge (1957) (see also Komissarov & Falle 1998, appendix A). The Synge
EoS used in our simulations includes protons, electrons and positrons. The composition of the plasma only changes due to fluid mixing as the production or annihilation of electron-positron pairs can be neglected due to the low gas density in the kiloparsec scale (see, e.g., Ghisellini et al. 1992). Assuming plasma neutrality, only one parameter is needed to fix the composition, e.g., the mass fraction of the leptons

\[ X_l = (\rho_{e^-} + \rho_{e^+})/\rho \]

where \( \rho_{e^-} \) and \( \rho_{e^+} \) are the local rest-mass densities of electrons and positrons, respectively. Using the Synge EoS instead of a constant-\( \gamma \) EoS requires about 50% more computation time, because the iterative (Newton-Raphson) computation of \( T(\varepsilon, \rho, X_l) \) (\( \varepsilon \) being the specific internal energy) involves Bessel functions.

### 2.2 The Parameter Space

The following notation holds throughout the paper. Subscripts \( b \) and \( m \) refer to the beam and the external medium, respectively. The speed of light is set to \( c = 1 \).

Assuming a uniform and static external medium a relativistic jet is fully described specifying, at a given inlet of radius \( R_b \), the density contrast \( \eta \) between the beam and the ambient medium, the pressure ratio \( K = P_b/P_m \) between the beam and the ambient gas, the beam speed \( v_b \) or equivalently the beam Lorentz factor \( W_b = 1/\sqrt{1 - v_b^2} \), the beam Mach number \( M_b = v_b/c_{sb} \) (where \( c_{sb} \) is the beam sound speed), and parameters that depend on the EoS. For an ideal gas EoS only the adiabatic index \( \gamma \) needs to be specified. In this case, a relativistic jet is freely scalable in size and density, but not in velocity due to the scale introduced by the speed of light. Using the Synge EoS (§2.1) \( \gamma \) is replaced by two parameters describing the chemical composition in the jet and in the external medium, e.g., the lepton mass fractions \( X_{lb} \) and \( X_{lm} \), respectively. The Synge EoS also introduces an extra mass scale (due to the electron and proton masses), and hence an additional (mass) parameter, e.g., \( m_0 = \rho_m R_b^3 \).

With this set of parameters \( \{\eta, K, M_b, W_b, m_0, X_{lb}, X_{lm}\} \) a jet model is scalable under the transformations \( t \to at, x \to ax, \rho \to \rho/a^3 \) (\( a = const \)). In the following this scale freedom will not be used, but \( \rho_m \) and \( R_b \) will be fixed instead of \( m_0 \).

Current observations provide only constraints for the parameter space. The kinetic luminosity \( L_{kin} \) of a typical FR II jet is \( 10^{46}\text{erg/s} \) (Rawlings & Sanders 1991; Daly 1995). The initial jet propagation speed can be restricted by observations of CSOs (Compact Source Object) which have linear sizes of less than \( 0.5\text{kpc} \) and which most likely represent a very
early evolutionary stage of a typical powerful radio source. Observed values of the hotspot propagation velocity are \( \approx 0.2c \) (Owsianik & Conway1998; Taylor et al.2000).

The density, temperature and composition of the external medium are \( \rho_m \approx 10^{-3} \text{g/cm}^3 \), \( T_m \approx 10^7 \text{K} \) and \( X_{lm} = m_e/(m_e + m_p) \approx 1/1837 \), respectively (e.g., Ferrari 1998). Values of \( R_b \sim 0.5 \text{kpc} \) can be inferred from observations of kiloparsec scale jets. But \( R_b \) is also constraint through other model parameters (see below). There are also observational constraints on the likely value of the beam Lorentz factor at parsec scale - \( W_b \approx 10 \), e.g., Ghisellini et al.(1993) –. Furthermore, measurements of the degree of polarisation as a function of frequency may be used to set an upper limit on the mean thermal electron density number which is \( n_e \approx 10^{-2} \text{cm}^{-3} \) (Perley, Willis & Scott 1979; Burch 1979; Ghisellini et al.1992).

We have fixed the values of \( L_{\text{kin}}, X_{lm}, \rho_m, T_m \) and the initial propagation speed in all our models. The latter is set equal to

\[
v_j^{1d} = \frac{\sqrt{\eta_R}}{\sqrt{\eta_R} + 1} v_b, \tag{1}
\]

with \( \eta_R = \rho_b h_b W_b^2 / \rho_m h_m \). This is the propagation velocity derived by Martí et al.(1997) for a pressure matched jet that moves in one dimension only (i.e., without sideways expansion). If the jet is not pressure matched, the propagation velocity can still be computed using the previous formula as long as the sound speed in the external medium fulfils \( c_{\text{sm}} << 1 \).

The density contrast \( \eta \) and the composition of the beam \( (X_{lb}) \) are not directly accessible to observations, i.e., these quantities are treated as free parameters. The beam Lorentz factor \( W_b \) cannot be chosen arbitrarily, because there are forbidden areas in the parameter space where the pressure ratio \( K \) and, most importantly, the beam temperature \( T_b \) have unrealistic values (see below). We used values of \( W_b \) in the range 6.6 – 8.

\( L_{\text{kin}} \) is obtained integrating the energy flux (see, Martí et al.1997, Eq. 20) over the beam cross section

\[
L_{\text{kin}} = (h_b W_b - 1) \rho_b W_b \pi R_b^2 v_b. \tag{2}
\]

This differs from the definition that other authors have used in the past – see e.g., Ghisellini (1998) – by the factor \( h_b W_b - 1 \) which accounts for the fact that we do not include the rest-mass energy in the energy density, as it cannot be extracted from the beam particles.

Having fixed \( L_{\text{kin}}, v_j^{1d}, X_{lb} \) and the external medium, the values of \( \eta, T_b \) and \( M_b \) are uniquely determined by \( W_b \) and \( R_b \) through relations (1) and (2). To clarify this point, let us define function
\[ C_1 := \eta h(\eta, T_b) = \frac{h_m}{W_b^2} \left( \frac{v_j^{1d}/v_b}{1 - v_j^{1d}/v_b} \right)^2 \] (3)

and

\[ C_2 := \eta (h(\eta, T_b)W_b - 1) = \frac{L_{\text{kin}}}{\rho_m W_b \pi R_b^2 v_b}. \] (4)

Then \( \eta = C_1 W_b - C_2 \). For \( T_b \) there exists no closed form, instead we have the relation:

\[ h(C_1 W_b - C_2, T_b) = \frac{C_1}{C_1 W_b - C_2}. \] (5)

From \( T_b \) and the EoS one can compute \( c_{sb} \) and thus the Mach number \( M_b \). Figure 1 shows the dependence of \( T_b \) and \( \eta \) on \( W_b \) and \( R_b \) in the case of leptonic jets. For large values of \( W_b \) there is only a small range of values of \( R_b \) around 0.35 \( \text{kpc} \) where solutions exist. In fact, the asymptotic value of the beam radius

\[ R_b^\infty = \sqrt{\frac{L_{\text{kin}}}{\rho_m \pi}} \cdot \frac{1 - v_j^{1d}}{v_j^{1d}} \] (6)

is obtained when \( W_b \to \infty \) and \( T_b = 0 \).

Computing \( R_b \) from (6) using the independently measured values of \( L_{\text{kin}}, v_j^{1d} \) and \( \rho_m \) from observations yields \( R_b \approx 0.35 \text{kpc} \), which is in agreement with standard values for kpc scale jets (e.g., Ferrari 1998). Fixing the values of \( v_j^{1d} \) and \( L_{\text{kin}} \), and for a constant large \( W_b \) (i.e., \( v_b \approx c \)), it turns out that

\[ k := \eta h_b = \text{const} \] (7)

and

\[ (kW_b - \eta) R_b^2 = \text{const}. \] (8)

The latter equation implies that for a given value of \( W_b \), \( \eta \) must increase with increasing \( R_b \). As \( k = \text{const} \) this means \( h_b \) and thus \( T_b \) must decrease. Eventually, if \( R_b \) increases, \( T_b \) becomes negative. This explains the physically forbidden area to the right of \( R_b^\infty \) in Fig. 1. A similar argument holds for the forbidden area to the left: decreasing \( R_b \) eventually leads to non-physical solutions with \( \eta < 0 \). Figure 1 also shows that for a given \( W_b \) there exists a maximum allowed \( \eta \) and for a given \( \eta \) there exists a maximum \( W_b \). This maximum \( W_b \) grows increasing the value of \( v_j^{1d} \). Let us also point out that from Eq. (4), increasing values of \( L_{\text{kin}} \) lead to larger values of \( R_b \) (notice the quadratic dependence). Hence, although there exits no solution with the seemingly reasonable values \( W_b = 10, \eta = 10^{-3} \) and \( R_b \approx 0.35 \text{kpc} \) (for the chosen parameters of the external medium, \( v_j^{1d} \), and \( L_{\text{kin}} \)), one can obtain models with these characteristics by increasing \( v_j^{1d} \) and \( L_{\text{kin}} \) up to 0.24 and \( 1.6 \cdot 10^{46} \text{erg/s} \), respectively.
Slightly smaller values of $W_b$ have been chosen to keep our parameters the closest possible to the observed values.

Most of the previous jet simulations assumed that the jets are pressure matched ($K = 1$), because (i) there are no direct measurements of $K$, and (ii) this reduces the parameter space. However, fixing $K$ implies that the properties of the external medium depend on the choice of the jet parameters. Therefore, one has to adjust the temperature of the external medium of every jet model such that $K = 1$. As we wanted to simulate typical FRII-jets in a typical environment, the properties of the external medium are fixed and the assumption $K = 1$ is abandoned. For $K > 1$ jets can still remain stable, because they are pressure confined by the cocoon, where the pressure is much larger than in the external medium. In the simulated models $K$ is in the range $1 - 200$. Under-pressured jets ($K < 1$) cause considerable numerical problems when they are evolved beyond $3 \cdot 10^6$ y, because the amount of ambient gas entrained into the beam becomes large. As the beams of our models are under-dense with respect to the external medium, dense blobs of matter from the ambient medium begin to pile up in front of the nozzle blocking the inflow. As we assume axisymmetry the material piles up around the jet axis.

### 2.3 Numerical models

The simulations are performed with the relativistic hydrodynamic, high-resolution shock-capturing code of MMI98. However, instead of a constant *gamma*-law EoS we have used the EoS described in Sect. 2.1. The code is well suited to solve the equations of relativistic hydrodynamics in cylindrical coordinates, i.e., for problems with axial symmetry. It has been extensively tested and applied previously (e.g., Martí et al. 1995, 1997, MMI98). Covering a time span of up to $6.6 \cdot 10^6$ y the set of models represent the longest simulations of relativistic jets performed so far.

We have considered three models (Table 1) including two leptonic ($X_{lb} = 1$) ones and one baryonic model ($X_{lb} = 10^{-3}; BC$). One of the leptonic models has a factor of 100 lower density $\eta = 10^{-5}$ (LH) than the other two models with $\eta = 10^{-3}$. The corresponding number densities of thermal electrons in the beam are $n_e \simeq 2 \cdot 10^{-5}$ and $n_e \simeq 2 \cdot 10^{-3}$, respectively. Note that both values are below the upper limit estimated from measurements of the degree of polarisation (see Sect. 2.2). The low $\eta$ leptonic model has been run in order
Table 1. Parameters of the three simulated jet models.

| Model: | BC | LC | LH |
|--------|----|----|----|
| $L_{\text{kin}}$ [erg/sec] | $10^{46}$ |   |  |
| $v_{1d}^j$ [c] | 0.2 |   |  |
| $X_{ib}$ | $10^{-3}$ | 1 | 1 |
| $\eta$ | $10^{-3}$ | $10^{-3}$ | $10^{-5}$ |
| $\varepsilon_b$ [$c^2$] | $5.69 \cdot 10^{-3}$ | 0.298 | 119 |
| $T_b$ [K] | $10^{10}$ | $10^9$ | $2.37 \cdot 10^{11}$ |
| $M_b$ | 16.4 | 2.38 | 1.71 |
| $W_b$ | 7.95 | 6.62 | 6.62 |
| $K$ | 1.4 | 91 | 217 |
| $\gamma_b$ | 1.42 | 1.50 | 1.33337 |
| $R_b$ [kpc] | 0.366 | 0.361 | 0.342 |

to check how extremely low values of $\eta$ affect the long term evolution of a relativistic jet (see, e.g., Birkinshaw 1991; Ferrari 1998).

All models have the same power and initial propagation velocity, and they all propagate into the same homogeneous external medium with $\rho_m = 10^{-3} m_p/\text{cm}^3$, $T_m = 10^7 \text{K}$, and $X_{lm} = m_e/(m_e + m_p) \approx 1/1837$. As the jet power is fixed, the lower the density the higher is the internal energy, i.e., the light jet is also the hottest ($\varepsilon \gg 1$). Note also that although the values of the injection temperatures (Table 1) seem to be rather large, they are unanimously determined once the fiducial conditions in the external medium, $L_{\text{kin}}, v_{1d}^j$ and $X_{ib}$ have been fixed.

We have not considered a light baryonic model (the counterpart of the light leptonic model) for three reasons. First, the temperatures within the beam and in some parts of the cocoon would reach $10^{14} \text{K}$. Thus, the dynamics of the model would be very similar to that of the hot leptonic model, i.e., radiation dominated. Second, at temperatures of $10^{14} \text{K}$ (i.e., Lorentz factors of $\approx 10^4$) leptons of the thermal plasma will contribute significantly to the synchrotron radiation power. Assuming an equipartition magnetic field ($\approx 60 \mu \text{Gauss}$ for our simulations; see Sect. 5.4), the synchrotron cooling time is $t_{\text{cool}} \approx 10^4 - 10^5 \text{y}$. This time is much smaller than the typical lifetime of the source, i.e., one has to consider synchrotron cooling in the simulations. Third, the light baryonic model is too hot. Its hot thermal plasma would have a detectable emissivity which is not observed in real sources, so far (see, e.g., Celotti et al.1998).

In order to track the evolution of the beam material the beam mass fraction $X = \rho_b/\rho$ is evolved by our code by means of an additional continuity equation. We also have to evolve the lepton fraction $X_l$ because, as explained in Sect. 2.1, the composition may change due to mixing of species within the flow.
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Table 2. Length and time units for each model. The third and fourth row give the total duration of simulation $t_{\text{max}}$ in code and physical units, respectively. In the last row the time $t_{\text{comp}} = 6.3 \cdot 10^6$ y is given in code units for each model.

| Jet | BC | LC | LH |
|-----|----|----|----|
| $R_b$ [kpc] | 0.366 | 0.361 | 0.342 |
| $R_b/c$ [y] | 1192 | 1176 | 1114 |
| $t_{\text{max}} [R_b/c]$ | 5300 | 5400 | 5950 |
| $t_{\text{max}}$ [y] | $6.3 \cdot 10^6$ | $6.4 \cdot 10^6$ | $6.6 \cdot 10^6$ |
| $t_{\text{comp}}[R_b/c]$ | 5300 | 5350 | 5650 |

Our computational domain spans a region of size $200R_b \times 500R_b$ (approximately 70 kpc $\times$ 175 kpc) in cylindrical coordinates $(r, z)$ with an uniform grid whose resolution is 6 cells per beam radius. This resolution is a compromise between a reasonable computing time per model (about 185 hours on a NEC SX-5 vector computer, running with a sustained performance of about 2.3 Gigaflops) and the maximum evolutionary time to be reached in a simulation (see Sect. 5.5). Axisymmetry is assumed along the $r = 0$ boundary, while the downwind boundaries at $r = 200R_b$ and $z = 500R_b$ are outflow boundaries. The boundary at $z = 0$ is reflecting.

3 RESULTS

3.1 Morphology and dynamics

Due to the selected model parameters the beam radius and, hence, the time units are slightly different for each model. They are given in the first and second row of Table 2. The maximum evolutionary time reached in each model is listed in rows three (in code units) and four (in years), respectively. The last row gives the evolutionary time (in code units) at which all models are compared. This time corresponds to the final evolutionary time of the shortest simulation (model LH) and is equal to $6.3 \cdot 10^6$ y.

Colour coded snapshots of the distributions of density, temperature and Lorentz factor, nearly at the end of the simulations, are displayed in Figs. 2–4. Additional contour lines mark the boundaries of the cocoon (see Sect. 3.3). At first glance, the morphology of all models is very similar, particularly concerning the overall cavity and beam shapes.

Figure 5 displays profiles of several variables along the symmetry axis. One notices that the differences among models in quantities which are most important for the dynamics ($\rho$, $\varepsilon$ and $P$) are small, at least far from the head of the jet. The main differences occur where the profiles intersect the biconical shocks and rarefactions within the beam which are quite different (in spatial distribution and strength) from model to model.
The differences in the adiabatic index in the cavity (from model to model) are the result of the different initial beam temperatures and compositions. The leptonic models reach lower values of the adiabatic index in the cavity than the baryonic one (i.e., they have $\gamma$ closer to $4/3$) because electrons and positrons become relativistic at lower temperatures than protons. Although model LC has an average temperature within the cavity (see Fig. 3) more than ten times smaller than in model BC, as a result of the mixing with the ambient medium, the effective adiabatic index in the cavity is smaller for model LC than for model BC. The reason being that model LC supplies with a thousand times more relativistic particles the cavity than model BC. The beam of model LH is the hottest and, thus, it has the lowest adiabatic index. It is even lower than model LC because the leptons of model LH are more relativistic (i.e., the beam temperature is higher).

At the jet head different phases of vortex shedding can be seen in Fig. 2. In model LC new vortices are forming, in model BC they have been just shed from the head, and in model LH the Mach disk has been replaced by a conical shock causing an acceleration of the head. However, these differences are present only temporarily. More important are the differences inside the beam and in the cocoon which we discuss next.

3.2 Beam

In the hot, light model LH the density in the beam is two orders of magnitude lower and the internal energy is, at least, two orders of magnitude larger than in the two colder, denser models BC and LC. However, model LH has a strongly pinched structure (at $z = 39R_b$ and $z = 63R_b$ in Fig. 4) while the colder models possess a continuous beam where the beam mass fraction is larger than 0.95 throughout the beam. The beam pinching, even disruption, is due to considerable mass entrainment in model LH, inside which one can find regions which are almost at rest. Hence, the light, hot jet appears to be the most unstable one.

The thermodynamic quantities change drastically at the first conical shock in model BC. The temperature rises from $10^{10}$K to $10^{12}$K, the pressure (initially model BC is pressure matched, i.e., $K \approx 1$) increases by two orders of magnitude and stays at values similar to those of the other models from there on ($K = \mathcal{O}(100)$). The internal energy and sound speed also strongly increase, while the relativistic Mach number drops from 130 to 20. All these changes are caused by the initially very low internal energy of model BC (see Tab. 1). The other two models initially have a much larger internal energy and the thermodynamic
quantities in the beam change less during the evolution. Actually, the properties of the recollimation shocks are strongly varying from model to model. On average, recollimations shocks have larger compression ratios and are more numerous in model LH than in the colder models (see Fig. 5a).

The sound speed approaches its maximum value \( c_s \approx 0.57 \) close to the injection nozzle in model LH. In the colder models \( c_s \) increases within the beam and the maximum value is reached at the hot spot. The deceleration (acceleration) at internal shocks (rarefactions) within the beam is much more violent in model LH (where Lorentz factors of up to 20 can be observed) than in the two colder models.

At the beam/cocoon interface a thin, hot layer forms in the cold models, which is thinner than that found in 3D simulations (Aloy et al.1999, 2000). This layer results from the interaction of the beam with the external medium, and appears naturally in some models of jet formation (Sol et al.1989). Its existence has been proposed by different authors (Komissarov 1990; Laing 1996; Laing et al.1999) in order to account for a number of observational characteristics of FRI radio sources. However, the physical nature of the shear layer is still largely unknown, as a study of its properties requires much better resolution and full 3D simulations. The structure of the shear layer is remarkably different in model LH, where the hot layer forms out of the beam/cocoon interface due to the extremely high specific internal energy of this model. Nevertheless, the variations of pressure and Lorentz factor across the shear layer are too small in model LH to have an effect on the emission. This is in contrast to Aloy et al.(2000), where the shear layer was much more extended and its emission properties were distinguishable.

Despite of the differences between models, all of them maintain relativistic beams up to distances of more than 80 kpc the average Lorentz factor being larger than 5. This has important theoretical consequences, in particular, for the most commonly accepted emission models of kiloparsec-scale X-ray jets as we will discuss in Sect 5.

### 3.3 Cocoon

The cocoon surrounding the beam is formed by beam matter that flows backward (relative to the beam) after being deflected at the hot spot. This beam matter partially mixes with the external medium, i.e., the cocoon does not contain pure beam material \((X \neq 1)\). Hence, we define the cocoon as the region containing material with a beam particle fraction \(0.1 < \)
We have used the beam mass fraction as the criterion, because the density within the cavity is similar in all models (see Fig. 2), i.e., the lower the value of $X$, the lower is the number of emitting particles and thus the total emitted intensity. The thresholds are a bit arbitrary, but we have checked that variations of the lower threshold (up to a decade) do not change significantly the shape of the cocoon. For a very small lower threshold ($X_{\text{min}} = 0.001$) the cocoon practically coincides with the complete cavity blown by the jet (see below). Increasing the value of the upper threshold would include too much of the beam itself as part of the cocoon. Other cocoon definitions have been given in, e.g., Cioffi & Blondin (1992).

According to our definition, the shape of the cocoon is remarkably different among the considered models. In model LH the cocoon is restricted to a narrow layer around the beam, while it is much more extended and matches the relativistic backflow region in the denser models (BC and LC). This is due to the lower beam density of model LH, which allows for a larger inertial confinement of the beam flow by the external medium without causing much mass entrainment (which would destroy the jet). Another consequence of the lower density of model LH is a lower beam mass fraction within the cavity outside the cocoon (about two orders of magnitude smaller than in the other models). The average cocoon temperature of models BC and LH is around $10^{11}$K the size of local fluctuations reaching up to two orders of magnitude. In model LC the average cocoon temperature is lower ($10^{10}$K) and there are almost no fluctuations (see Fig. 3). The temperature distribution of models BC and LH is very similar (although they have a different density and composition), whereas that of model LC (with the same density as model BC and the same composition as model LH) is very different. In Sect. 5 we will explain how the varying particle density causes this effect.

A backflow with mildly relativistic velocities occurs in thin shells starting from the head of the jet and limiting the cocoon in models BC and LH. The amount of beam material in this backflow is smaller in model LC than in the other models (see, e.g., Fig. 4).

4 EVOLUTION

4.1 Definitions

For the discussion of the jet evolution we introduce the following quantities: (1) the jet length $l_j$ being the z-coordinate of the contact discontinuity; (2) the cocoon length $l_{cc}$ being the difference between $l_j$ and the smallest z-coordinate of the fluid belonging to the cocoon (according to the definition given in Sect. 3.3); (3) the cavity length $l_c$ being the z-coordinate
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of the bow shock on the jet axis; (4) the average cocoon radius \( r_j \); (5) the average cavity radius \( r_c \); (6) the cocoon aspect ratio \( A_j = l_{cc}/r_j \); (7) the cavity aspect ratio \( A_c = l_c/r_c \); (8) the hotspot pressure \( P_{hs} \) being the average pressure in the 10 cells upstream of \( z = l_j \); (9) the average cavity pressure \( P_c \); (10) the jet velocity \( v_j \) being the velocity of the contact discontinuity between the jet and the external medium; and (11) the bow shock velocity \( v_c \) on the jet axis.

The position of the contact discontinuity \( (l_j) \) is always near the bow shock \( (l_c, \text{ see Fig. 6a}) \). But whereas \( l_c(t) \) is increasing monotonically, \( l_j(t) \) looks rather ‘noisy’ due to the vortex shedding mechanism, which makes it difficult to detect (numerically) the position of the contact discontinuity. The detection algorithm searches for the first cell on the jet axis (beginning at \( z = 0 \)) where \( X < 0.5 \). When a vortex sheds off, the distance of this cell from the inlet is abruptly reduced. As \( A_j \) and \( v_j \) depend on \( l_j \), they display an oscillatory behaviour, too (Figs. 6b and 6d, respectively). Other quantities like \( r_j \) and \( P_{hs} \) do also oscillate because of the highly dynamical processes that affect them (vortex shedding and mixing with the external medium), in particular when the jets enter their second evolutionary phase (see Sect. 4.2). \( P_{hs} \) and \( v_j \) have been smoothed (averaged over 25 and 50 cells, respectively) in Figs. 6a and 6c, but their almost constant values during the first 100 time units are not caused by this smoothing process.

4.2 Evolutionary phases

In all models two evolutionary epochs can be distinguished. During the first epoch \( (t < t^{1d} \approx 1.2 \cdot 10^5 \text{y}; \text{ see Fig. 6}) \) multidimensional effects are not yet important and the jet propagates with a velocity close to the 1D estimate (1). The Mach disk is only slightly disturbed, the shedding of vortices is negligible, and both the hotspot pressure and the jet velocity are nearly constant (Figs. 6c, 6d). The aspect ratios of the cocoon and of the cavity (Fig. 6b) increase with time, while the average cavity pressure decreases (Fig. 6c).

The second epoch starts when the first larger vortices are produced near the jet head. The beam cross section increases and the jet decelerates progressively. The deceleration is repeatedly interrupted by short acceleration phases caused by the temporary replacement of the Mach disk by a conical shock, which occurs e.g., in model BC at \( t \approx 400, 1200 \) and \( 5100 \), respectively (Fig. 7). The velocity increase during the acceleration phases can exceed 100%.
Table 3. Cavity lengths (in kpc) and aspect ratios at the end of the one-dimensional epoch \((t = 1.1 \cdot 10^5 \, \text{y};\) first two rows) and near the end of the simulations \((t = 6.0 \cdot 10^6 \, \text{y};\) third and fourth rows) for all models. The last row shows the time averaged velocity during the time interval \([1.1 \cdot 10^5, 6 \cdot 10^6]\) (in units of \(c\)).

| Phase | Jet | BC | LC | LH |
|-------|-----|----|----|----|
| 1D    | \(l_c\) | 7.57 | 7.10 | 6.83 |
|       | \(A_c\) | 3.37 | 3.09 | 2.89 |
| 2D    | \(l_c\) | 106.8 | 100.6 | 97.4 |
|       | \(A_c\) | 4.32 | 4.03 | 3.79 |
|       | \(\tau\) | 0.055 | 0.052 | 0.050 |

The acceleration phases also give rise to a higher hotspot pressure. Such pressure can be estimated assuming that the pressure of the external medium is negligible. Then the ram pressure of the external medium \(\rho_m v_c^2\) is equal to the thermal pressure behind the bow shock, and hence about equal to the hotspot pressure as the pressure is continuous across the contact discontinuity (Fig. 8). There is a phase shift between \(P_{hs}\) and \(\rho_m v_c^2\), because pressure fluctuations are generated at the contact discontinuity (vortex shedding) which arrive at the bow shock later.

5 DISCUSSION

In spite of an extremely different composition and internal energy, the differences in morphology and dynamics of the jets are small (particularly, between models BC and LC, both having the same density contrast). This is unexpected, as the jets of our models have a fixed power which is carried by particles of different mass, i.e., the energies per particle are quite different. A leptonic jet has \(\approx 10^3\) times more particles than a baryonic jet of the same mass density. Therefore, its kinetic energy per particle is smaller, i.e., its temperature is smaller, but its specific internal energy (erg/g) is larger. We also expected that differences in the adiabatic index of the models would manifest itself.

It turns out, however, that fixing the power and initial speed of the jet as well as the properties of the ambient medium, the development of the jet seems to be quite well defined (at least for the jet parameters considered and for the evolution time covered by our simulations). However, a more careful inspection of the results reveals that some remarkable differences exist between our models which we discuss in the following.
5.1 Comparison with the extended Begelman-Cioffi model

A simple theory of the evolution of the cocoon of extragalactic jets was presented in BC89. The authors argue that the cocoon has not yet reached pressure equilibrium with the surrounding medium in many sources, and that high pressure confines the jet keeping it highly collimated as seen in extragalactic sources. BC89 assume that both the bow shock velocity $v_c$ and the power transported into the cavity $L_c$ are not time-dependent, and that the pressure of the external medium is negligible. Then the average cavity pressure is given by

$$P_c = \frac{(\gamma_c - 1)L_c}{v_c F_c},$$

where $F_c = \pi r_c^2$ is the cross section of the cavity and $\gamma_c = c_p/c_v$ is the constant average adiabatic index in the cavity. The cavity pressure causes an expansion of the cavity with a velocity $\dot{r}_c$, i.e.,

$$P_c = \rho_m \dot{r}_c^2.$$  \hspace{1cm} (10)

This implies $1/r_c \propto \dot{r}_c$ and hence

$$r_c \propto t^{1/2}, \quad F_c \propto t, \quad P_c \propto t^{-1}, \quad l_j/r_c \propto t^{1/2}. \hspace{1cm} (11)$$

We have extended the model of BC89 (eBC, hereafter) replacing the assumption $v_c = \text{const}$ by the more realistic one $v_c \propto t^\alpha$ and further assume that $\dot{r}_c \propto t^\beta$. Then (9) together with (10) yields

$$\frac{1}{t^{\alpha + 2(\beta + 1)}} \propto t^{2\beta} \quad \Rightarrow \quad \beta = -1/2 - \alpha/4, \hspace{1cm} (12)$$

and thus

$$r_c \propto t^{1/2 - \alpha/4}, \quad F_c \propto t^{1-\alpha/2}, \quad P_c \propto t^{-1-\alpha/2}, \quad l_j/r_c \propto t^{3/2 + 5\alpha/4}. \hspace{1cm} (13)$$

The only free parameter of the eBC model is the value of $\alpha$. A deceleration ($\alpha < 0$) leads to a faster radial expansion, a slower decrease of the cocoon pressure and to a slower increase of the aspect ratio. For $\alpha = -2/5$ the cavity evolves self-similarly (i.e., with constant aspect ratio).

We have extracted the parameter $\alpha$ and the exponential dependence of $P_c$ and $l_j/r_c$ from our simulation results by fitting the position of the bow shock as a function of time with a power law (Tab. 4). Using the fitted values of $\alpha$ we have also computed the exponents for $P_c^{\text{BC}}$ and $(l_j/r_c)^{\text{eBC}}$ according to the extended BC89 model. The fitting procedure is not very reliable for the 1D epoch as it involves only a few data points, while velocity fluctuations cause problems for fitting the 2D epoch. Thus, differences between the models
Comparison of the simulation results with the extended BC89 model. Columns one and two give the evolutionary epoch and the model, respectively, and column three shows the value of the free parameter $\alpha$ obtained from the simulations by fitting the position of the bow shock as a function of time with a power law. The fourth column provides the fitted exponential time dependence of the cavity pressure, while the fifth column gives corresponding exponent of the eBC model ($= -1 - \alpha/2$). Columns 6 and 7 show the exponent describing the time dependence of the aspect ratio obtained from the simulations and the eBC model($1/2 + 5\alpha/4$), respectively.

| Phase | Jet | $\alpha$ | $P_c$ | $P_c^{eBC}$ | $l_j/r_c$ | $(l_j/r_c)^{eBC}$ |
|-------|-----|---------|-------|------------|-----------|------------------|
| 1D    | BC  | -0.113  | -0.945| -0.944     | 0.449     | 0.358            |
|       | LH  | -0.219  | -0.906| -0.891     | 0.358     | 0.226            |
| 2D    | BC  | -0.355  | -0.722| -0.823     | 0.058     | 0.056            |
|       | LH  | -0.363  | -0.779| -0.818     | 0.059     | 0.045            |

The values of $P_c^{eBC}$ and $(l_j/r_c)^{eBC}$ obtained from the simulations and from $\alpha$ via the eBC model agree reasonably well. The exponent for the aspect ratio is slightly larger than that predicted by the eBC model (especially for the 1D epoch), i.e., the radial expansion of the cavity is slower than expected from the deceleration of the jet head. This discrepancy arises from the assumption of the BC89 model that the whole power of the jet is used to increase the pressure behind the bow shock (which is responsible for the sideways expansion of the cavity). However, part of the jet power is used to fuel the relativistic backflow and its kinetic energy is not available to expand the cavity in radial direction, i.e., the aspect ratio (jet length over cavity radius) grows faster in the simulations than in the BC89 model.

During the 2D epoch the cavity evolution is almost self-similar and the cavity aspect ratio remains nearly constant. This situation will only change when the cavity pressure becomes equal to the pressure of the external medium and the radial expansion ends. Extrapolating from Fig. 6 this will happen after $\sim 10^6$ time units or $10^9$ y, which exceeds the estimated ages of observed jets ($10^7$ y – $10^8$ y) by one order of magnitude.

The evolution of the jets proceeds in two epochs independent of the composition of the jet. This is in agreement with previous results obtained by MMI98 and Komissarov & Falle (1998). The 1D epoch of our models corresponds to the intermediate phase of Komissarov & Falle (1998), while the second epoch may be identified with their self-similar phase. They also find that the growth rate of the aspect ratio is smaller during the second epoch. However, their growth rates are slightly larger than ours in both epochs. This discrepancy is probably due to different jet parameters. Their relativistic models are 30 times denser than ours, have

are probably not significant. Note in this respect that although the value of $\alpha$ varies by a factor of two between the models, the variation of $P_c^{eBC}$ is at most 5% and that of $(l_j/r_c)^{eBC}$ less than 25% as both quantities can be extracted more reliably. The values of $P_c^{eBC}$ and $(l_j/r_c)^{eBC}$ obtained from the simulations and from $\alpha$ via the eBC model agree reasonably well. The exponent for the aspect ratio is slightly larger than that predicted by the eBC model (especially for the 1D epoch), i.e., the radial expansion of the cavity is slower than expected from the deceleration of the jet head. This discrepancy arises from the assumption of the BC89 model that the whole power of the jet is used to increase the pressure behind the bow shock (which is responsible for the sideways expansion of the cavity). However, part of the jet power is used to fuel the relativistic backflow and its kinetic energy is not available to expand the cavity in radial direction, i.e., the aspect ratio (jet length over cavity radius) grows faster in the simulations than in the BC89 model.

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a little bit smaller Lorentz factor, are pressure matched, and possess a finite opening angle. The transition to the second epoch does not occur in those models of Komissarov & Falle (1998) which have an opening angle larger than 10°.

5.2 Evolutionary differences

The average jet velocities during the 2D epoch are listed in Tab. 3 together with the lengths and the cavity aspect ratios at the end of the 1D epoch and at the end of the simulation. At first glance there seems to be a clear trend in the 1D epoch: with increasing internal energy and decreasing Mach number the jets become slower and wider (the mass flux into the cocoon increases). As already found by MMI98, the jet velocities are noticeably larger than predicted by the 1D estimates \( v_j^{1d} = 0.2c \) at the beginning of the simulation (see Figs. 6d and 7). This discrepancy is caused by imposing a reflecting boundary condition at \( z = 0 \) (see Sect. 2.3) and by the different pressure ratios \( K \). After some ten time units the jet head has propagated far enough from the grid boundary to be no longer affected by it. The average cavity pressure decreases and the pressure contrast between the jet and the cavity becomes closer to one in all the models (Fig. 6). After the jets has decelerated at the end of the 1D epoch their velocities and the 1D estimates agree very well (Fig. 7).

During the 2D epoch the jet of model BC propagates considerably faster than those of models LC and LH, while the jet of model LC is only slightly faster than that of model LH. This result can be easily understood. The relativistic Mach number of model BC is larger, i.e., the angle between the conical shocks inside the beam and the beam axis is smaller. Hence, these shocks are less efficient in decelerating the flow. This also affects the propagation speed of the head, because the terminal Mach disk is temporarily replaced by conical shocks an effect pointed out already by Martí et al. (1997). The noisy bow shock velocity (Fig. 7) is also a result of this process. However, as the heads of all jet models are hot and the Mach number is low in this region, this effect does not lead to very strong differences between the computed models.

The differences in the cavity aspect ratios correlate with the average propagation speeds. The jet of model BC is the fastest and the longest one (Tab. 3). The cavity radius is similar for the cold models and bigger for model LH. The models expand radially almost at the same rate, because the average cavity pressure is essentially equal for all three models although a significant part of the kinetic energy which is contained in the relativistic backflow cannot
be used to inflate the cavity. In fact, we included model LH in our investigation in order to check whether an increased fraction of internal (non-directed) energy would lead to less kinetic (directed) energy in the cocoon.

5.3 The influence of different compositions

From a numerical point of view the most important effect of a varying composition is the non-constant adiabatic index requiring an extension of the original MMI98 code (see Sect. 2.3). The composition dependent $\gamma$ calculated from the Synge EoS (see Sect. 2.1) however, does not lead to large differences in the evolution and morphology of the jets. The average $\gamma$ in the cocoon is around 1.4 and $\gamma = 5/3$ in the external medium for all three models. Due to the initial conditions, the adiabatic index within the beam is different for every model, but these differences have no observational relevance as $\gamma$ is not an observable quantity.

We expected that the largest internal energy content in the leptonic models would cause a larger expansion of the beam than in the baryonic model, leading to a faster deceleration of the head and differences in the aspect ratio of the cavity (more spherical in the leptonic models). However, this effect is compensated by the large overpressure of the cavity that efficiently confines the jet laterally.

The variation of the particle density with the composition gives rise to huge differences in the temperature. At the same density, an $e^\pm$-gas contains three orders of magnitude more particles than an ionised hydrogen gas. The particle density is proportional to $\eta \cdot X_{lb}$ and, therefore, it is much higher in model LC than in the other two models. Hence, the same amount of energy is distributed over many more beam particles in model LC than in models BC and LH, leading to an about three orders of magnitude lower beam temperature (Fig. 3). However, this does neither explain the flat temperature profile of model LC, nor the identical average cavity temperatures of models LH and BC (whose hotspot temperatures differ by a factor of ten). To explain the temperature differences in the cocoon, particles from the external medium must be considered too. There are $N_{cm} = 7 \cdot 10^{64}$ particles from the external medium and $N_{cb} = 3 \cdot 10^{65}$ particles from the beam in the cavity of model LC at $t = 6.0 \cdot 10^6$y. The dominant fraction of the internal energy of the cavity is also carried by beam particles. Hence, the mixing of beam and external medium has little impact on the internal energy per particle, i.e., the temperature remains at the hotspot value.

Concerning models LH and BC, particles from the external medium dominate the particle
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density in the region of the cavity excluding the dense shell behind the bow shock (cavity interior), i.e., \(N_{c,b}/N_{c,m} < 1\), but the internal energy is still dominated by beam particles. This leads to large changes and fluctuations of the temperature during the mixing of particles from the beam and the external medium (Fig. 3). The temperatures near the jet head and near the inlet are very different in these two models. Although the hotspot temperatures differ, the average cocoon temperatures are similar because the same amount of energy is distributed over a similar number of particles. In addition, since the particle density is much lower than in model LC, the average cocoon temperature is much higher.

In general, when \(N_{c,b}/N_{c,m} > 1\) the cavity interior has a flat temperature profile and we talk of an isothermal cavity. When \(N_{c,b}/N_{c,m} < 1\) the average temperature in the cavity interior is heterogeneous and independent of the beam composition and density. In this case we talk of a non-isothermal cavity. Another effect is caused by the different time dependencies of \(N_{c,b}\) and \(N_{c,m}\). The number of particles from the external medium in the cavity interior is proportional to its volume, and the cavity interior has a roughly constant density during the whole evolution (e.g., it changes by less than 20% during \(6.0 \cdot 10^6\) yr in case of model LH). Hence, from the eBC model one can derive \(N_{c,m} \propto t^{2+\beta} \cdot t^{1+\alpha} = t^{2+\alpha/2}\). While the number of beam particles increases in the cavity interior linearly with time \(N_{c,b} \propto t\), the ratio \(N_{c,b}/N_{c,m} \propto t^{-1-\alpha/2}\) decreases, i.e., a transition from an isothermal to a non-isothermal cavity may occur during the jet evolution.

Figure 9 shows the evolution of \(N_{c,b}/N_{c,m}\) for the three models. Model LC has an isothermal cavity throughout the simulation, while the cavities of the other two models are always non-isothermal. Although, we have not found the above mentioned transition, Fig. 9 suggests that one should be able to find initial jet parameters (between those of models LH and LC) where such a transition will occur during the simulation.

In an isothermal cavity the external particles can be neglected, i.e., \(N_c \approx N_{c,b} \propto t\). As the cavity volume is proportional to \(t^{2+\alpha/2}\), the number density within the cavity \(n_c \propto t^{-1-\alpha/2}\). However, as \(P_c \propto t^{-1-\alpha/2}\), this implies that the temperature \(T_c \propto P_c/n_c\) does not depend on time. This explains why the cocoon temperature is constant in \(z\)-direction and also in \(r\)-direction (Fig. 3). In a non-isothermal cavity particles from the external medium dominate, i.e., \(n_c = \text{const}\) and the average cocoon temperature decreases according to \(T_c \propto t^{-1-\alpha/2}\).
5.4 Comparison with observations

Comparing the hydrodynamic properties at the end of our simulations with those deduced from observations of Cyg A, we find that the hot spot pressure ($P_{hs} \approx 0.001 \rho_m c^2 \approx 1.5 \cdot 10^{-9} \text{ dyn cm}^{-2}$) agrees within a factor of 2 with the value reported by Carilli et al. (1996).

Comparing the hotspot advance speeds of CSO and FR II sources implies a deceleration of the hotspot propagation at early epochs. This deceleration indeed occurs in our simulations. The propagation velocity after the deceleration ($v_j \approx 0.05 c$) is in good agreement with the estimates of Daly (1995). The equipartition magnetic field of $\approx 60 \mu\text{G}$ in the cavity and $\approx 600 \mu\text{G}$ in the hotspot are larger than those reported for Cyg A (Carilli et al.1996), although they agree with the values obtained for the radio lobes and hotspots of younger powerful radio sources (Ferrari 1998). This agreement is easily explained considering that the evolutionary time covered by the simulations is relatively short.

Up to now it is impossible to determine hydrodynamic properties of jets directly from the observed radio emission. Thus, in order to compare the simulated jets with observed ones we proceed in ‘opposite’ direction and compute the non-thermal synchrotron radio emission of our hydrodynamic models. We follow the same approach as in Aloy et al. (2000) assuming that the energy density weighted with the beam particle fraction is distributed among the emitting non-thermal particles according to a power law $N(E) \propto E^{-\sigma}$ (without cut-offs). The magnetic field is considered to have equipartition strength and a ad hoc structure: aligned with the flow velocity and with a negligible random component (because we are not concerned in polarisation properties). The simulated radio maps obtained from the hydrodynamic models (Fig. 10) include only those computational zones where the density is smaller than $0.1 \rho_m$ to exclude emission from the external medium. We have chosen a viewing angle of $45^\circ$ to compute the synthetic radio maps in order to avoid excessive Doppler boosting of the beam which would out-shine completely the diffuse emission from the cocoon.

Two main features known from observations are present in all models: radio lobes with hotspots and a one-sided, knotty jet (particularly evident in model LH). The knots are associated with internal shocks in the beam. Emission is dominated by the hotspots, where the pressure is maximum. Figure 10 shows that the models have different jet/cocoon emission ratios, as expected from their different cocoon morphology (Sect. 3.3). While models BC and LC show some emission from the cocoon, resembling the observed lobes, the emission is dominated by the beam in model LH.
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In actual sources the lobes of the jet and counter-jet are often observed to be equally prominent. Our simulations show larger lobes (i.e., more cocoon emission) for the counter-jet (particularly for models BC and LC). This can be explained by the relativistic backflow present in the simulated models, because Doppler boosting enhances the emission from the backflow regions in the counter-jet and dims that of the approaching jet lobe. The relativistic backflow also limits the radial expansion of the lobes and prevents the beam from inflating large lobes with gas moving at sub-relativistic speeds. Therefore, the similarity of the lobes in actual radio sources suggests that there are no relativistic backflow regions. Their presence in the simulated models is most likely due to the axial symmetry imposed in our 2D simulations. Three dimensional simulations (Aloy et al.1999) show much smaller backflow velocities, typically $\sim 0.25c$, and hence do not lead to substantial Doppler beaming in the counter-jet cocoon.

Very recently, the new generation of space-based X-ray observatories has allowed for the detection of X-ray emission from kiloparsec scale jets (Schwarz et al.2000; Chartas et al.2000; Sambruna et al.2001). The observed X-spectra are not satisfactorily explained by standard radiation mechanisms, i.e., synchrotron self-Compton (Tavecchio, Maraschi & Ghisellini 1998; Schwarz et al.2000; Tavecchio et al.2000) and synchrotron radiation. The latter can arise from the same population of particles that produce the radio emission (Schwarz et al.2000; Marshall et al.2001), or from a second, much more energetic population of electrons co-spatial with the one responsible for the radio emission (Röser et al.2000; Sambruna et al.2001). Alternatively, inverse Compton scattering of beamed photons of the cosmic microwave background (CMB) has also been proposed (Tavecchio et al.2000; Celotti et al.2001). For this mechanism to work, it is necessary that the beam of the jet is, at least, mildly relativistic, because the amplification of the CMB radiation increases with the Doppler factor, $\delta \equiv [\Gamma (1 - \beta \cos \theta)]^{-1}$ (Tavecchio et al.2000). Our models show relativistic beam velocities ($\Gamma \approx 5$) out to 80 kpc (see Sect. 3.2). Whether this is a result of the imposed axisymmetry of our models or a physical feature may only be disentangled by performing three-dimensional simulations (see Sect 5.5).

Thermal bremsstrahlung from the gas confining the radio-optical jet was very soon disregarded as the source of the observed X-ray radiation from the jets of powerful sources (Sambruna et al.2001; Schwarz et al.2000), although the data can be fitted by a thermal bremsstrahlung spectrum (Chartas et al.2000). In order to explain the observations invoking thermal bremsstrahlung, a large electron density ($n_e \sim 2 \text{ cm}^{-3}$) is required. This is, however,
inconsistent with the observed (relatively low) rotation measure \((n_e < 3.7 \times 10^{-5}/(B L)); \ B \) and \(L \) are the magnetic field strength and the path length, respectively; Schwarz et al. 2000). The estimate for the electron number density, is based on the assumption that (i) there are no magnetic field reversals within the VLA beam, and (ii) the confining thermal material is not relativistic and relatively cold. But, our simulations show the presence of a very hot relativistic plasma (at least in model BC), which is not accurately described by the classical formulas.

Given that the densities are very similar in all our models and that there are huge temperature differences inside the cavity (Fig. 3), we have analysed whether relativistic thermal bremsstrahlung can be used to distinguish between the cavities inflated by jets of different composition. We follow the work of Nozawa, Itoh & Kohyama (1998), which includes the appropriate Elwert (1939) factor to compute the relativistic thermal bremsstrahlung cross section. Their method is accurate if one can neglect the thermal motion of protons, i.e., at temperatures below \(10^{11}\) K. For larger temperatures the relativistic thermal bremsstrahlung cross section is still accurate within an order of magnitude (Itoh, private communication). Above \(10^{9}\) K electron-electron, electron-positron, positron-ion, and positron-positron bremsstrahlung processes are important (see, e.g., Novikov & Thorne 1973). These processes should have been included when computing the total bremsstrahlung emissivity of our models, where temperatures are as high as \(10^{12} - 10^{13}\) K in the shear layer confining the jet, in the hotspot, and in the cocoon of model BC. However, as the emissivity of these processes is comparable with that produced by the interaction of electrons with protons (Dermer 1986; Svenson 1982), the emissivity of the latter process provides a lower bound of the total bremsstrahlung emissivity, which is accurate to an order of magnitude up to temperatures of \(\sim 10^{13}\) K.

The bremsstrahlung power \(P_\nu^{\text{brems}}\) per unit of frequency \((\nu)\) in the comoving frame is proportional to \(n_en_i/\sqrt{T}\) \((n_i\) is the number density of ions in the medium\). Considering a viewing angle of 90°, i.e., beaming effects are unimportant, \(P_\nu^{\text{brems}}\) is dominated in all models at relatively low frequencies \((h\nu\leq 100\ \text{keV})\) by the very dense shell at the cavity boundary behind the bow shock containing shocked external medium. Although the non-shocked external medium is cooler than the shell, \(P_\nu^{\text{brems}}\) is three orders of magnitude smaller than in the shell, because the shocked material is much denser than the ambient medium. The bremsstrahlung emission from the cavity’s interior is very weak in models BC and LH. This is different for model LC, where the temperatures in the cocoon are lower and
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The electron densities are larger than in the other models. The resulting bremsstrahlung emissivity of the cavity’s interior is almost as strong as that of the external medium.

The dominance of the X-ray emission from the shell is also evident from the surface brightness, i.e., from the integral of $P^\text{brems}_\nu$ computed along the line of sight (Fig. 11). The total emission X-ray maps of all models look very similar, as the shell outshines the cavity at a frequency of 10 keV (Fig. 11 top panel). Note that our analysis does not take into account the emission of intra-cluster gas, which would be dominant for a radio source located in a galaxy cluster. The dominance of the shell is in agreement with the work of Heinz, Reynolds & Begelman (1998). They also find that the shell’s dominance decreases as the source evolves because the bow shock decreases in strength when the pressure inside the cavity decreases due to the expansion. McNamara et al. (2000) find a reduced X-ray surface brightness when observing the radio lobes of Hydra A, an effect which is evident in our models. However, McNamara et al. (2000) also point out that there is no evidence for shock-heated gas surrounding the radio lobes, and hence suggest that the cavity expands subsonically. This does not contradict our results, as we consider FRII sources (Hydra A is a FRI source) and because our models represent a still relatively young evolutionary stage of a powerful radio source. FRI sources reach the transonic regime relatively soon and, therefore, one does not expect to find a strong shock surrounding the radio lobes. Instead the radio lobes flare into the external medium. This may explain the lack of X-ray emission from the region surrounding the cavity. On the other hand, it is expected that a further evolution of the source will lead to a decrease of the cavity pressure, and consequently to a weakening of the cavity bow shock (which in turn reduces the X-ray emission of the shell).

We have also computed the surface brightness per unit of frequency at 10 MeV (Fig. 11 bottom panel). We choose this high frequency, which is beyond the observational range of \textit{CHANDRA} ($\approx 0.1 - 10 \text{ keV}$), because only when considering hard X-rays one begins to see the emitting fluid inside the cavity and the emission of the model differs most. Increasing the frequency from 10 keV to $100 \text{ keV} \lesssim \nu \lesssim 1 \text{ MeV}$ the interior of the cavity shows up, although the bow shock surrounding the head of the jet is still dominant. At even higher frequencies ($\gtrsim 1 \text{ MeV}$), the hotspot and the high temperature cocoon are visible. Finally, at $h\nu = 10 \text{ MeV}$ (bottom panel in Fig. 11), only models BC and LH show some emission while the exponential cut-off of the bremsstrahlung spectrum reduces the emissivity of model LC, as the temperature of the cavity interior of model LC, $kT \simeq 500 \text{ keV}$, is much smaller than the considered frequency. This might provide a way to distinguish observationally – once fu-
ture detectors have the appropriate dynamical range and sensitivity – leptonic from baryonic jets (with the same kinetic power and density contrast). If one observes a transition from a shell dominated source to a source dominated by the interior of the cavity when increasing the observation frequency from $h\nu \approx 1$ MeV to $h\nu \approx 10$ MeV, the jet feeding the radio lobe is most probably rich in baryons. Conversely, if this transition is not detected and one finds instead a significant suppression of the emission, the jet is mostly made of leptons.

Typical luminosities of very powerful X-ray, large-scale jets are of the order $10^{45}$ erg sec$^{-1}$, e.g., Chartas et al. (2000) report a X-ray luminosity in the $2 - 10$ keV band of $L_X \approx 6.3 \cdot 10^{45}$ erg sec$^{-1}$ for PKS 0637-752. The X-ray luminosities produced by our jet models due to thermal bremsstrahlung are far below the observed values. This rules out thermal bremsstrahlung as the dominant mechanism for the X-ray emission of large-scale jets. However, several factors may enhance the total thermal luminosity. First, the size of our jets is still relatively small as compared with fully evolved FR II sources (the larger the jet the larger is the luminosity). Second, the kinetic power of our models, although being large, is still below the one inferred e.g., for PKS 0637-752 ($10^{48}$ erg sec$^{-1}$; Tavecchio et al. 2000). Our simulations show that the jet itself radiates only a small fraction of the total power per unit of frequency emitted by the whole cavity ($1.5 \cdot 10^{22}$ erg sec$^{-1}$Hz$^{-1}$ for model BC at 10 keV). At 10 MeV the emitted power from the cavity is $2.7 \cdot 10^{16}$ erg sec$^{-1}$Hz$^{-1}$ and is dominated by the hotspot which contributes $8.0 \cdot 10^{14}$ erg sec$^{-1}$Hz$^{-1}$.

5.5 Caveats

Which are the astrophysical and numerical limitations of our simulations?

We have used a uniform external atmosphere, but extragalactic jets after having crossed a galactic halo will propagate into a much more diffuse intergalactic (or intra-cluster) medium. Both the pressure and the density decrease with the distance from the galactic nucleus. A density declining atmosphere will cause a widening of the jet (and of the cavity), and a substantial deceleration provided that the jets are not inertially confined. The effects of (continuous or abrupt) density changes in the external medium and its impact on the long term evolution of Newtonian jets has been investigated, e.g., by Hooda et al. (Hooda et al. 1994). Such simulations involve additional, not very well determined model parameters and are more complex making them less attractive from the computational point of view, in
particular as our grid resolution (zones per beam radius) is much larger than that in Hooda et al. (1994).

Our simulations are purely hydrodynamic, i.e., they do not consider the effects of magnetic fields. Strong magnetic fields might confine the shocked jet material in an extended *nose cone* preventing it to be continuously deposited into the cocoon. An episodic release of thermal material in the cocoon would have important consequences for the stability of the jet. However, it is quite unlikely that at kiloparsec scales strong magnetic fields do exist. Instead, jets are expected to transport weak and mainly randomly oriented magnetic fields (e.g., Ferrari 1998).

The grid resolution employed (six zones per beam radius at the injection nozzle) is far from being appropriate to account for phenomena like mass entrainment from the external medium. However, it is sufficient to describe the gross morphological features and to capture the average cocoon dynamics. One should consider that, in order to compute up to a maximum physical time, $t_{\text{end}}$, the required run time is proportional to $(\text{cells} / R_b)^3$. Hence, increasing the grid resolution by a factor of two requires about 1200 hours of computing time to reach $t_{\text{end}} \approx 6 \cdot 10^6 \text{y}$, and 9600 hours to simulate a typical lifetime of a FR II source ($t_{\text{end}} \approx 10^7 \text{y}$). In order to check whether the chosen resolution was sufficient to capture, at least the qualitative morphology, we performed a set of shorter simulations $t_{\text{end}} \approx 2.3 \cdot 10^4 \text{y}$ at different grid resolutions (Fig. 12). With four zones/$R_b$ the short acceleration phases caused by the vortex shedding are missing, while they are present with six zones/$R_b$. At even higher resolution the qualitative changes are small.

Finally, we have restricted ourselves to axisymmetric jet flows. This is a considerable restriction, because we force the jet to propagate along the symmetry axis. As, e.g., Aloy et al. (2000) showed, three-dimensional jets may wobble and their effective head area may change appreciably. Consequently, two possibilities arise: (i) the effective area grows (as the result of, e.g., the *dentist drill effect*; Cox, Gull & Scheuer (1991)) leading to a faster deceleration, or (ii) the wobbling of the Mach disk reduces the cross sectional surface of the jet head leading to an acceleration (Aloy et al. 2000). It is difficult to extrapolate the results of Aloy et al. (2000) and to decide whether or not such mechanisms operate in the long run, and what their consequences are. Thus, three-dimensional numerical simulations are required to clarify this issue.

Because of the assumed axisymmetry we encountered numerical problems when simulating under-pressured jets (see Sect. 2.2). These problems would not appear in three-
dimensional simulations as the jet fluid will bypass the blobs of external medium material trying to block the inlet.

6 SUMMARY AND CONCLUSIONS

We have performed the longest and best resolved large-scale simulations of relativistic jets up to date. The simulations extend up to an evolutionary age of $6 \cdot 10^6$ y which is only a few times smaller than the ages of typical powerful radio sources. The simulated jets differ in their density, temperature and composition, but have similar gross properties (kinetic luminosity and thrust). Hence, they may represent a relatively early stage of the observed extragalactic jets and their radio lobes. We have considered three models including a hot and a cold jet made of a pure (leptonic) electron-positron plasma (models LH and LC) and one baryonic model BC consisting of an electron-proton plasma. The leptonic hot model LH has a factor of 100 lower density ($\eta = 10^{-5}$) than the other two models ($\eta = 10^{-3}$).

Although there is a difference of about three orders of magnitude in the temperatures of the cavities inflated by the simulated jets, we find that both the morphology and the dynamic behaviour are almost independent on the assumed composition of the jets. Only the leptonic hot model LH behaves differently because of its very light beam. This reflects the well known influence of the density parameter $\eta$ according to which light jets are less stable than heavy ones. Model LH suffers from considerable mass entrainment into the beam, which eventually leads to the disruption of the beam. The instability grows on a time scale of a few million years, i.e., the evolution of model LH is inconsistent with the sizes and ages of the jets of typical powerful radio sources. In addition, the initial conditions chosen for model LH are quite extreme, and most likely cannot be found in actual jets.

In all models two evolutionary epochs can be distinguished. During the first epoch multidimensional effects are not yet important and the jet propagates with a velocity close to the 1D estimate. The Mach disk is only slightly disturbed, the shedding of vortices is negligible, and both the hotspot pressure and the jet velocity are nearly constant. The aspect ratios of the cocoon and of the cavity increase with time, while the average cavity pressure decreases. The second epoch starts when the first larger vortices are produced near the jet head. The beam cross section increases and the jet decelerates progressively. The deceleration is repeatedly interrupted by short acceleration phases caused by the temporary replacement of
the Mach disk by a conical shock. The velocity increase during the acceleration phases can exceed 100%.

The evolution of the cocoon and cavity aspect ratios of all models is in agreement with an extended version of the simple theoretical model proposed by Begelmann & Cioffi (BC89) for the inflation of the cocoon. The cavity aspect ratio of the simulated models is too large according to current observations, i.e., the simulated jets are too elongated. This situation may change when the simulations are either extended even further in time or when the assumption of axisymmetry is relaxed. However, both approaches would require prohibitive amounts of computational resources and axisymmetry is, in general, a good approximation to study extragalactic jets.

The size and shape of the cocoon depends on $\eta$, i.e., on the ratio of the densities in the jet and in the ambient medium. According to our definition of the cocoon – the region containing material with a beam particle fraction $0.1 < X < 0.9$ – we find that a light jet (model LH) has a very thin cocoon restricted to a narrow layer girding the beam. Jets with heavier beams (models BC and LC) possess more extended cocoons, more similar to those of observed jets. The differences in the cocoon sizes become particularly evident when comparing the synthetic radio maps that we have computed for all models. The cocoon aspect ratio found for the leptonic hot model LH is unrealistically large, which provides another argument against its realization in nature (see above).

The propagation velocities and the hotspot pressures of our jets agree very well with the analytical and observed values. The strength of the equipartition magnetic field comfortably matches the expectations for a young powerful radio source (whose magnetic fields are thought to be stronger than those of evolved radio sources). The beam velocities of the models are relativistic ($\Gamma \approx 4$) at kiloparsec scales. Provided this is not an artefact caused by the assumed axisymmetry, the results strongly support the ideas of Tavecchio et al. (2000) and Celotti et al. (2001) that the X-ray emission of several extragalactic jets may be due to relativistically boosted CMB photons.

The radio emission of all models is dominated by the contribution of the hotspots. The two cold models (BC and LC) have very similar radio-morphologies with a continuous beam connecting the inlet with the hotspot. This is different from model LH where the synthetic radio maps show a discontinuous and knotty beam even at a viewing angle of 45° (where the effects of Doppler boosting are still small).

We find that the thermal bremsstrahlung emission is strongly suppressed at very high
energies (≈ 10 MeV) in the leptonic cold model because the average temperature of its cavity (∼ 10^{10} K) is much lower than that of the two other models (∼ 10^{11} K). The baryonic and the leptonic hot model both show a gradual transition in the characteristics of their thermal bremsstrahlung emission. At low energies the emission is dominated by the shell, while the interior of the cavity dominates at high energies. All models exhibit a depression in the X-rays surface brightness of the cavity interior in agreement with recent observations (McNamara et al.2000). These observations also show a lack of strong X-ray emission from the boundaries of the radio lobes of Hydra A. This does not contradict our results as Hydra A is a FRI source, while we have been simulating powerful radio sources. In a powerful source the bow shock is a strong shock and, hence, it may produce substantial X-ray emission. FRI sources flare into the external medium and, most probably, they do not have strong bow shocks.

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Figure 1. Values of $\eta$ (solid lines) and $T_b$ (dashed lines) for leptonic jets ($X_{lb} = 1$) as a function of $R_b$ and $W_b$ after fixing $L_{kin}$, $v_J$, $X_{lb}$ and the external medium according to Tab. 1. There are no physical solutions in the areas without $T_b$-contour lines. The thick, solid lines correspond to values of $\eta$ (from top to bottom) of $10^{-4}$, $10^{-3}$, $10^{-2}$. The thick, dashed lines correspond to values of $T_b$ (from left to right) of $10^{10}$K, $10^{9}$K and $10^{8}$K. For baryonic jets the diagram would be similar, except for the values at the thick, dashed lines which would be three orders of magnitude higher: $10^{13}$K, $10^{12}$K and $10^{11}$K, respectively – the reason being that for the same value of $\eta$, one would have a value of $X_{lb}$ three ten folds larger. The squares mark the situation of the three simulated models in this diagram (see 2.3).

Figure 2. Snapshots of the logarithm of the density ($\log_{10}(\rho/\rho_n)$) for the models BC (top panel), LC (central panel) and LH (bottom panel) at $t \approx 6.3 \cdot 10^6$y. The black lines are iso-contours of the beam mass fraction with $X = 0.1$ (outermost) and $X = 0.9$ (innermost). These values correspond to the boundaries of the cocoon and the beam, respectively. Please, note that the colored version of the figure is only provided in the electronic version of the journal.

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Figure 3. Snapshots of the logarithm of the temperature ($\log_{10}(T)$, T in Kelvin) for the models BC (top panel), LC (central panel) and LH (bottom panel) at $t \approx 6.3 \cdot 10^6$y. The black lines are iso-contours of the beam mass fraction with values 0.1 (outermost) and 0.9 (innermost), that correspond to the limits of the cocoon and the beam, respectively. Please, note that the colored version of the figure is only provided in the electronic version of the journal.
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Figure 4. Snapshots of the Lorentz factor $W$ for the models BC (left panel), LC (central panel) and LH (right panel) at $t \approx 6.6 \cdot 10^8$y for model LH and $t \approx 6.3 \cdot 10^8$y for models LC and BC. The black lines are iso-contours of the beam mass fraction with values 0.1 (outermost) and 0.9 (innermost), that correspond to the limits of the cocoon and the beam, respectively. Please, note that the colored version of the figure is only provided in the electronic version of the journal.

Figure 5. Profiles along the jet axis ($r = 0$) at $t \approx 6.3 \cdot 10^8$y of (a) density (in units of $\rho_m$), (b) pressure (in units of $\rho_m c^2$), (c) specific internal energy density (in units of $c^2$), (d) local sound speed (in units of $c$), (e) temperature (in Kelvins), (f) Lorentz factor and (g) beam mass fraction. The units given at the abscissas are in kpc.

Figure 6. Time evolution of (a) the jet length $l_j$, average cavity radius $r_c$, average cocoon radius $r_l$, (b) the aspect ratios (length/radius) of cocoon and cavity, (c) the hotspot pressure $P_{hs}$, average cavity pressure $P_c$ and (d) jet velocity $v_j$ for models BC (dashed), LC (solid) and LH (dotted), respectively. The jet velocities of models LC and LH have been divided by 10 and 100, respectively, in order to enhance the readability of the figure.

Figure 7. Bow shock velocity $v_c$ for the three models.

Figure 8. Hotspot pressure $P_{hs}$ (solid) and ram pressure $\rho_m v_r^2$ of the bow shock (dashed).

Figure 9. Ratio of beam particles to external particles in the cocoon.

Figure 10. Simulated total intensity (in logarithmic arbitrary units) radio maps for models BC, LC and BH (from top to bottom) after $6 \cdot 10^8$y, and for a spectral index of $-0.7$. Both, the jet (to the right) and counter-jet emission are shown for a viewing angle of 45° in all three models. To allow for a better comparison with observed sources the images have been convolved with a circular Gaussian beam with FWHM of $10 R_b$. The same limiting “noise” level is used in all frames to allow for a better comparison of the models.

Figure 11. Surface brightness per unit of frequency for model BC at 10 keV (top) and 10 MeV (bottom), respectively. No plots are shown for the other two models, as the plots for model LH and the plot of model LC at 10 keV would be indistinguishable from those of model BC, and as model LC shows no bremsstrahlung emission at 10 MeV. The colour scales are normalised to maximum values of $4.5 \cdot 10^{-25} \text{erg sec}^{-1} \text{Hz}^{-1} \text{cm}^{-2}$ (top) and $2.8 \cdot 10^{-29} \text{erg sec}^{-1} \text{Hz}^{-1} \text{cm}^{-2}$ (bottom), respectively. Please, note that the colored version of the figure is only provided in the electronic version of the journal.

Figure 12. Beam velocity averaged over cross sectional cuts through the beam perpendicular to the jet axis for different grid resolutions. The solid line corresponds to a resolution of 6 zones per beam radius (working resolution). The dashed and dashed–dotted lines correspond to test resolutions of 4 and 10 zones per beam radius, respectively.
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