On possible ‘cosmic ray cocoons’ of relativistic jets

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ABSTRACT
We consider effects on an (ultra-) relativistic jet and its ambient medium caused by high energy cosmic rays accelerated at the jet side boundary. As illustrated by simple models, during the acceleration process a flat cosmic ray distribution can be created, with gyroradia for highest particles’ energies reaching the scales comparable to the jet radius or the energy density comparable to the ambient medium pressure. In the case of efficient radiative losses a high energy bump in the spectrum can dominate the cosmic ray pressure. In extreme cases the cosmic rays are able to push the ambient medium off, providing a ‘cosmic ray cocoon’ separating the jet from the surrounding medium. The considered cosmic rays provide an additional jet breaking force and lead to a number of consequences for the jet structure and its radiative output. In particular the involved dynamic and acceleration time scales are in the range observed in variable AGNs.

Key words: acceleration of particles – cosmic rays – galaxies: jets – quasars: general–instabilities

1 INTRODUCTION

Shock waves are widely considered as sources of cosmic ray particles in relativistic jets ejected from active galactic nuclei (AGNs). In the present paper we consider an alternative, till now hardly explored mechanism involving particle acceleration at the velocity shear layer, which must be formed at the interface between the jet and the ambient medium (cf. discussion, in a different context, of the turbulence role at a tangential discontinuity by Drobysh & Ostryakov 1998). One should note that there is a growing evidence of interaction between the jet and the ambient medium, and formation of boundary layers, both in observations (e.g. Attridge et al. 1999, Scarpa et al. 1999, Perlman et al. 1999) and in modelling (Aloy et al. 1999). In the next section we discuss this acceleration mechanism in some detail. We point out possible regimes of turbulent second-order Fermi acceleration at low particle energies, next dominated by the ‘viscous’ acceleration at larger energies, and by acceleration at the tangential flow discontinuity at highest energies. In section 3 we shortly consider highest energies allowed in this model by radiative and/or escape losses. Then, in section 4, we discuss time dependent spectra of cosmic rays accelerated at infinite planar flow discontinuity. In the presence of efficient radiative losses a usually formed flat power-law distribution is ended with a bump (in some conditions a nearly monoenergetic spike) followed by a cut-off. Dynamic consequences of cosmic ray pressure increase at the jet boundary are discussed in section 5. In particular a cosmic ray cocoon can be formed around the jet changing its propagation and leading to an intermittent jet activity. Final remarks are presented in section 6.

One should be aware of a partly speculative presentation character of this paper. Till now, besides casual remarks, the considered complicated physical phenomenon was hardly discussed in the literature. One can mention in this respect a discussion of radiation-viscous jet boundary layers by Arav & Begelman (1992) and a discussion of possible cosmic ray acceleration up to ultra-high energies by Ostrowski (1998a). With the present paper we would like to open the considered physical mechanism to more detailed modelling and quantitative considerations.

2 PARTICLE ACCELERATION AT THE JET BOUNDARY

For particles with sufficiently high energies the transition layer between the jet and the ambient medium can be approximated as a surface of discontinuous velocity change, a tangential discontinuity (‘td’). If particles’ gyroradia (or mean free paths normal to the jet boundary) are comparable to the actual thickness of this shear-layer interface it becomes an efficient cosmic ray acceleration site provided the considered velocity difference $U$ is relativistic and the sufficient amount of turbulence is present in the medium (Ostrowski 1990, 1998a). The problem was extensively discussed in early eighties by Berezhko with collaborators (see the review by Berezhko 1990) and in the diffusive limit by Earl et al. (1988) and Jokipii et al. (1989). However, till now...
no one considered the situation with highly relativistic flow characterized with the Lorentz factor \( \Gamma \equiv (1-u^2)^{-1/2} \gg 1 \) and, thus, our present qualitative discussion is mostly based on the results derived for mildly relativistic flows.

Any high energy particle crossing the boundary from, say, region I (within the jet) to region II (off the jet), changes its energy, \( E \), according to the respective Lorentz transformation. It can gain or lose energy. In the case of uniform magnetic field in region II, the successive transformation at the next boundary crossing, \( II \to I \), changes the particle energy back to the original value. However, in the presence of perturbations acting at the particle orbit between the successive boundary crossings there is a positive mean energy change:

\[
\langle \Delta E \rangle = \eta_E (\Gamma - 1) E . \tag{2.1}
\]

The numerical factor \( \eta_E \) depends on particle anisotropy at the discontinuity. It increases with the growing magnetic field perturbations’ amplitude and slowly decreases with the growing flow velocity. The last factor will be particularly important for large \( \Gamma \) flows. For mildly relativistic flows, in the strong scattering limit particle simulations give values of \( \eta_E \) as substantial fractions of unity (Ostrowski 1990). For large \( \Gamma \) we will assume the following scaling

\[
\eta_E = \eta_0 \frac{2}{\Gamma} , \tag{2.2}
\]

where \( \eta_0 \) is defined by the magnetic field perturbations’ amplitude at \( \Gamma = 2 \). In general \( \eta_0 \) depends also on particle energy. During the acceleration process, particle scattering is accompanied with the jet’s momentum transfer into the medium surrounding it. On average, a single particle with the momentum \( p \) transports across the jet’s boundary the following amount of momentum:

\[
\langle \Delta p \rangle = \eta_p (\Gamma - 1) U p , \tag{2.3}
\]

where the \( z \)-axis of the reference frame is chosen along the flow velocity and the value of \( p \) is given as the one before transmission. The numerical factor \( \eta_p \) depends on scattering conditions near the discontinuity and in the highly perturbed conditions (in mildly relativistic shocks) it can reach values being a fraction of unity also. At large \( \Gamma \) we expect \( \eta_p \approx \eta_E \). As a result, there acts a drag force per unit surface of the jet boundary and the opposite force at the medium along the jet, of the magnitude order of the accelerated particles’ energy density. Independent of the exact value of \( \eta_E \), the acceleration process can proceed very fast due to the fact that average particle is not able to diffuse – between the successive energizations – far from the accelerating interface. One should remember that in the case of shear layer or tangential discontinuity acceleration - contrary to the shock waves - there is no particle advection off the ‘accelerating layer’. Of course, particles are carried along the jet with the mean velocity of order \( U/2 \) and, for efficient acceleration, the distance travelled this way must be shorter than the jet breaking length.

The simulations (Ostrowski 1990) show that in favourable conditions the discussed acceleration process can be very rapid, with the time scale given in the observer frame (\( \equiv \) the region II rest frame) as

\[
\tau_{\text{ad}} = \alpha \frac{r_s}{c} , \tag{2.4}
\]

where \( r_s \) is the characteristic value of particle gyroradius in the ambient medium. The introduced acceleration time is coupled to the acceleration length \( l_{\text{ad}} \sim c r_s \) due to particle advection along the jet flow. For efficient scattering the numerical factor \( \alpha \) can be as small as \( \sim 10 \) (Ostrowski 1990). A warning should be risen in this place. The applied diffusion model involves particles with infinite diffusive trajectories between the successive interactions with the discontinuity. Thus reaching stationary conditions in the acceleration process requires infinite times, leading to the infinite acceleration time. However, quite flat spectra, nearly coincident with the stationary spectrum, are generated in short time scales given by Eq. 2.4 and these distributions are considered in the present discussion. One may note that in analytic evaluation of \( \tau_{\text{ad}} \) for the ultra-relativistic jet, applying Eq-s (2.1) and (2.2), large \( \Gamma \) factors cancel each other. For the mean magnetic field \( B_k \) given in the Gauss units and the particle energy \( E_{\text{GeV}} \) given in GeV (1 GeV \( \equiv 10^{18} \text{ eV} \)) the time scale (2.4) reads as

\[
\tau_{\text{ad}} \sim 10^5 \alpha E_{\text{GeV}} B_G^{-1} [s] . \tag{2.5}
\]

Let us remind that in the case of a non-relativistic jet, \( U \ll c \), the acceleration process is of the second-order in \( U/c \) and a rather slow one.

For low energy cosmic ray particles the velocity transition zone at the boundary is expected to appear as a finite-width turbulent shear layer. We do not know of any attempt in the literature to describe the internal structure of such layer on the microscopic level. Therefore, we limit the discussion of the acceleration process within such a layer to quantitative considerations only. From rather weak radiation and the observed effective collimation of jets in the powerful FR II radio sources one can conclude, that the interaction of presumably relativistic jet with the ambient medium must be relatively weak. Thus the turbulent transition zone at the jet boundary must be limited to a relatively thin layer. Within such a layer two acceleration processes take place for low energy particles (by ‘low energy particles’ we mean the ones with the mean radial free path \( \lambda \) much smaller than the transition layer thickness, \( D \)). The first one is connected with the velocity shear and is called ‘cosmic ray viscosity’ (Earl et al. 1988). The second is the ordinary Fermi process in the turbulent medium. The acceleration time scales can not be evaluated with accuracy for these processes, but – for particles residing within the considered layer – we can give an acceleration time scale estimate

\[\star\] The expression (2.4) and the following discussion is valid for an accelerated particle. A small fraction of external particles reflected from the jet can reach a large energy gain, \( \Delta E/E \sim 1^2 \), but these particles do not play a principal role in the cosmic ray energy balance.

\[\dagger\] Below we use also another energy units with respective indices GeV and TeV.

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\[ \tau_i = \frac{r_k}{c} \frac{e^2}{V^2 + (U_V^2)^2} , \]

where \( V \) is the turbulence velocity (\( \sim \) the Alfvén velocity for subsonic turbulence) and \( D \) is the shear layer thickness. The first term in the denominator represents the second-order Fermi process, while the second term is for the viscous acceleration. One expects that the first term can dominate at low particle energies, while the second for larger energies, with \( \tau_i \) approaching the value given in Eq. (2.4) for \( \lambda \sim D \). If the second-order Fermi acceleration dominates, \( \lambda < D(V/U) \), the time scale (2.6) reads as

\[ \tau_i \sim 10^7 E_{\text{TeV}} B^{-1}_G V_3^{-2} \quad [s] , \]

where \( V_3 \) is the turbulence velocity in units of 3000 km/s. Depending on the choice of parameters this scale can be comparable or longer than the expansion and internal evolution scales for relativistic jets. In order to efficiently create high energy particles for the further acceleration by the viscous process and the tangential discontinuity acceleration one have to assume that the turbulent layer includes high velocity turbulence, with \( V_3 \) reaching values substantially larger than 1. Then the scale (2.7) may be much reduced, also because of oblique shocks formed in the turbulent layer and the accompanying first order Fermi acceleration processes. For the following discussion we will assume that such effective pre-acceleration takes place, but the validity of this assumption can be estimated only \textit{a posteriori} from comparison of our conclusions with the observational data. Another possibility is that a population of high energy particles exist in the medium surrounding the jet due to some other unspecifed acceleration processes in the central object vicinity.

The cosmic ray energy spectra generated with the above mechanisms at work are expected to be very flat (Section 4; see also Ostrowski 1998a). With such particle distribution the dynamic influence at the jet and its ambient medium can be due to effects of the highest energy cosmic rays, immediately preceding the spectrum cut-off. Because of the short acceleration time scale (2.4) expected for such particles the acceleration process can provide particles (protons) with energies reaching ultra high energies. Without radiative losses one can obtain particles with \( r_g \sim \) the jet radius, \( R_j \), near the cut-off in the spectrum (Ostrowski 1998a).

For standard jet parameters the considered particle energies may reach \( \sim 10^{19} \) eV. Let us note that the existence of such cosmic rays was suggested by Mannheim (1993) to explain \( \gamma \)-ray fluxes from blazars.

### 3 ENERGY LOSSES

To estimate the upper energy limit for accelerated particles, at first one should compare the time scale for energy losses due to radiation and inelastic collisions to the acceleration time scale. The discussion of possible loss processes is presented by Rachen & Biermann (1993). The derived loss time scale for protons can be written in the form

\[ T_{\text{loss}} \simeq 5 \cdot 10^9 B_G^{-2} (1 + Xa)^{-1} E_{\text{TeV}}^{-1} [s] , \]

where \( B_G \) is the magnetic field in Gauss units, \( a \) is the ratio of the energy density of the ambient photon field relative to that of the magnetic field and \( X \) is a quantity for the relative strength of \( p \gamma \) interactions compared to synchrotron radiation. For cosmic ray protons the acceleration dominates over the losses (Eq-s 2.5, 3.1) up to the maximum energy

\[ E_{\text{max}} \approx 2 \cdot 10^9 \alpha^{-1} [B_G(1 + Xa)]^{-1/2} . \]

This equation can easily yield a large limiting \( E_{\text{max}} \sim 1 \) with moderate jet parameters (e.g. \( B_G \approx 1 \), \( Xa \sim 10^2 \), and \( \alpha \approx 10 \)). However, one should note that the particle gyroradius provides the minimum scale for the acceleration region’s spatial extent (Ostrowski 1998a). Thus, for the actual particle maximum energy \( E_{\text{max}} \) the jet radius should be larger than the respective particle gyroradius \( r_g(E_{\text{max}}) \). E.g., for \( R_j = 10^{16} \) cm and \( B_G = 1 \) the particle energy satisfying the condition \( R_j = r_g \) equals \( E_{\text{max}} \approx 10 \) EeV, what is consistent with the above estimate based on Eq. 3.2.

### 4 ENERGY SPECTRA OF ACCELERATED PARTICLES

The acceleration process acting at the tangential discontinuity of the velocity field leads to the flat energy spectrum and the spatial distribution expected to increase their extension with particle energy. Below, for illustration, we propose two simple acceleration and diffusion models describing these features. For low and high energy particles we consider the time dependent acceleration process at, respectively, the plane shear layer or tangential discontinuity, surrounded with infinite regions for particle diffusion. In the discussion below all particles are ultra-relativistic with \( E = p \).

#### 4.1 A turbulent shear layer

At first we consider ‘low energy’ particles wandering in an extended turbulent shear layer, with the particle mean free path \( \lambda \propto p \). With the assumed conditions the mean time required for increasing particle energy on a small constant fraction is proportional to the energy itself, and the mean rate of particle energy gain is constant, \(< \dot{p} > = \text{const.} \). Let us take a simple expression for the synchrotron energy loss, \(< \dot{p} > \propto p^2 \), to represent any real process acting near the discontinuity. With \(< \dot{p} > \equiv < \dot{p} >_{\text{gain}} - < \dot{p} >_{\text{loss}} \) the transport equation for the particle momentum distribution function \( n \equiv n(t,p,x) \) has the following form

\[ \frac{\partial n}{\partial t} + \frac{\partial}{\partial p} [\dot{p} n] + \frac{\partial}{\partial x} [k_{\perp} \frac{\partial n}{\partial x}] + \left( \frac{\partial n}{\partial t} \right)_{\text{esc}} = Q . \]

where \( x \) measures the distance perpendicular to the shear layer and the escape term at highest energies is represented by \( \left( \frac{\partial n}{\partial t} \right)_{\text{esc}} \).

For the jet boundary acceleration, the jet radius and the escape boundary distance provide energy scales to the process. Another scale for particle momentum, \( p_{\text{esc}} \), is provided as the one for equal losses and gains, \(< \dot{p} >_{\text{gain}} = < \dot{p} >_{\text{loss}} \). As a result, a divergence from the power-law and a cut-off have to occur at high energies in the spectrum. At small
energies, where the diffusive regions are extended and losses non-significant, the considered solution should be close to the power-law.

We used a simple Monte Carlo simulations of the acceleration process to solve Eq. (4.1). In the equation we assume a continuous particle injection, uniform within the considered layer, \( Q = \text{const} \). The diffusion coefficient \( \kappa \) is taken to be proportional to particle momentum, but independent of the spatial position \( x \). With neglected particle escape through the shear layer side boundaries and the considered uniform conditions \( \partial_{t} \rightarrow 0 \) and the spatial diffusion term in Eq. 4.1 vanishes. For the escape term \( \partial_{x} \left( \frac{\partial}{\partial x} \right) \), we simply assume a characteristic escape momentum \( p_{\text{max}} \). At figures 1 and 2 we use \( p_{c} \) as a unit for particle momentum, so it defines also a cut-off for \( p_{c} < p_{\text{max}} \). At Fig. 1, at small momenta the spectrum has a power-law form – in our model \( n(t,p) \propto p^{-2} \) with a cut-off momentum growing with time. However, at long time scales, when particles reach momenta close to \( p_{c} \), losses lower the value of \( < \dot{p} > \) leading to spectrum flattening and piling up particles at \( p \) close to \( p_{c} \). Then, a low energy part of the spectrum does not change any more and only a narrow spike at \( p \approx p_{c} \) grows with time. Let us also note that in the case of efficient particle escape, i.e. when \( p_{\text{max}} < p_{c} \), the resulting spectrum would be similar to one of the short time spectra in Fig. 1, with a cut-off at \( \approx p_{\text{max}} \) (cf. Ostrowski 1998a).

\[ \text{Figure 1. Spectra } N(t,p) \equiv d n/d \log p \text{ of accelerated particles within the boundary shear layer. The results are given in a sequence of times } t_{1} < t_{2} < \ldots < t_{0} \text{ (} t_{i+1} = 10 t_{i} \text{).} \]

4.2 Tangential discontinuity acceleration

An illustration of the acceleration process at the tangential discontinuity have to take into account a spatially discrete nature of the acceleration process. Here, particles are assumed to wander subject to radiative losses outside the discontinuity, with the mean free path proportional to particle momentum \( p \) and the loss rate proportional to \( p^{2} \). At each crossing the discontinuity a particle is assumed to gain a constant fraction \( \Delta \) of momentum (cf. Eqs 2.1, 2.2):

\[ p' = (1 + \Delta) p \quad , \quad (4.2) \]

\[ \frac{1}{p} - \frac{1}{p_{m}} = \text{const} \cdot \Delta t \quad . \quad (4.3) \]

The time dependent energy spectra obtained within this model are presented in Fig. 2, where we choose units in a way to put constant in Eq. (4.3) equal to one and the particle mean free path equals at two considered models at \( p = p_{c} \). Comparison of the results in two models allows to evaluate the modification of the acceleration process by changing the momentum dependence of the particle diffusion coefficient. For slowly varying diffusion coefficient (represented here with a ‘\( \lambda = \text{const} \)’ model) high energy particles which diffuse far away off the discontinuity and loose there much of their energy still have a chance to diffuse back to be accelerated at the discontinuity. In the model with \( \kappa \) quickly growing with particle energy (here the ‘\( \lambda = C \cdot p' \)’ model) such distant particles will decrease their mobility in a degree sufficient to break, or at least to limit their further acceleration. One should note that in both models the spectrum inclination at low energies is the same (here \( n(p) \propto p^{-2} \)).
5 CONSEQUENCES OF THE JET’S ‘COSMIC RAY COCOON’

Cosmic ray distributions in Fig-s 1 and 2 reveal a few spectral components: a flat power-law section at small energies, followed either with a smooth transition to the cut-off, or at first a hard component (‘bump’ or ‘spike’) proceeding the final cut-off. The later case occurs when the radiative losses are enough efficient to pile up particles before the loss dominated high energy range. Spectra without such hard component will appear in cases when particles escape from the acceleration region at low energies, or when the acceleration time scale is longer than the involved dynamical scales (for the jet expansion or slowing down, e.g.). One may note that in the models discussed by us the power-law section of the spectrum has the form $n(p) \propto p^{-2}$.

Normalization of the spectrum at low momenta essential for dynamical considerations is defined by the injection efficiency. This parameter can not be derived from available models or observations and it is treated as a free parameter in the present considerations. The cosmic ray pressure at the jet boundary, $P_{cr} = \int p n(p) dp$, grows with growing injection efficiency and extension of the spectrum in energy. Additionally, the high energy bump can substantially contribute to $P_{cr}$. Let us review a few possibilities arising due to cosmic ray population at the jet boundary, as illustrated at Fig. 3.

Dynamical effects caused by cosmic rays depend on the ratio of $P_{cr}$ to the ambient medium pressure, $P_{ext}$, at the boundary. If, at small particle energies, the acceleration time scale is longer than the jet expansion time or particle escape is efficient at small energies, the formed energetic particle population cannot reach sufficiently high energy density to allow for dynamic effects in the medium near the interface. Then, it acts only as a small viscous agent near the boundary, decreasing slightly gas and magnetic field concentration (cf. Arav & Begelman 1992). In such cases we call the occurring cylindrically distributed cosmic ray population a ‘weak cosmic ray cocoon’. Then, if accelerated particles are electrons or can transfer energy to electrons, a uniform cosmic ray electron population may be formed along the jet leading to the observed synchrotron component with slowly varying spectral index and break frequency. Density of such radiating electrons is expected to have maximum in a cylindrical layer at the jet boundary.

If acceleration dominates losses at small (injection) energies, then the time-dependent high energy part of the spectrum can bear a power-law form with a growing cut-off energy, like the short time distributions at Fig-s 1 and 2. After losses become significant a few possibilities appear. If it happens at low energies, when the acceleration process is limited to the turbulent shear layer, a power-law with a growing sharp spike preceding a cut-off energy will appear. If in such case the increasing cosmic ray pressure in a cocoon could reach values comparable to the medium pressure, a substantial modification of the jet boundary layer is expected. Below we will discuss various possibilities arising in such cases of the ‘dynamic cosmic ray cocoons’. If the acceleration at the tangential discontinuity resembles our models in section 4.2, then, depending on the conditions near the jet, the cosmic ray pressure may stabilise at an intermediate $P_{cr} < P_{ext}$, or grow to form the dynamic cosmic ray cocoon.

Let us consider a possible scenario of dynamic interaction of high energy cosmic rays with the jet and the ambient medium. The particles are ‘injected’ and further accelerated at the jet boundary. Growing number of such particles results in forming the cosmic ray pressure gradient outside the jet pushing the ambient medium apart. Additionally, an analogous gradient may be formed directed into the jet, helping to keep it collimated. The resulting rarefied medium or partly emptied of the magnetized plasma space near the jet boundary will decrease the acceleration efficiency. Thus the cosmic ray energy density may build up only to the value comparable to the ambient medium pressure, when it is able to push the magnetized plasma away. Because the diffusive escape of charged particles from the cosmic ray cocoon is not expected to be efficient (contrary to photons considered by Arav & Begelman 1992), in some cases the blown out volumes could be quite large, reaching the values comparable to $R_j$ or even to the local vertical scale of gas. The accumulated cosmic rays can be removed by advection – in the form of cosmic rays’ filled bubbles or cosmic ray dominated winds – outside the active nucleus into regions of more tenuous plasma, or simply outside the jet at larger distances from the central source.

The jet moving in a space filled with photons and high energy cosmic rays (cf. Fig. 3) is subject to the braking force due to scattering this species (e.g. Sikora et al. 1998; for the photon breaking). For cosmic rays with $\lambda \sim R_j$ both types of particles penetrate relatively freely inside the jet and the breaking force is exerted more or less uniformly within its volume, in rough proportion to the electron (or pairs’) density for the photon breaking and to the turbulent magnetic field energy density for the cosmic ray breaking. If cosmic ray cut-off energy is lower, with the equivalent $\lambda < R_j$, the cosmic ray breaking force acts within the jet.

Figure 3. A schematic view of the situation considered near the central source: ‘BH’ is a central black hole, ‘CR’ denotes a cosmic ray particle accelerated at the jet boundary, $\gamma$ is a photon radiated from the central source vicinity.
boundary layer of width $\lambda$. From Eq. (2.3) we estimate the cosmic ray breaking force per unit jet length to be

$$f_{b,cr} = 2\pi \eta_0 (\Gamma - 1) P_{ex} R_j,$$

(5.1)

where we consider $\lambda \leq R_j$ and we put $U = 1$. From the above discussion, in the stationary conditions one can put $P_{ex} \approx P_{ext}$, where $P_{ext}$ is the external medium pressure. For a jet with a (relativistic) mass density $\rho_j$, with Eqs (2.2,3) and $\eta_0 = \eta_b$, the jet breaking length due to cosmic rays is

$$L_{b,cr} = R_j \frac{\Gamma}{4\eta_0} \frac{\rho_j c^2}{P_{ext}}.$$

(5.2)

For example, assuming $P_{ext} = \rho_j c^2$ and $\eta_0 = 0.25$, we obtain $L_{b,cr} = \Gamma R_j$.

Because of dynamic (‘d’) form of pushing out the ambient medium and following it cosmic rays’ escape, the backreaction of this process at particle acceleration is expected to make the full process unstable, with an intermittent behaviour seen in longer time scales. The full configuration with the ‘heavy’ ambient gas supported with the ‘light’ gas of ultra-relativistic particles in the cosmic ray cocoon is expected to be subject to the Raileigh-Taylor (‘RT’) instability. A related characteristic time scale can be roughly estimated as

$$t_{RT} \sim \left( \frac{L}{2\pi g} \right)^{1/2}.$$

(5.3)

where $g$ is the gravitational acceleration and $L$ the scale of instability (e.g. for $g = 10^2$ cm/s$^2$, $L = 10^{17}$ cm and additional requirement of the sound velocity comparable to $c$ the time scale $t_{RT}$ is estimated to be below 1 yr).

Another type of instability can be generated by the time dependence of the non-linear acceleration process. Continues injection of (low energy) seed particles to the acceleration process can be continued till the cosmic ray pressure becomes equal to $P_{ext}$. Then, the ambient plasma is pushed away from the jet and its interaction with the jet boundary surface diminishes. It must lead to decrease of the injection efficiency, if acceleration at the turbulent surface layer is responsible for the process. Then, cosmic ray energy density contained in still accelerated highest energy particles increases until they manage to escape from the jet vicinity, allowing for re-establishment of original conditions. It allows the ambient medium to ‘fall down’ at the jet to start a new phase of intensive interaction between the jet and the ambient magnetized plasma, initiating efficient injection of low energy seed particles. The process should be accompanied with intense kinetic energy dissipation processes and a radiation flare at all frequencies. In the flare phase one can expect substantial weakening of the cosmic ray jet breaking mechanism allowing for larger jet velocity and forming internal shocks. Next, the full process could repeat with a time scale comparable to the time required for removing highest energy particles from the system. For the inequality $t_{RT} > \tau_{id}(E_{\text{max}})$ a continues (diffusive, or as a wind) particle escape will govern the process. Then one may expect smaller variations of the output radiation flux.

The presented discussion assumed the cylindrical symmetry of the unstable flow, which may be not true. However, any large amplitude perturbation of the conditions near/in the jet can not be much smaller than the spatial scale $R_j$ and the respective observer’s time scale shorter than $R_j/c$.

6 CONCLUSIONS AND FURTHER SPECULATIONS

Limited to the hydrodynamic approach the present discussion is intended to provide an alternative view of the AGN central activity related to the jet outflows. The acceleration of cosmic rays up to extremely high energies occurs in a natural way at the relativistic jet boundary if there are effective preliminary acceleration mechanisms providing seed particles with mean free paths comparable to the width of a boundary layer. With the assumption that such processes work efficiently we discussed several possible consequences for the conditions in the spatial volume containing the jet. Here, the main factor playing the role is a flat spectrum cosmic ray population carrying substantial energy density in highest energy particles. These particles may dynamically influence the jet flow and conditions in the surrounding medium, without direct radiative effects. A possibility is considered of the jet timely separated from the ambient gas by the layer filled with cosmic rays and ambient photons. During such a phase the electromagnetic radiation produced in the jet can more easily escape from the AGN centre, to reach observer situated close to the jet axis direction. Also, the plasma self-absorption frequency can be decreased for such observer if the optical depth of the upper plasma layers does not dominate the output. As we expect larger intensity of generated low-frequency radiation when the ambient medium directly pushes on the jet, the discussed picture should be characterized with a positive correlation of the radiation intensity with the self-absorption frequency shift to larger values (see Bottcher 1999 for a recent discussion of such shifts within the standard jet picture). Also, for flares in some BL Lac objects with a week-month time scale, the beginning of the flare should be seen approximately at the same time at all non-absorbed frequencies. Only later evolution of the introduced disturbance will lead to shifts of flare maxima at different frequencies. One should remember that we do not include into this discussion other radiation sources (accretion disc, corona) in the active nucleus vicinity.

The instabilities related to the discussed process may lead to temporary variations of jet flow velocity and the degree of the jet surface perturbation. Perturbations in jet flow introduced by these instabilities may also lead to shock wave formation with its observational consequences. As a result the processes accelerating lower energy cosmic rays and cosmic ray electrons are expected to have fluctuating nature with time scales estimated in Eq. (5.3) for large changes. The processes occurring inside the jet can be characterized with the observer’s time scale a factor of $\Gamma$ shorter. However, in the situation with the jet perturbation introduced by the external process, the actual time scale will be intermediate between the internal, Lorentz contracted one and the external perturbation scale.

In the above discussion we avoided considering the acceleration of electrons (or pairs). In the mentioned model of Mannheim (1993) energetic electrons arise as a result of cas-
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cading of pairs resulting from the energetic proton interactions inside the jet. One can consider also different scenarios providing the cosmic ray electrons. E.g., in the space close to the jet boundary a large power can be stored in the highly anisotropic population of cosmic ray protons. Such distribution is known to be unstable and it leads to creation of long electromagnetic plasma waves. Damping of such waves by pairs may be very efficient acceleration process providing cosmic ray electrons (cf. Hoshino et al. 1992, in a different context). Of course, the short plasma waves at the jet boundary can be also generated by velocity shear.

In the presented evaluations we often consider the situation with particles starting to play a dynamic role in the system when their energies reach scales yielding gyroradia \( r_g \sim R_j \). If particles become dynamically important at lower energies, with \( r_g << R_j \), all considered time scales should be respectively scaled down. Then, the jet breaking force due to cosmic rays is acting only at the external layers of the jet, generating magnetic stresses along it.

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REFERENCES

Aloy M.A., Ibáñez J. M.*, Martí J.M.*, Gomez J.L., Müller E., 1999, ApJ Lett., accepted [astro-ph/9906425]
Attridge J.M., Roberts D.H., Wardle J.F.C., 1999, ApJ, 518, L87
Aran N., Begelman M.C., 1992, ApJ, 401, 125
Bednarz J., Ostrowski M., 1996, MNRAS, 283, 447
Berezhko E.G., 1990, Preprint Fractional Acceleration of Cosmic Rays, The Yakut Scientific Centre, Yakutsk.
Böttcher M., 1999, ApJ Lett, accepted (ASTRO- PH/9901165)
Drobysh A. Yu., Ostyakov V. M., 1998, Astronomy Rep., 42, 679
Earl J.A., Jokipii J.R., Morfill G., 1988, ApJ, 331, L91
Hoshino M., Arons J., Gallant Y.A., Langdon A.B., 1992, ApJ, 390, 454
Jokipii J.R., Kota J., Morfill G., 1989, ApJ, 345, L67
Mannheim K., 1993, A&A, 269, 67
Ostrowski M., 1990, A&A, 238, 435
Ostrowski M., 1997, in Proc. Int. Conf. ‘Relativistic Jets in AGNs’, eds. M. Ostrowski, M. Sikora, G. Madejski, M. Begelman; Cracow (p. 153)
Ostrowski M., 1998a, A&A, 335, 134
Ostrowski M., 1998b, in Proc. Vulcano Workshop ”Frontier Objects in Astrophysics and Particle Physics”, eds. F. Giovanelli & G. Mannocchi, in press [astro-ph/9808234]
Perlman E.S., Biretta J.A., Fang Z., Sparks W.B., Macchetto F.D., 1999, AJ, accepted [astro-ph/9901170]
Rachen J.P., Biermann P., 1993, A&A, 272, 161
Scarpa R., Urry C.M., Falomo R., Treves A., 1999, ApJ, accepted [astro-ph/9906463]
Sikora M., Sol H., Begelman M.C., Madejski G., 1998, A&A Suppl, 120, 579