BURSTING PARTICLE ACCELERATION IN RADIO JETS

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I. INTRODUCTION

This work follows on the papers by Henriksen, Bridle and Chan (1982) (HB) and by Eilek and Henriksen (1983). These papers introduced a comprehensive model of (hydrodynamic) turbulence driven Alfven wave, resonant, particle acceleration. Very similar ideas were introduced independently by Fedorenko (1980), but the details were sketchy and the universality of the resulting spectral index was not realized. Eichler (1979) has come closest to these ideas previously (again independently) in a very original paper directed at ion acceleration in solar flares. He includes a back reaction on the turbulent eddies however which requires an adhoc treatment and is probably unnecessary (we replace this by the Lighthill noise analysis: HBC) and he omits synchrotron and adiabatic losses. Moreover, the quasi-linear theory is not used in calculating the distribution function. Nevertheless, it is reassuring that the same intuitive physics, even to the self-similar universality, should arise independently.

This mechanism establishes in a volume the Alfven wave spectral density necessary for the particle acceleration, as a balance between MHD wave emission from the turbulent vortices throughout the volume (Lighthill noise) and the damping losses to relativistic and thermal particles. The wave number dependence of the wave spectral density $W_A(k)$, is fixed both by the spectrum of the fluid turbulence (e.g. 'young' that is, steep; Kolmogorov or Kraichnan) and by the wave number dependence of the damping (e.g. Eilek 1979). The stochastic resonant interaction of this spectrum with relativistic electrons (second order Fermi process) can, assuming isotropy and quasi linear theory, be described as a diffusion in energy ($\varepsilon = cp$ for relativistic particles) space with coefficient $D = Gp^{\tilde{\nu}}$ (e.g. Lacombe, 1979). The index $\tilde{\nu}$ is that of the Alfven wave spectral density $W_A(k) = W_o k^{-\tilde{\nu}}$. The quanty G depends on the normalizing factors of the particle and wave spectrum, but not on energy or wave number. Eilek and Henriksen (EH)

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have argued that there are feed-back mechanisms operating alternately on the timescales of the synchrotron lifetime and the acceleration time which yield $\tilde{\nu} = 3$ and $S/G = 6-s_t$ asymptotically. Here S is the synchrotron (Compton losses may be included by adding the effective B of the radiation field to B) loss coefficient (B is magnetic field strength and other symbols are standard)

$$S = \left(\frac{4e^4}{9m_e^4 c^6} \right) B^2, \tag{1}$$

and s_t gives the wave number dependence of the Alfven wave driving power/unit wave number (per unit volume) due to the turbulence as (κ is the eddy wave number)

$$I_a(k) = I_o(k/\kappa_T)^{-s_t} . \tag{2}$$

The quantity I_o depends only on the turbulent flow properties, and κ_T is the Taylor scale of the turbulence. The index $s_t = 1.5$ for Kolmogorov turbulence, 1 for Kraichnan (i.e. 'old', equipartition turbulence) and 2 for young turbulence. They have also shown that under these conditions, the solution for the electron distribution function is

$$f = f_o\ p^{-s}, \tag{3}$$

where s tends asymptotically to $S/G = 6-s_t$ over substantial portions of energy space. This converts to a synchrotron spectrum $I_\nu \propto \nu^{(s_t-3)/2}$ (e.g. $\propto \nu^{-0.75}$ for Kolmogorov).

Thus far however, the theory has been discussed purely locally, without reference to convection or to the expected positional dependence of the synchrotron brightness. In the extragalactic jet sources, such variation with position is the defining property (both parallel and perpendicular to the jet axis). In this paper, therefore, we study how the EH theory applies in the presence of convection and of axial variations in the local properties of the turbulence. The well studied sources NGC315 and NGC6251, are referred to for illustration.

II. THEORETICAL FRAMEWORK

We may write the equation for the electron distribution function, $f(\underset{\sim}{r}',t')$, in the local proper frame of the jet in the coordinates of the background (galaxy) frame $(f(\underset{\sim}{r},t) = f(\underset{\sim}{r}',t'))$ as

$$\frac{\partial f}{\partial t} + \underset{\sim}{u}\cdot\underset{\sim}{\nabla} f - \left(\frac{\underset{\sim}{\nabla}\cdot\underset{\sim}{u}}{3} \right) p \frac{\partial f}{\partial p} = \frac{S}{p^2} \frac{\partial}{\partial p} (p^4 f)$$

$$+ \frac{G}{p^2} \frac{\partial}{\partial p} \left[p^{2+\tilde{\nu}} \frac{\partial f}{\partial p} \right] \quad \ldots . \tag{4}$$

For definiteness, we proceed by substituting the CH velocity field

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(Chan and Henriksen, 1980) for the jet velocity $\underset{\sim}{u}$ so that

$$u = u_z \left(\frac{r}{R_j} \frac{dR_j}{dz} , \frac{u_\phi}{u_z} \frac{r}{R_j} , 1 \right) \tag{5}$$

where (r,ϕ,z) are cylindrical coordinates relative to the jet axis, $R_j(z)$ is a fiducial boundary stream line usually identified with a radius determined from the transverse brightness profile and $u_z = u_z(z)$, $u_\phi = u_\phi(z)$. With this ansatz the operative equation becomes

$$\frac{\partial f}{\partial t} + u_z \frac{\partial f}{\partial z} - \frac{1}{3R_j^2} \frac{d}{dz} (R_j^2 u_z) \, p \frac{\partial f}{\partial p}$$

$$= \frac{S}{p^2} \frac{\partial}{\partial p} (p^4 f) + \frac{G}{p^2} [p^{2+\tilde{\nu}} \frac{\partial f}{\partial p}] \tag{6}$$

This equation becomes complete when $u_z(z)$, $S(z)$, $G(z)$ and $\tilde{\nu}$ are given.

In the turbulence driven model of (EH)

$$G = \frac{(s - 2)}{2\pi\tilde{\nu}(\tilde{\nu}+2)} \left(\frac{\Omega m_e}{c} \right)^{(s-2-\tilde{\nu})} \frac{I_o}{f_o} \kappa_T^{s_t} ,$$

$$I_o = \eta_A \left| \frac{6-4m}{3-m} \right| \rho v_a^3 (E_t/\rho v_a^2 R)^{\frac{3}{3-m}} , \tag{7}$$

$$\tilde{\nu} = s + s_t - 3,$$

where Ω is the classical gyrofrequency $(eB/m_e c)$, E_t is the turbulent kinetic energy density, $s_t = 3(m-1)/(3-m)$, and m is the spectral index of the hydrodynamic turbulence (e.g. m = 5/3 for Kolmogorov), v_a is the appropriate Alfven speed, η_A is a number of order unity and $R >> 1$ is a factor which allows for the fraction of turbulent energy in the resonant region (below the Taylor scale) $R \gtrsim (\kappa_T/\kappa_o)^{(m-1)}$. EH show $\tilde{\nu} = 3$ for a form stable spectrum when synchrotron losses and resonant acceleration dominate equation (6), which determines s as $6-s_t$. However this need not be so in other regimes. We also note that if the explicit model for the jet turbulent energy density given by HBC is adopted, then

$$I_o = \eta_A \left| \frac{6-4m}{3-m} \right| R^{\frac{3}{3-m}} \rho v_a^{-s_t} (u_z \, dR_j/dz)^{6/(3-m)} . \tag{8}$$

The dependence on f_o in G reflects the coupled integro-differential nature of the full problem, which we have dealt with by replacing f with $f_o p^{-s}$ in the integrals for G. We must always check for self-consistency a posteriori however. Moreover, this dependence can lead to self-limiting strong acceleration wherein G decreases as f_o increases, and hence to a relaxation oscillation.

We proceed now to investigate various separate physical regimes of equation (6).

a) Adiabatic Expansion

In this case Sp $z/u_z \ll 1$ and $Gp^{(\tilde{\nu}-2)} z/u_z \ll 1$ so that (6) becomes

$$\frac{\partial f}{\partial t} + u_z \frac{\partial f}{\partial z} - \frac{1}{3R_j^2} \frac{d}{dz} (R_j^2 u_z)\, p \frac{\partial f}{\partial p} = 0.$$

This has the general solution

$$f = F(t - \int \frac{dz}{u_z} , R_j^2 u_z p^3) \qquad (9)$$

where F is formally arbitrary but is usually determined by the initial conditions. If, for example, at one point in the jet a burst of relativistic electrons is produced with a spectrum of the form $f_o(t_o)p^{-s}$, then subsequently

$$f = f_o(t - \int \frac{dz}{u_z})(R_j^2 u_z p^3)^{-\frac{s}{3}} \qquad (10)$$

where $t_o \equiv t - \int dz/u_z$ is the (convective) 'retarded time'. We normally observe at some fixed t.

b) Bursting Acceleration

In this case $Gp^{(\tilde{\nu}-2)} \gg u_z/z$, Sp, so that equation (6) becomes <u>in the appropriate range of p</u>

$$\frac{\partial f}{\partial t} = \frac{G}{p^2} \frac{\partial}{\partial p} [p^{2+\tilde{\nu}} \frac{\partial f}{\partial p}],$$

and the electron spectrum is dominated by the local acceleration.

The general solution for the Laplace transform of f, $\overline{f}(\omega,p) \equiv \int_0^\infty e^{-\omega t} f(t,p)\, dt$ is (when $\tilde{\nu} = 3$), $\overline{f}(\omega,p) = \frac{1}{G} \int_{p_{min}}^{p_{max}} p'^2 dp'\, G_\omega(p,p') f_o(p')$,

$$G_\omega(p,p') = \frac{2I_4(2\sqrt{\omega/Gp_>})K_4\ (2\sqrt{\omega/Gp_<})}{(pp')^2},$$

$p_>, p_<$ are respectively the greater or lesser of p,p', f_o is the initial distribution and I_4, K_4 are Bessel functions. Asymptotes may be found by using $f(\infty,p) = \lim\limits_{\omega\to o} \omega f$, but it is normally more convenient to find these as limits to the self-similar behaviour. When however, $f_o(p) = (N_o/4\pi p_o^{\ 2})\delta(p-p_o)$, then for $p > p_o$

$$f(\infty,p) = \lim_{\omega\to o} \left\{ \frac{\omega}{4G}\ \frac{N_o}{4\pi} \right\} p^{-4}.$$

This supports our interpretation of the asymptote (13b) below. For $p < p_o$, the result is

$$f(\infty,p) = \lim_{\omega\to o} \left\{ \frac{\omega}{4Gp_o}\ \frac{N_o}{4\pi p_o^{\ 3}} \right\}.$$

We look for the asymptotic behaviour by setting

$$x \equiv Gp^{(\tilde{\nu}-2)}t$$
$$g \equiv 4\pi p^3 f \qquad (11)$$

so that the self-similar equation to be solved is

$$\frac{dg}{dx} = (\tilde{\nu}-2)\frac{d}{dx}[(\tilde{\nu}-2)x^2\frac{dg}{dx} - 3gx].$$

This has the general solution (see also EH and Lacombe, 1979)

$$g = x^{3/(\tilde{\nu}-2)}\exp(-x^{-1}/(\tilde{\nu}-2)^2)\ [A_1 + A_2\int\frac{dx\ \exp(x^{-1}/(\tilde{\nu}-2)^2)}{x^{(2\tilde{\nu}-1)/(\tilde{\nu}-2)}}] \qquad (12)$$

and the asymptotic region is found with $x >> 1$. In this latter limit the dominant terms of the two modes are

$$f_1 \sim \frac{A_1}{4\pi}(Gt)^{\frac{3}{\tilde{\nu}-2}} \quad .... (a)$$

and

$$f_2 \sim \frac{A_2}{4\pi}(Gt)^{-1}\ p^{-(\tilde{\nu}+1)} \quad ... (b). \qquad (13)$$

The physical mode is (13b) as it gives an upward flux of particles in momentum space and corresponds to the non-trivial steady state $f \propto p^{-(\tilde{\nu}+1)}$. The solution is consistent with $s = \tilde{\nu}+3 - s_t$ <u>for all $\tilde{\nu}$</u>

if $s_t = 2$ (young turbulence), so that the initial spectrum can be expected to be maintained ($\tilde{\nu} \simeq 3$).

The mode (13a) gives a downward flux of particles in momentum space and corresponds to the trivial steady state f = const. It requires an initial supply of high energy particles to be realized ($p < p_o$ in the Laplace Transform example above).

We note finally that the inverse dependence on G in (13b), is due to f(p) varying inversely with the rate at which particles are accelerated beyond p.

c) Synchrotron Loss Dominated

Here we require that $Sp \gg u_z/z$, $Gp^{(\tilde{\nu}-2)}$, so that (6) is simply

$$\frac{\partial f}{\partial t} = \frac{S}{p^2} \frac{\partial}{\partial p} (p^4 f).$$

This equation has the general solution

$$f = (F(t - 1/Sp)p^{-4},$$

where F is arbitrary. However the initial value problem is not well behaved past t = 1/Sp for strict initial power laws ($f \propto (1/Sp-t)^{s-4} p^{-4}$ if $f_o \propto p^{-s}$), so that it is generally more useful (EH) to use the asympote

$$f = \frac{1}{S} \frac{A_1}{4\pi p^4} \frac{1}{(t - 1/Sp)} . \tag{14}$$

for $t \gg 1/Sp$. The particles stream downward in momentum space.

d) Form Stability

This case has been discussed at length in EH who argue for $\tilde{\nu} = 3$ and S/G = s. We require further that $Sp\ z/u_z$ and $Gp\ z/u_z \gg 1$, so that (6) becomes

$$\frac{\partial f}{\partial t} = \frac{S}{p^2} \frac{\partial}{\partial p} (p^4 f) + \frac{G}{p^2} \frac{\partial}{\partial p} [p^5 \frac{\partial f}{\partial p}] .$$

A true steady state may be achieved locally in two modes;

$$f_1 = A_1\ p^{-S/G}$$

and (15)

$$f_2 = F_o/(4\pi G(S/G - 4))p^{-4},$$

if $S/G \neq 4$. The self-similar asymptotes have two corresponding modes

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which are at large $x \equiv Gpt$

$$f_1 \sim A_1 (Gt)^{3-S/G} p^{-S/G}$$
$$f_2 \sim A_2 (Gt)^{-1} p^{-4} \qquad (16)$$

with $S/G = 6 - s_t = s = 3 + 2\alpha$.

We expect this form stable condition to apply in the mean in a region where adiabatic losses are not important. Turbulent bursting acceleration followed by synchrotron decay may well be superimposed on this mean behaviour in a form of relaxation oscillation. The remarkable conclusion to be drawn from equations (10), (13b), (14) and (16) is that the spectral index is not expected to vary greatly in any of these processes, and may be expected to be close to $\alpha = 0.5$ in the bursting mode (young turbulence and synchrotron loss).

III. PREDICTED SYNCHROTRON BRIGHTNESS VARIATIONS

We may calculate the surface brightness along a pathlength $d(\theta)$ (distance d in the source at an angle θ to the magnetic field) produced by optically thin synchrotron radiation of relativistic electrons as

$$I_\nu = c_5(\alpha)\ 4\pi c^{2\alpha} \left(\frac{\nu}{2c_1} \right)^{-\alpha} (B \sin\theta)^{\alpha+1} f_o\ d(z), \qquad (17)$$

with

$$d_{\parallel}(z) = 2R(z) \operatorname{cosec} \theta$$
$$d_{\perp}(z) = 2R(z) \sec \theta,$$

where θ is the angle between the line of sight and the magnetic field (pitch angle), and $\parallel$ and $\perp$ refer to the direction of the field relative to the axis of the jet. There are various cases. We express our results in terms of the observed index $\alpha \equiv (s-3)/2 = (3-s_t)/2$.

a) Adiabatic Expansion

If at any point (the 'source point') in the jet, a steady distribution function becomes adiabatic then equation (10) gives $f \propto (R_j^2 u_z)^{-(1+\frac{2\alpha}{3})}$ subsequently. If moreover $B \propto z^{-1}$ and u_z, dR_j/dz are constant, then (17) gives $I_\nu \propto z^{-(2+7\alpha/3)}$, which agrees with a well known result. Note however, that if $u_z \propto z^{-1}$ (Landau-Squires jet) then $I_\nu \propto z^{-(1+5\alpha/3)}$. But such a velocity decline is probably excessively wasteful of energy in the large scale jets.

In the general case of equation (10) we find

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$$I_\nu = f_o \left(t - \int \frac{dz}{u_z}\right)(R_j^2\, u_z)^{\frac{-(1+2\alpha)}{3}}\, B^{(1+\alpha)}\, d(z). \qquad (18)$$

The time dependent factor will not change rapidly until $z = O(u_z t)$ where t is the age of the outburst. This may be significantly close to the source point (in terms of real brightenings) but it is unlikely to negate the steady adiabatic decline over many scale in z.

b) Bursting Mode

Suppose that G increases suddenly at a point in the jet due to a sudden increase in the turbulent input power I_o through the burst factor $(\kappa_T/\kappa_o)^{1/2}$ (see e.g. EH) or otherwise (e.g. rapid variation in ρ - a cloud encounter - or rapid increase in R_j). One expects a phase in which particles are pumped up in energy space according to (13b). We see readily from equation (19) that $I_\nu \propto (z/G)\, B^{1+\alpha}$ in such a phase, where G might be calculated initially from equation (7) and (8) keeping f_o constant. However with G large the time scale of this phase is very short, and G will vary rapidly as the number of relativistic particles (f_o) increases. Hence we proceed directly to the 'trickle down' phase of the outburst when synchrotron losses are bringing the particles back down in energy space. Either equation (14) (if S >> G for a time) or equation (16) should apply in this phase. From the first of the form stable modes (16), and (17) we find $f \propto G^{-2\alpha}$ and,

$$I_\nu \propto (d(z)/G^{2\alpha})\, B^{1+\alpha}$$

or, since $G \propto S \propto B^2$ in this mode,

$$I_\nu \propto d(z) B^{(1-3\alpha)}. \qquad (19)$$

Consequently, if $d(z) \propto z$, $I_\nu \propto z$, $z^{1.5}$ or z^2 according as B is constant, $\propto z^{-1}$ or $\propto z^{-2}$. We note that the actual magnetic field is unlikely to be the equipartition magnetic field in these bursts because of the sudden production of relativistic particles.

This brightening phase must end at p when

$$Sp\, z/u_z \geq 1 \qquad (20)$$

at which point an adiabatic decline will begin, in the absence of additional bursting. Identifying such behaviour at sufficiently high frequency can lead to a useful lower limit for u_z.

Should the strictly form stable mode (see 15) be achieved in a region, we see that f_o = constant there. This yields from (17) that

$$I_\nu \propto d(z)\, B^{1+\alpha}. \qquad (21)$$

Therefore if $d(z) \propto z$, $I_\nu \propto z, z^{-\alpha}, z^{-(1+2\alpha)}$ according as B = constant, $\propto z^{-1}$ or $\propto z^{-2}$. This gives brightness declines from the burst peak which are less steep than adiabatic and quite symmetrical in form relative to the rising edge. In a region of relaxation ascillation where first G dominates and then S dominates, this probably corresponds to the mean behaviour.

IV. NGC6251 AND NGC315

The large source NGC315 (Willis et al. 1981; HBC and unpublished data of Bridle, Fomalont, Palimaka and Henriksen) has a main jet which is remarkably smooth in its brightness variation. There is a rapid rise to a peak at $z \sim 10'$, followed by $I_\nu \sim z^{-0.6}$ out to $z \sim 20''$. This is followed by $I_\nu \propto z^{-1.6}$ out to $z \sim 100''$, and after $z \sim 120''$ there is a nearly adiabatic decline. This latter decline is associated with $dR_j/dz \to 0$. HBC have interpreted this brightness behaviour simply as proportional to the turbulent driving associated with a dynamical model, which corresponds to the strict steady state of the form-stable mode (EH). It is tempting to interpret NGC315 with our more detailed theory here, as a single slow burst. In the steady form-stable mode interpretation we must have $B \simeq$ constant at the base, then $B \propto z^{-1}$ from 10" to 20" and $B \propto z^{-1.6}$ from z = 20" to $\sim$ 100" (we have used $\alpha = 0.6$). This latter variation does not correspond well to the equipartition field variation (Willis et al. 1981) which tends to be flatter (as one might expect if relativistic particles are being produced).

If we interpret the adiabatic decline at $z \sim 120''$ ($1'' \equiv .314$ kpc) according to equation (20) and use $B = B_{eq} \simeq 3\ \mu G$, then $u_z \gtrsim 4.4 \times 10^2\ \nu_{max}^{1/2}\ cms^{-1}$, where ν_{max} is strictly the maximum frequency at which the decline is seen. This gives $u_z \gtrsim 4.4 \times 10^7\ cms^{-1}$ if $\nu_{max} \sim 10^{10}$ Hz.

The spectacular one-sided source NGC6251 has been thoroughly studied by Perley, Bridle and Willis (1983). It has a complex bursting structure which we suggest is due to many rapid events of the same general type as the single slow burst in NGC315. We will leave a detailed discussion to a later work. Here we observe that the various $I_\nu(z)$ observed in the different segements of the jet can generally be understood in terms of our ideas above. Thus, the inner jet has $I_\nu \propto z$ for $9'' < z < 13''$, so that $B \simeq$ constant based on the form stable mode. The mean behaviour out to z 100" is $I_\nu \propto z^{-1.5}$, which, as in NGC315, requires a form stable steady state and $B \propto z^{-1.6}$. The central, faint region of the jet has $I_\nu \propto z^{-2.2}$ ($120'' < z < 180''$) which suggests a form stable steady state with $B \propto z^{-2}$.

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When $dR_j/dz \neq$ constant, it is better to use the steady form stable brightness variation in the form $I_\nu \propto B(R_j)^{(1+\alpha)} R_j$. Thus a $B \propto R_j^{-1.6}$ gives symmetrical $I_\nu \propto R_j^{-1.5}$ (I_ν rising with decreasing R and vice versa) as is observed. The same is true with $B \propto R_j^{-2}$ (or a little steeper) to give symmetrical $I_\nu \propto R_j^{-2.2}$. The form stable brightening mode (eq. (19)) may be distinguished from this steady behaviour by brightness peaks that coincide with maximum radii, rather than the converse as above, and as is mostly observed. Thus these variations can be mainly dynamically driven (through R_j), with steady turbulence.

There are two anomalous regions at z = 50' and at z = 210" which have the adiabatic slope (α = 0.6) both rising and falling. That is, approximately, $I_\nu \propto R_j^{-(2+7\alpha/3)} = R_j^{-3.4}$. Applying equation (20) is not too useful for these burst however, because $(\Delta z)_{burst} \ll z$ and only a very low, lower limit is obtained at radio frequencies.

The oscillations following the burst at z = 20" out to z ≃ 110" appear to be almost dynamical 'ringing' of R_j while the turbulence is relatively undisturbed from the form stable steady state except at z ≃ 50" where the compression and expansion is so rapid as to be adiabatic. The brightness oscillations in the outer jet region appear to be slipping in phase relative to R_j. This suggests that the turbulence is becoming intermittent, with new particles produced again in bursts.

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DISCUSSION

Norman: How does this generalize to the important particle acceleration processes generated in shocks?

Henriksen: This is a volume mechanism, whereas the shock wave acceleration is more discrete. Moreover, our mechanism is second order Fermi acceleration rather than first order. Shock waves in laboratory jets (not necessarily the same regime) tend to fade into volume turbulence. However, the shock wave mechanism can also produce a universal spectrum and it may well operate in exceptionally dissipative regions.

Vlahos: Coronal streamers are jets that are closer to us. Have you scaled down your process of particle acceleration to such objects, and what are the results?

Henriksen: No. I have not attempted to apply this to the sun. It should be done. The mechanism becomes inefficient, however, if the thermal particle density is 'too high'. However, the scaling is not obvious, as it depends on magnetic field strength, spatial scales, turbulent intensity, and so on.

Sturrock: What damps the Alfvén waves? Do the effects discussed by Hasegawa play an important role?

Henriksen: The Alfvén waves are damped by the resonant interaction with the relativistic particles primarily, and also by interaction with thermal particles at high wave numbers. Dr. Hasegawa's effects seem to be relevant to the high wave number regime.

Sturrock: What determines which turbulence spectrum (Young, Kolmogorov, Kraichnan) is appropriate?

Henriksen: Kolmogorov appears to be the fully developed initial range of turbulence before equiparition with the magnetic field has occurred on all scales of interest. We know (e.g. De Young, Ap. J. 241, 81, 1980) that this equiparition develops first at small scales and only later at larger scales. Thus Kraichnan turbulence is 'old'. The young turbulence represents a state where relatively more energy is in the larger spatial scales.

Kennel: Can you tell anything about the variety of impulsive acceleration mechanisms by observing the spectral index, and its variation, downstream of the acceleration region?

Henriksen: Yes. We predict the range $0.5 \leq \alpha \leq 1.0$, with the normal condition being $\alpha \simeq 0.75$. This is frankly the main reason for studying the mechanism. We don't expect significant downstream variation in α (except for localized anomalies) until synchrotron losses dominate and have depleted the store of high energy particles. The 'trickle down' of the reservoir is the 'relaxation' mode that can give $\alpha = 0.5$ in the absence of strong turbulence.