COSMIC RAY ACCELERATION BY LARGE SCALE GALACTIC SHOCKS

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1/ Introduction. It is now ten years since we first heard, at an ICRC, that the mechanism of diffusive shock acceleration may account for the existence of galactic cosmic rays (Axford, et al. 1977 ; see also Rapporteur paper by Cesarsky, 1977) ; and that the first reports on this subject were published all over the world. Our understanding of this mechanism has greatly progressed since then ; detailed application to stellar wind shocks (Cesarsky and Montmerle 1983, Webb et al. 1985) and especially to supernova shocks (Blandford and Ostriker 1980, Axford 1980) have been developed.

Existing models can usually deal with the energetics or the spectral slope, but the observed energy range of cosmic rays is not explained. In particular we had shown (Cesarsky and Lagage 1981, Lagage and Cesarsky 1983) that supernova shocks cannot accelerate protons beyond $\sim 10^{14}$ eV ; this, limitation was brought up by shock lifetime and by geometrical consideration on the limited size of the shock. Therefore it seems worthwhile to examine the effect that large scale, long-lived galactic shocks may have on galactic cosmic rays, in the frame of the diffusive shock acceleration mechanism.

The spiral density wave theory predicts the existence of large scale galactic shocks, but these shocks are slow, with velocities u in the few tens of km/s range (Roberts et al. 1975). Large scale shocks may also surround evolving stellar associations, and lead to the formation of the HI supershells observed in our and other galaxies (Bruhweiler et al. 1980) ; here again, except for ~ the first million years, shocks are slow, with velocities u well below 100km/sec (Tomisaka et al. 1981). The diffusive shock acceleration mechanism only operates if $u > v_A$, where v_A, the Alfven velocity in the ionized part of the medium of density n_e cm^{-3}, is equal to 2.2 $B_{\mu G}/\sqrt{n_e}$ km/sec. This condition is difficult to fulfill for slow shocks, and consequently we disregard them here.

Large scale fast shocks can only be expected to exist in the galactic halo. We consider three situations where they may arise : expansion of a supernova shock in the halo, galactic wind, galactic infall ; and discuss the possible existence of these shocks and their role in accelerating cosmic rays.

2. Supernova shocks in the galactic halo. Since a shock propagating down a density gradient, towards lower densities, can be accelerated, Axford (1980) conjectured that shock waves, resulting from supernovae exploding at some height above the galactic disk, wander through the halo at larger speeds.

Falle et al. (1984) have made numerical studies of strong explosions in plane stratified media. They have found that the results concerning the acceleration of the rising shock are very well reproduced by an approximate analytical solution developed by Lambach and Probstein in 1969. Assuming an exponential density profile of scale height l, the results can be written in terms of dimensionless variables : τ, proportional to time, and Z, proportional to the shock position. The shock velocity is then represented by $U = dZ/d\tau$. Indeed, as expected, the shock starts to accelerate at a certain height (Figure 1). However, the space above the galactic disk is filled with a more tenuous gas, of scale height l>>l, so that the shock is not likely to be reaccelerated to velocities as high as implied by the results presented by Falle et al. A more realistic representation of the behaviour of density with height is given by : $n=n_0 e^{-z/l}+n_1$, with $n_0=1$ cm^{-3} and l=100pc. We have calculated U(τ) in this case, following the method of Laumbach and Probstein, for two cases : $n_1=10^{-3}$ cm^{-3} and $n_1=10^{-6}$ cm^{-3}. As shown in fig. 1, the shock is never accelerated if $n_1=10^{-3}$ cm^{-3} ; even $n_1=10^{-6}$ cm^{-3} is sufficient to stop the acceleration of the shock after some time. Thus the maximum rigidity that can be attained by cosmic rays is only increased by a small factor.

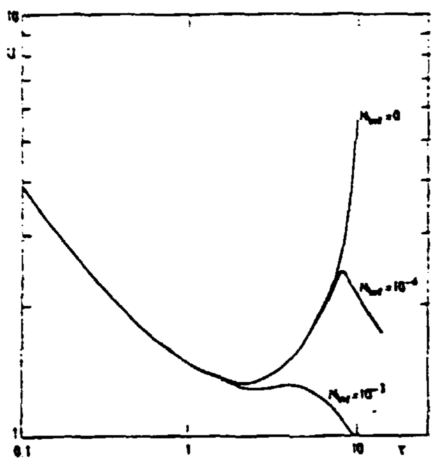

Fig. 1. Propagation of a supernova shock in the halo.

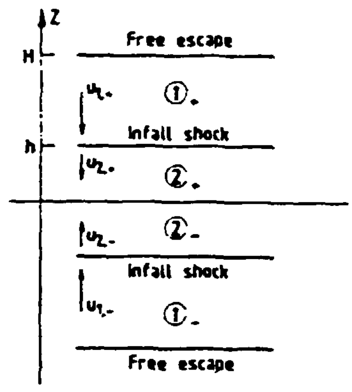

Fig. 2. Geometry for the infall model.

3. Terminal shock of a galactic wind - or galactic infall shock ?

The nature of the halo of our galaxy is not well understood yet ; in particular, it is not known whether a great fraction of the energy expended by supernovae (~10^{42} erg/sec) has to be dissipated in the halo, powering a galactic fountain or a wind, or whether it is mostly radiated in the ultraviolet and EUV in the disk.

There are also reasons for invoking just the opposite process : infall of matter onto the galactic disk. Observationally, there are the high velocity clouds : Oort's suggestion (1970) that they are extragalactic material falling onto the galaxy has never yet been proven wrong. Mirabel and Morras (1984) find evidence for a net influx of at least 0.2 M☉/yr towards the galaxy, with velocity >140 km/s ; the energy budget involved is not far from that of supernovae. The present limits on the X-ray luminosity of the halo of our galaxy and of several similar galaxies are compatible with such an accretion rate, especially considering that the wavelength range at which most of the energy is radiated is very model dependent (Cox and Smith 1976).

a) Galactic wind.

Let us consider particles diffusing in the vicinity of a spherical shock. We assume that the magnetic field is axially symmetric, so that only two components of the diffusion tensor must be defined : $K_{/\!/}$ and $K_\perp$ along and across the field. In a direction perpendicular to the shock, the diffusion coefficient is : $K_r=K_{/\!/}\,x^2+K_\perp(1-x^2)$, where x=cos θ and θ is the angle between the radial and the magnetic field direction.

The importance of the curvature of a shock has been considered by many authors in two cases : supernova shocks (Prischep and Ptuskin 1981 ; Krymsky and Petukhov 1980) and stellar winds (Webb et al. 1985 ; see also Drury 1983). A general conclusion is that the acceleration is quenched when on either side of the shock, $K_r>R_s\,u$ (1), where R_s is the shock radius and u the shock velocity : since in general K_r is expected to increase with energy this condition limits the energy that can be attained by particles accelerated by the shock.

(i) Bohm limit. Upstream of a shock, the anisotropy of cosmic rays (~u/c) leads to prompt generation of resonant hydromagnetic waves, (Bell 1978) ; therefore, the $//$ diffusion coefficient may be very small, and close to its minimum, K_B, corresponding to a scattering mean free path equal to the Larmor radius. In that case, $K_r=K_{//}=K_\perp=K_B$. For supernovae, the energy limit due to inequality (1) is then always less significant than that due to arguments on the supernova lifetime (Lagage and Cesarsky 1983).

Jokipii and Morfill (1985) assumed that the galaxy emits a strong wind, and considered cosmic ray acceleration by the shock at its boundary. They assumed that $K=K_B$, and calculated maximum rigidities using the lifetime argument. But they neglected condition (1), which, in this case, is more stringent than lifetime, so that the maximum rigidities R_{max} which they list are too high by ~two orders of magnitude (e.g. for protons, if the shock radius is 100Kpc and u_1=500km/sec, $R_{max}<2.10^{16}$ eV and not equal to 3.10^{18} eV).

(ii) $K_{//} \neq K_B$. However, as noted by these authors in a subsequent paper (Jokipii and Morfill 1986), higher energies may be obtained if $K_{//} \neq K_B$. For scattering due to resonant interactions with hydromagnetic waves the quasi-linear theory predicts that : $K_\perp = K_{//}/1+(K_{//}/K_B)^2$.
Thus, for a given field inclination, it is easy to define the diffusion tensor that maximizes R_{max}: $K_{//}=K_B[\ (9x^4-10x^2+1)^{1/2}-(3x^2-1)]^{1/2}/x\sqrt{2}$.

For $\theta>70°$, $K_{//}$ can be greater than K_B. For θ~90°, $K_{//}$~K_B/x. The angle θ is limited to θ_{max}=arc tg (c/u), which is the maximum θ for which it is possible to find a frame where the electric field vanishes. Thus the maximum rigidity can be increased with respect to the Bohm limit value, by at most a factor ~ c/u, yielding R_{max}~$6.10^8 R_s(\mathrm{Kpc}) B_{\mu G}$, independently of u. Indeed, Jokipii and Morfill (1986), who consider various cases and treat them numerically, find R_{max} (protons) ~10^{19} eV, which is as expected for an optimum magnetic field inclination.

A problem with a galactic wind shock is that the galactic disk is upstream of the shock ; thus it only sees the cosmic rays from the shock precursor. If K_r is energy dependent, only the cosmic rays of highest energy reach the galactic disk.

b) Galactic accretion shock.
If instead the galaxy is surrounded by an accretion shock which accelerates cosmic rays, the disk is downstream from the shock and therefore it receives the whole flux of accelerated particles, further compressed and energized while they travel in the halo.

We have used a formalism developed by Cox and Smith (1976) to check that a shock does develop for a wide variety of condition. The height of this shock is relatively small : ~10Kpsec. In the case of an accretion shock, the cosmic ray flux at the galactic disk depends not only on the shock characteristics and geometry, but also, or even mainly, on the values of the diffusion coefficients throughout the volume considered, and on the boundary conditions. We illustrate the type of solutions that can be obtained by an example.

As the scale of the shock is expected to be less than the radial extension of the disk, we have assumed a cylindrical geometry (see Fig. 2). Particles of momentum p_0 are injected at a rate Q into the accretion shock (z = h). Then, they are subjected to diffusion, convection, acceleration and they eventually can leave the system either via nuclear losses in the disk (rate $v_1.\delta(z)$) or by escaping the galaxy at a free boundary located at z = H. Let N(z,p) be the differential energy spectrum of the accelerated particles ; the equation governing the propagation of these particles is the usual one:

$$\frac{\partial}{\partial z}(K_i.\partial N/\partial z - u_i.N) + 1/3.\frac{\partial u}{\partial z}.\frac{\partial}{\partial p}(pN) - N.v_1.\delta(z) + Q.\delta(p-p_0).\delta(z-h) = 0$$

where K_i and u_i are respectively the diffusion coefficient and the convective velocity of region i.

Out of the shock, the solution reads : $N = A_i.\exp(u_i.z/K_i) + B_i$

The integration constants A_i and B_i are calculated via the boundary conditions : at z=H $N_1=0$; at z=h $N_1=N_2$ and

$$[K_1.\frac{\partial N_1}{\partial z}-u_1.N_1-K_2.\frac{\partial N_2}{\partial z}-u_2.N_2-1/3.(u_1-u_2).\frac{\partial}{\partial p}(p.N_1)]_{z=h}=-Q\delta(p-p_0) ;$$

$$\text{at } z=o, [-u_1.N_2-u_2.N_2-2./3.u_2.\frac{\partial}{\partial p}(p.N_2)-v_1.N_2]_{z=0}=0$$

These conditions lead to the following system of equations :

$$\alpha_1.p.A_2+\beta_1.p.B_2 = d(p.A_2+p.B_2)/d\mathrm{Log}(p/p_0)$$
$$\alpha_2.p.A_2+\beta_2.p.B_2 = d(\gamma_2.p.A_2+p.B_2)/d\log(p/p_0)-3/(u_2-u_1).Q.\delta(p-p_0)$$

where:

$$\alpha_1 = 3/2.v_1/u_2,\ \alpha_2 = -3/(u_2-u_1).[u_1.\exp(u_2.h/K_2)/(1-\exp(u_1.(h-H)/K_1) +1/3.(u_2-u_1).p.\frac{\partial}{\partial p}\exp(u_2h/K_2)]$$

$$\beta_1 = 3.(1+v_1/(2.u_2)),\ \beta_2 = 3/(u_2-u_1)).[u_2-u_1/(1-\exp(u_1.(h-H)/K_1)]$$

In the case K independent of momentum the system can be solved analytically ; the solution is :$pN = 9/(u_2-u_1)/(m_1-m_2).Q.[(p/p_0)^{m_1}-(p/p_0)^{m_2}]$,where:

$$m_{1,2} = -(\alpha_2-\beta_2-\alpha_1+\gamma_2.\beta_2)\pm[(\alpha_2-\beta_2-\alpha_1+\gamma_2.\beta_2)^2-4.(1-\gamma_2)(\alpha_1.\beta_2-\alpha_2.\beta_1)]^{1/2}$$

The velocity entering in the term describing losses can be estimated as $h_{disk}.\sigma_p.c$ whose numerical value is about 3.10^4 cm/s. When convection dominates, it has to be compared to u_2 ($\sim 5.10^7$ cm/s) and to h/K_2 when diffusion dominates. Then, in the model considered, nuclear losses do not destroy particules efficiently . If, furthermore, $(-u_1.H/K_1)$ is large ,m_1 simplifies to 0 and ,as soon as p is greater than a few p_0, N is proportional to p^{-1} .

When K depends on p, we expect that as long as $(-u_2.h/K_2)$ remains large, the spectrum is a power law of index 1. At higher energies, the spectrum depends on the shape of K(p).

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