Particle Acceleration and Turbulence in Cosmic Ray Shocks: Possible Pathways beyond the Bohm limit

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Abstract: The diffusive shock acceleration is discussed in terms of its potential to accelerate CRs beyond the “knee”. One idea is to resonantly generate a turbulent magnetic field via accelerated particles much in excess of the background field. We identify difficulties of this scenario and suggest two separate mechanisms.

The first mechanism is based on a nonlinear modification of the flow ahead of the shock supported by particles already accelerated to some specific (knee) momentum. The particles gain energy by bouncing off converging magnetic irregularities frozen into the flow in the shock precursor and not so much by re-crossing the shock itself. The acceleration rate is determined by the gradient of the flow velocity and turns out to be formally independent of the particle mean free path. The velocity gradient is set by the knee-particles. The acceleration rate of particles above the knee does not decrease with energy, unlike in the linear acceleration regime. The knee (spectrum steepening) forms because particles above it are effectively confined to the shock only if they are within limited domains in the momentum space, while other particles fall into ”loss-islands”, similar to the ”loss-cone” of magnetic traps. This also maintains the steep velocity gradient and high acceleration rate.

The second mechanism is based on the transfer of particle generated Alfven waves to longer scales via interaction with strong acoustic turbulence in the shock precursor. This work is reported in a separate paper.

Introduction

The production of CRs above the “knee” (about 10^{15}eV) in the same accelerator as its part below the knee remains a popular topic but a difficult goal [9, 8, 15]. Of course, SNRs are almost certainly responsible for the CR (at least electron) acceleration below the knee, which is documented in several ways. Recent high resolution observations with HESS [2] are also consistent with the presence of a break in the spectrum in the TeV energy range. The spectral break must reveal valuable information about acceleration in general. This include, but is not limited to, quasi-abrupt changes (at the break energy) in (i) dynamics of resonant waves that confine particles [12, 13, 5] (ii) acceleration regime [8], and (iii) particle confinement regime [11].

The important quantities that regulate DSA are the strength and spectral distribution of the turbulent magnetic field, δB. It is usually assumed to saturate at the level of the ambient field, $\delta B \sim B_0$, which thus produces pitch-angle scattering at the rate of Ω (the gyrofrequency), which limits the particle mean free path along the field to a distance of the order of gyroradius. This constitutes the so called “Bohm diffusion limit”. Under these circumstances, the mean field B_0 sets the acceleration rate and the maximum particle energy. Due to the resonance condition $kp = const$, confinement of higher energy particles requires that *longer waves must be excited*.

In order to accelerate beyond the knee several suggestions have been made. Most of them invoke the generation of $\delta B \gg B_0$ [3, 8, 15, 16]. Physically, such generation is deemed possible since the free energy source is the pressure gradient of accelerated particles, which may reach a significant fraction of the shock ram energy. An important question that still remains, however, is *how realistic the* *high saturation level is*? Earlier studies [17] and [1] suggest that due to particle trapping and turbulent mirroring the instability saturates at levels $\delta B \sim B_0$. Thus, the questions of the saturation level and the mechanism of confinement of high energy particles remain unanswered.

Acceleration in the CR precursor

To demonstrate why and when the mechanism that we suggest becomes faster than the standard one, we first consider why the latter is slow. Upon completing one acceleration cycle, i.e., crossing and re-crossing the discontinuity, a particle gains momentum

$$\frac{\Delta p}{p} \sim \frac{\Delta U}{c} \tag{1}$$

where ΔU is the relative velocity between the upstream and downstream scattering centers $\Delta U = U_1 - U_2 \sim U_1$. Thus, the acceleration requires the number of cycles $N_{cycl} \sim c/U_1 \gg 1$. The particle acceleration time can be estimated as

$$\tau_{acc} \simeq \frac{\kappa\left(p\right)}{U_1^2} \sim \lambda c/U_1^2 \sim \tau_{col} c^2/U_1^2 \tag{2}$$

where λ and τ_{col} are the particle mean free path (m.f.p.) and collision time, respectively. It is also useful to note that the acceleration time is of the order of time needed to the flow to cross the diffusion zone ahead of the shock, $\tau_{acc} \sim L_{dif}\left(p\right)/U_1$, where $L_{dif} = \kappa\left(p\right)/U_1$ is the particle diffusion length. Hence

$$\tau_{acc} : \tau_{cycl} : \tau_{col} \sim \frac{c^2}{U_1^2} : \frac{c}{U_1} : 1$$

which means that out of the c^2/U_1^2 wave-particle collisions, needed to gain a momentum $\Delta p \sim p$, only c/U_1 are productive in terms of the energy gain. Most of the collisions are wasted. In addition, $\tau_{col} \propto p$ growth with momentum and the acceleration slows down.

Fortunately, the curse of the long cycle ($\sim$upstream residence time of particles) becomes a blessing when the number and energy of accelerated particles increase. They diffuse ahead of the shock and the plasma flow develops an extended CR precursor (CRP, thereafter) of the length $L_p \sim \kappa\left(p_*\right)/U_1 = L_{dif}\left(p_*\right)$, where p_* is the particle momentum corresponding to the maximum contribution to the pressure of accelerated particles. The flow in the CRP gradually slows down towards the main shock (subshock) and the acceleration occurs primarily in the precursor precisely due to the increasingly long upstream residence time. The acceleration rate can be also obtained from the standard diffusion convection equation

$$\frac{\partial f}{\partial t} + U\frac{\partial f}{\partial z} - \frac{\partial}{\partial z}\kappa\frac{\partial f}{\partial z} = \frac{1}{3}\frac{\partial U}{\partial z}p\frac{\partial f}{\partial p} \tag{3}$$

In the smooth part of the flow we thus have for the acceleration rate

$$\dot{p} \simeq -\frac{1}{3}p\frac{\partial U}{\partial z} \tag{4}$$

if one considers characteristics of eq.(3) ignoring diffusion effects [14]. An important consequence of the last equation is that the acceleration rate is formally independent of the particle diffusion coefficient (or m.f.p.) as opposed to the linear case where such dependence is the main factor making the acceleration slow. The caveat, however, is that the derivative $\partial U/\partial z$ in eqs.(4) and (3) still depends on κ. Indeed estimating $\partial U/\partial z$ as

$$\frac{\partial U}{\partial z} \simeq \frac{U_1}{L_p} \simeq \frac{U_1^2}{\kappa\left(p_*\right)} \tag{5}$$

one sees that the acceleration time is $\tau_{acc} \sim \kappa\left(p_*\right)/U_1^2$. However, the acceleration rate does not depend on the current particle momentum as in the linear case but only on p_* i.e., pressure dominant momentum. Therefore, particles with momenta $p \sim p_*$ are accelerated approximately at the same rate as in the linear theory while particles with $p > p_*$ accelerate faster rate than in linear theory, and may in principle be *accelerated faster than the Bohm rate* and, most importantly, at a *momentum independent rate*. The overall acceleration process is illustrated in Fig.1, depending on whether or not the pressure dominant momentum p_* is stopped from growing after some $t = t_*$. If it is, particle momentum grows exponentially, otherwise it continues to grow linearly in time.

One simple reason for the saturation of $p_*\left(t\right)$ is a geometrical one. The thickness of the CR cloud ahead of the subshock, $L_p\left(p_*\right)$ cannot exceed the accelerator size, i.e., some fraction of the shock radius R_s, in a SNR, for example, e.g., [6, 4]. Thus we may identify p_* with the maximum momentum achievable in a SNR but limited by the geometrical constraints not by the lifetime of the SNR. What is also required here is that, along with or, independent of this geometrical limit to p_*, the character of *particle confinement* to the shock *changes* at p_* when p_* reaches some critical value. This should lead to a break in the spectrum at $p \sim p_*$ such that its slope at $p > p_*$ is steeper than p^{-4} and flatter at $p < p_*$. Then the main contribution to the particle pressure would come from particles *around* p_*. A *gas of relatively weak shocks* that emerges in the precursor as a result of acoustic Drury instability [7, 10] provides the required scattering environment. The key here is mirroring/trapping of particles that leads to Levy Flight (fractional kinetic) rather than diffusive-type particle kinetics [11]. An example of such kinetics is shown in Fig.2.

The fundamental reason for the formation of the spectral break ("knee") at $p \sim p_*$ is that particles with $p > p_*$ are effectively confined to the shock precursor only while they fall within limited domains in momentum space, while other particles fall into "loss-islands", similar to the "loss-cone" of magnetic traps. This complex structure of momentum space is due to the roughness character and distribution of the scattering magnetic irregularities (i.e. discontinuities, shocklets, mirrors) which leads to fractional particle kinetics. The losses steepen the spectrum above the break at p_*, which also prevents the shock width from increasing with the maximum particle energy that may continue to grow. The maximum momentum is estimated to be [11]

$$p_{max} \sim \frac{c}{u_{sh}} \frac{L}{L_p} p_* \qquad (6)$$

where L/L_p is the ratio of the distance between the scattering centers (weak shocks) to the precursor length (roughly estimated to be $\lesssim 1/10$) and p_* is the maximum momentum achieved during the standard phase of the DSA, which becomes the break point or "knee" of the final spectrum. The spectral index between p_* and p_{max} is steeper than the "standard" p^{-4} and its slope depends on details of particle interaction with scatterers. When $p_{max}(t)$ becomes $> p_*$ the acceleration regime

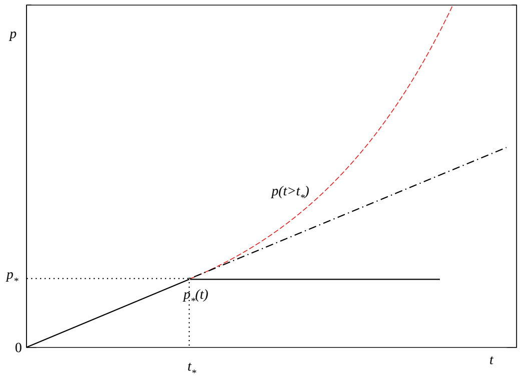

Figure 1: Schematic representation of acceleration process. Up to the time t_* the acceleration proceeds at the standard rate, which is similar during both the linear and nonlinear phases (solid line). The particle pressure-dominant momentum $p_*(t)$ grows as the maximum momentum of the entire spectrum. If for $t > t_*$ the pressure containing momentum p_* remains flat (solid line), particles with $p > p_*$ maintain the same acceleration rate that they had when their momentum was equal to p_* (dashed line). The dash-dotted line shows for comparison, how the maximum momentum $p_{max} = p_*$ would continue to grow after $t = t_*$, had $p_*(t)$ not stopped growing at $t = t_*$.

changes from the linear to exponential, so that the acceleration time is

$$\tau_{acc}(p_{max}) \sim \tau_{NL}(p_*) \ln \frac{p_{max}}{p_*} \qquad (7)$$

where $\tau_{NL}(p) \simeq 4.2\kappa_B(p)/u_{sh}^2$ (with the Bohm diffusion κ_B) is the nonlinear acceleration time [14].

Summary and Discussion

The principal conclusion of this paper is that particle acceleration inside the cosmic ray shock precursor may very well be faster and proceed to higher energies during SNR shock evolution than the conventional DSA theory suggests. The requirement for such enhanced acceleration is some limitation on the momentum of particles that contribute the most to the CR pressure in order to prevent further inflation of the shock precursor. This

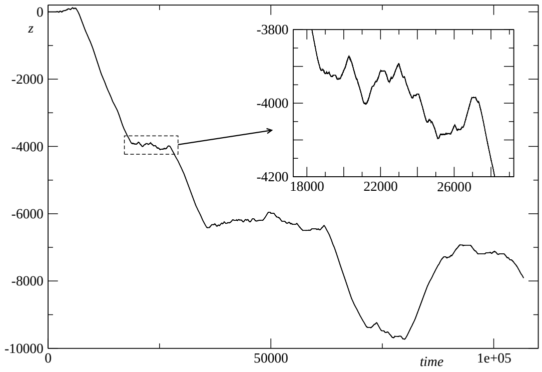

Figure 2: Particle trajectory represented as $z(t)$. It starts and remains long close to the origin ($z = 0$). Such long rests will be repeated many times at different locations. Overall, particle transport is organized in clusters (five such clusters are shown) where the transport is strongly suppressed. The clusters are connected with long jumps where the transport is ballistic.

can be achieved by spatial limitations of momentum growth, as discussed in detail in e.g., [4], by the wave compression and blue-shifting of the wave spectrum in the CRP [12], and by the change of acceleration regime after the end of the free expansion phase, [8]. As a result of these limitations, a break on the particle spectrum is formed and maintained at $p = p_*$ beyond which particles do not contribute to the pressure significantly, but are still accelerated at the same rate as particles with momenta $p < p_*$. The total acceleration time up to the maximum momentum p_{max} (given by eq.[7]) is only logarithmically larger than the acceleration time to reach the spectral break (knee) at $p \simeq p_*$ and the maximum momentum itself is also pinned to the break momentum p_* through eq.(6). The break momentum is simply the maximum momentum in the standard DSA scheme, whatever physical process stops it from growing further.

Acknowledgements

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