J. Astrophys. Astr. (2011) **32,** 427–435

Particle Acceleration at Shocks: Insights from Supernova Remnant Shocks

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Abstract. I review some basic properties of diffusive shock acceleration (DSA) in the context of young supernova remnants (SNRs). I also point out some key differences with cosmological, cluster-related shocks. DSA seems to be very efficient in strong, young SNR shocks. Provided the magnetic fields exceed some hundreds of μGauss (possibly amplified by CR related dynamics), these shocks can accelerate cosmic ray hadrons to PeV energies in the time available to them. Electron energies, limited by radiative losses, are likely limited to the TeV range. Injection of fresh particles at these shocks is poorly understood, but hadrons are much more easily injected than the more highly magnetized electrons. That seems supported by observational data, as well. So, while CR protons in young SNRs may play very major roles in the SNR evolution, the CR electron populations have minimal such impact, despite their observational importance.

Key words. Particle acceleration—supernova remnant shocks—galaxy cluster shocks.

1. Introduction

Much of this workshop is focused on the role of high energy particles or cosmic rays (CR) on the scales of galaxy clusters and beyond. Currently the only convincing direct evidence for the existence of such CR comes from radio electron synchrotron emission; e.g., halos and relics. The origins and energy sources for those electrons are not yet clear. Both first order acceleration at shocks and second order acceleration in cluster turbulence are candidate contributors (e.g., Brunetti & Lazarian 2007; Pinzke & Pfrommer 2010). A key step in understanding how those environments contribute in the cosmological context is to understand how the relevant physics works generally, of course. The best shock acceleration laboratories we have outside the heliosphere are supernova remnant shocks. Young supernova remnants (SNR), and Type Ia SNR, in particular, would seem to be the best candidates, since Type Ia events are less likely to involve complex circumstellar environments and there is no relic central compact object (e.g., a pulsar) to confuse the picture. We should, of course, keep in mind some potentially important distinctions between SNR shocks and cosmic structure shocks. For example, although both types of shocks will involve shock speeds $\sim$1000 km s^{-1}, young SNR shocks are typically very high Mach number, while the most important structure shocks in terms of energy dissipation generally have relatively low Mach numbers (e.g., Ryu *et al.* 2003).

In this spirit, I review below some elements of what we do understand about CR acceleration in strong SNR shocks as well as some of the important physics and astrophysics that is less well established. I start with a very short overview of the so-called Diffusive Shock Acceleration (DSA) physics, followed by a summary of some key features of DSA in SNRs that we do and do not know. In this I emphasize differences between CR hadrons (mostly protons) and CR leptons (mostly electrons). This is important, since virtually all the direct information we have about CR in SNRs is information about the CR electrons, while most of the energy transferred by the SNR shock to CR should go to the CR hadrons. Finally, I will touch on the issue of magnetic field amplification in CR-modified shocks, which is currently thought to be as an important feature of such shocks.

2. A cartoon view of DSA

DSA is the commonly accepted paradigm of CR acceleration in SNRs. Over the past quarter century DSA theory has progressed substantially. The basic outline is well established, although there remain important holes in our understanding that impact our ability to predict and to understand relevant observations. There are a number of good reviews of DSA (e.g., Malkov & Drury 2001). Here I provide only a short, cartoon picture of the essential physics.

The underlying key DSA concepts as illustrated in Fig. 1 are: (1) In collisionless, astrophysical plasma shocks the bulk of the plasma is approximately thermalized through collective interactions on scales much less than the collision lengths and times for particles much faster than the bulk flow through the resulting shock transition (roughly the postshock sound speed in a strong shock). (2) Sufficiently suprathermal charged particles, with their larger interaction lengths, can pass through the dissipative shock transition (thickness δ). However, the CR are likely to be scattered upstream (region 1) or downstream (region 2) within a length scale $x_{D,1,2} = \kappa(p)_{1,2}/u_{1,2}$ into trajectories that take them back through the shock transition. In this expression $\kappa(p) = (1/3)\lambda v$ is the spatial diffusion coefficient of a particle with scattering length λ and velocity v, while u is the flow speed of scattering centers with respect to the shock. (3) Each pair of shock crossings gives such a CR particle a fractional momentum boost that is proportional to the velocity difference between upstream and downstream scattering centers, $\Delta u = u_1 - u_2$; that is, $\Delta p = (4/3)(\Delta u/v)p$. (4) The CR, although a very small fraction of the plasma particles, can collectively contribute substantial pressure due to their very large individual energies. Indeed, in *strong* shocks, even though the CR probably represent less than 0.1% of the number flux through the shock the CR pressure at the shock, dominated by relativistic protons, can approach the momentum flux into the shock; i.e., $P_{CR} \sim \rho_1 u_1^2$. (P_{CR} will be much smaller than this in weak shocks, however (e.g., Kang & Ryu 2011).) Because of their large scattering lengths, the CR diffuse upstream a characteristic distance $x_{D,1}$, producing a shock pressure precursor that

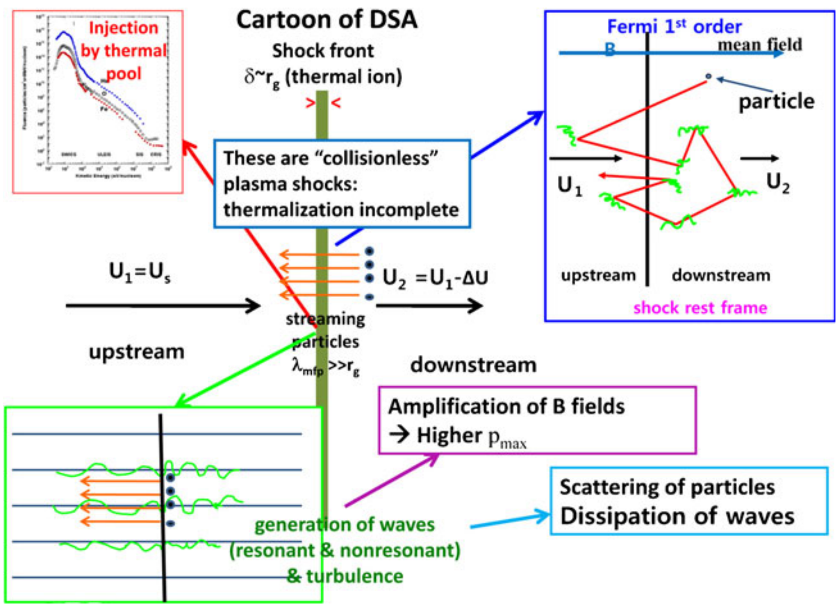

Figure 1. Cartoon of diffusive shock acceleration.

decelerates and heats the inflowing plasma, while enhancing turbulence and magnetic fields. The precursor modifications also add to the total compression through the full shock transition, while weakening the gas subshock on account of a hotter, slower flow going into the subshock.

The local magnetic field plays an essential role in all this, since particle trajectories, including scattering, are mediated by the field. Details depend on the field geometry, but at the simplest level the key parameter is the particle gyro-radius, $r_g = (p_\perp/m)(1/\omega_B) = R(p)/B$, where $\omega_B = ZeB/mc$ is the nonrelativistic cyclotron frequency of a particle with charge $q = Ze$ and $R(p) = p_\perp c/Ze \sim \sqrt{E_k^2 + 2E_k mc^2}/(Ze)$ is particle rigidity. The rigidity parameter is an especially convenient metric to distinguish particle behaviors. The thickness, δ, of the bulk plasma transition or 'gas subshock' will be of the order of the (post-shock) thermal proton gyroradius, while gyroradii (and resonant scattering lengths) for much more energetic protons will be much greater, allowing those particles to cross easily through the subshock. The downstream flow away from the shock is always subsonic with respect to the shock normal, so some suprathermal protons could easily recross the shock in the absence of magnetic fields and postshock turbulence. The magnetic turbulence greatly reduce the chances for protons to recross (e.g., Malkov 1998), although as their velocities increase, so do their chances to 'swim' upstream to recross the subshock to the upstream side. Those that do provide a small population of low energy CR seed particles. This so-called 'thermal leakage injection' process is moderated by the thermalization process itself (e.g., Vladimirov *et al.* 2006) as well as modifications to the shock brought on by higher energy CR. The fraction of the proton flux passing through shocks that are injected by the thermal leakage process is not entirely clear and probably depends on both the shock strength and the geometry of the magnetic field. Simulations indicate that DSA will saturate in strong shocks when the injected fraction exceeds something of the order $\sim 10^{-4}$ (e.g., Berezhko *et al.* 1995), emphasizing again that only a very small population of CR is required to impact the shock dynamics.

On the other hand, thermal leakage is rigidity dependent, and for nonrelativistic particles of kinetic energy E_k thermal leakage of electrons should be negligibly small compared to protons. Electrons will generally need to have energies $\sim m_p/m_e$ higher than postshock thermal energies, so be relativistic in cases of interest before they can participate in DSA. Thus, electron CR injection is expected to be much less efficient than proton CR injection and probably requires some 'pre-acceleration'.

A further significant consequence of DSA on strong shock formation is the likelihood that CR streaming against the plasma flow into the shock will amplify magnetic field fluctuations through resonant and nonresonant instabilities (e.g., Bell & Lucek 2001; Bell 2004). The amplification may, in fact, be fast and strong enough to enhance the effective magnetic field by one or more orders of magnitude over the upstream magnetic field. Such field amplification may be essential in allowing SNRs to accelerate CR to the PeV energies needed to explain the break in the observed CR spectrum at the so-called 'knee', and also to account for thin X-ray synchrotron rims observed in several young SNRs.

3. The DSA CR spectrum in SNRs

The rate at which a particle gains momentum through DSA is $\dot{p} \approx \Delta p/\tau_d$ where τ_d measures the time between successive upstream shock crossings; $\tau_d = 8\langle\kappa/(uv)\rangle$, with the angular brackets indicating an upstream/downstream averaged κ/u. (So, $\dot{p} \propto p\Delta u\langle u/\kappa\rangle$.) At the same time energetic particles escape through downstream diffusion at a rate that scales with $u_2/x_{D,2}$. The balance between these two rates determines the steady-state logarithmic slope of the CR momentum distribution, $f(p)$; namely $\partial \ln f(p)/\partial \ln p = -q = -3u_1/\Delta u$. For strong, unmodified shocks with weak magnetic fields $\Delta u \approx (3/4)u_1$, so $q \approx 4$. This is independent of the details of the CR scattering, *except* that the characteristic upstream turn-around distance, $x_{D,1} \propto \kappa(p)/u_1$, is momentum dependent, so in a CR modified shock Δu will be momentum dependent. In most models $\lambda(p)$ is an increasing function of p, so $\Delta u/u$ increases with p and the steady-state CR distribution function, $f(p)$, develops a concave form.

The most common diffusion model is Bohm-like diffusion, $\lambda \propto r_g$, which in the relativistic regime gives $\kappa = \eta pc^2/(ZeB)$, with $\eta > 1/3$ taken as a constant. Then assuming $\kappa_1/u_1 = \kappa_2/u_2$ and a strong, unmodified shock

$$\dot{p} \approx \frac{1}{8\eta}\left(\frac{u_s}{c}\right)^2 \omega_{B1,p} m_p c, \tag{1}$$

where u_s is the shock speed. For a steady, plane shock this allows CR to be accelerated over a time interval, t, to a maximum energy,

$$E_{max} \approx 35\ \mathrm{TeV}\frac{Z}{\eta}\left(\frac{u_s}{3000\,\mathrm{km/s}}\right)^2\left(\frac{B_1}{\mu\mathrm{G}}\right)\left(\frac{t}{1000\,\mathrm{yr}}\right). \tag{2}$$

The normalizations were chosen to reflect the fact that young SNRs in typical ISM environments (ambient densities, $n_o \sim 0.01$–1 cm^{-3}) will decelerate from initial shock speeds $u_o \sim \sqrt{E_o/M_{ej}} \sim 10^4$ km s^{-1} on time scales of a few hundred to $\sim$1000 years. On the other hand, equation (2) is not strictly applicable for a decelerating shock. The net result from equation (1) for a Type Ia SN explosion ejecting 1.4 $M_\odot$ and evolving adiabatically is a time-asymptotic maximum energy given by (e.g., Edmon *et al.* 2011)

$$E_{\max} \approx 12\ \mathrm{TeV}\frac{Z}{\eta}\left(\frac{E_o}{10^{51}}\right)^{1/2}\left(\frac{B_1}{\mu\mathrm{G}}\right)\left(\frac{\mathrm{cm}^{-3}}{n_o}\right)^{1/3}. \tag{3}$$

An insight from comparison of equations (2) and (3) is that the peak CR energy in a decelerating shock is approached relatively early in the shock evolution. Further, CRs trapped inside the SNR will cool adiabatically as the SNR expands.

Another important point from equations (2) and (3) is that for such SNR blasts to accelerate CR protons close to the 'knee', so $\sim$PeV, as seems to be necessary, either the shocks must persist at speeds exceeding several thousand km s^{-1} for much longer than a thousand years (requiring very low ambient densities) or the magnetic fields in the shocks must be much stronger than the few micro Gauss values typical of the ISM. Both potential solutions have been discussed, but because other evidence outlined in the following section also suggests that the magnetic fields in young SNRs may commonly approach milliGauss levels, much of the recent literature has focused on the possibility that CR generated instabilities in strong, modified shocks are capable of amplifying ambient fields by an order of magnitude or more (e.g., Riquelme & Spitkovsky 2009). I will return below to the issue of field amplification.

4. Electrons

Most of this discussion has centered on CR protons, since they are expected to capture most of the energy transferred to CR in shocks. In reality the existing direct evidence primarily, if not exclusively, reveals the existence of CR electrons in SNRs. Radio synchrotron emission from $\sim$GeV electrons has been observed in SNRs for a long time, and in the past decade or so compelling evidence has accumulated for keV X-ray synchrotron emission from $\sim$TeV electrons. The GeV to TeV γ-ray emissions observed in several SNRs are likely due to inverse Compton emissions from TeV electrons and/or decay of π^0 secondaries in collisions between CR protons and thermal plasma protons. The roles of the two processes are still in debate, although the γ-ray spectrum inferred by very recent Fermi detection of GeV γs from RX J1713.7-3946 (Abdo *et al.* 2011) combined with previous H.E.S.S. TeV data (Aharonian *et al.* 2006) strongly suggest inverse Compton emission from TeV electrons in that object. In particular, π^0 decays should produce a relatively strong feature in the GeV spectrum (above the pion production threshold) that is not seen in the Fermi RX J1713.7-3946 data.

If relativistic electrons are present at a shock, DSA operates on them in exactly the same way as protons of the same rigidity, R (which is the energy per unit charge when they are relativistic). On the other hand, because the electron mass is so much smaller than the proton mass, there are two important differences in the expected

CR electron population. First, since the radiated synchrotron power of a particle scales as $P \propto (Ze)^6 \; B^2 R^2/(mc^2)^4$, electrons of a given rigidity, R, will radiate their energy $\sim 10^{13}$ times more rapidly than protons of the same rigidity (energy). This strongly limits the energy of CR electrons during DSA. For Bohm diffusion the acceleration rate, $\dot{E}_{\rm DSA}$, in a strong shock according to equation (1) scales with $Zeu_{\rm s}^2 B$, independent of R. The energy at which $\dot{E}_{\rm DSA}$ balances synchrotron cooling, P, is then (e.g., Edmon *et al.* 2011)

$$E_{e,\rm max} \approx 100\ {\rm TeV}\sqrt{\frac{\mu{\rm G}}{B_1}}\left(\frac{u_{\rm s}}{3000\ {\rm km/s}}\right). \tag{4}$$

This all assumes that there is time for the electrons to reach this energy, so one must check to see that $E_{e,\rm max}$ from equation (4) is less than $E_{\rm max}$ from equation (2) or equation (3).

So, an ambient field exceeding $\sim 100\,\mu$G helps DSA to produce PeV protons in young SNRs, and hence such a field would limit the electron population there to energies below about 10 TeV. Indeed such strong fields and thus such limited electron energies have been invoked to explain the very thin keV X-ray rims in several young SNRs. For Tycho's SNR, for example, the NE rim observed by Chandra (Bamba *et al.* 2003) is about 0.01 pc thick, representing a fractional thickness, $\Delta r/r \sim 0.002$. Assuming this represents the distance electrons radiating keV photons can travel downstream of the shock, Völk *et al.* 2005 estimated a postshock magnetic field $\sim$200–300 μG in Tycho. The associated electron energies would be (equation (4)) $\sim$1–10 TeV. Note the curious property of equation (4) that the highest frequency synchrotron emission, $\nu_{s,\rm max} \propto E_{e,\rm max}^2 B$, is *independent* of the magnetic field strength. This is a well-known property of DSA with Bohm diffusion, so it is applicable to strong shocks anywhere, and not just to SNRs.

The second important consequence of the small electron/proton mass ratio, as mentioned previously, is the resulting difficulty of electron injection from the thermal population at the shock. Electrons clearly are injected somehow from thermal plasma during passage through a SNR shock transition, although the number of such electrons is expected to be very much smaller than the number of protons. In a simple thermal leakage model that number would be as much as a factor $\sim \exp(-m_{\rm p}/m_{\rm e})$ smaller, so negligible.

Several alternate electron injection ideas have been presented that depend on some kind of pre-acceleration to increase rigidities for some electrons. For example, in some models ions reflected off the shock back into the upstream plasma generate right-handed Whistler waves that can resonantly interact with electrons in the shock precursor (e.g., Levinson 1996). The wave electric field component aligned with the large-scale magnetic field may be able to accelerate some electrons to sufficient rigidities to pass through the sub-shock transition and so 'inject' them into the CR population. An alternate approach to electron injection by Morlino (2009) utilizes the fact that most heavy atomic species entering a shock from the ISM will only be partially ionized. If those ions are accelerated by DSA to mildly relativistic energies before they are collisionly or photo-ionized to higher states, the electrons ejected in subsequent ionization events, which would be relativistic in the shock frame, can serve as a seed CR electron population. Morlino estimates that this can provide a seed electron population as large as $\sim 10^{-4}$ the seed proton CR population.

To account for observed synchrotron emissions from young SNRs with such small electron injection efficiencies, models typically call for magnetic fields interior to the forward shock in excess of 100 μG (e.g., Edmon *et al.* 2011). This is consistent with our previous discussion. It should be noted, however, that CR electron populations this small will not produce sufficient inverse Compton emission to account for the GeV and TeV γ-rays seen in several remnants. In that event photons from π^0s produced by CR proton collisions with thermal plasma provide an alternate source of γ-rays and, indeed, have been sought as the 'smoking gun' for DSA in SNRs for some time (Drury *et al.* 1994). As noted earlier, however, recent observations of RX J1713.7-3946 seem to require an inverse Compton γ-ray explanation, so a more substantial CR electron population and a weaker magnetic field than suggested above (e.g., Zirakashvili & Aharonian 2010). But, even here, the CR electron/proton ratio is probably less than 1%. Thus, in young SNRs CR electrons are very likely not dynamically important. For the CR electrons to become dynamically important in a collisionless shock there probably would have to be a substantial pre-existing CR electron pressure in the upstream plasma.

5. Magnetic field amplification

As outlined above, the body of evidence, while not unequivocal, mostly now points to magnetic fields inside the forward shocks of young SNRs that are at least an order of magnitude stronger than can be accounted for by simple compression of likely ambient fields. Several plasma and MHD instabilities have been suggested that may be capable of amplifying magnetic fields as they pass through a shock precursor. The best known of these was pointed out by Bell 2004, and derives from the charge imbalance inferred from the electron/proton CR ratio discussed in the previous section. With that, CR protons streaming ahead of the shock carry an electric current, which must be balanced by a return current in the thermal plasma. The associated $j \times B$ body force on the plasma will amplify fluctuations in the local magnetic field. The linear instability peaks at small scales compared to the CR gyroradius (e.g., Blasi & Amato 2007), although the scales involved grow as the instability becomes nonlinear. In idealized simulations Bell's instability seems capable of amplifying the field within a precursor by as much as an order of magnitude on fairly short time scales.

There are other CR-induced mechanisms that may also be capable of amplifying magnetic fields within modified shocks, including purely hydrodynamical processes. Beresnyak *et al.* (2009) pointed out, for example, that the CR pressure gradient within the shock precursor will rapidly amplify pre-existing turbulent motions inside the precursor. The magnetic field can then be quickly amplified by that enhanced turbulence on scales less than the diffusion scale, $x_{D,1}$ defined in §1. Ryu *et al.* (1993) pointed out some time ago that the weak diffusive coupling between CRs and thermal plasma should cause a strong shock precursor to be unstable to the amplification of inflowing turbulence.

6. Summary

Diffusive Shock Acceleration is evidently very efficient in strong shocks around young supernova remnants. There is a fairly broad consensus that protons from the overtaken ambient medium can be injected in small, but sufficient numbers to capture a major fraction of the energy dissipated in such shocks. Electrons are obviously also accelerated in these shocks, since we see clear observational evidence for their existence there. In fact, to date only CR electrons have been unequivocably identified by their emissions in SNRs. The number of CR electrons needed to explain observed nonthermal emissions in SNRs, is, however, quite small, and small enough that the dynamical influence of the electrons is probably negligible, in contrast to proton CRs.

The rate of particle acceleration depends on the velocity jump experienced by particles being accelerated and on the strength of the scattering, which will generally depend on the strength and orientation of the local magnetic field. In the SNR setting, accompanying energy losses by protons and other hadrons are negligible, so the maximum energies achieved basically depend on how fast and how long the acceleration takes place. The faster the shock and the stronger the magnetic field the higher the expected energies. Electrons, on the other hand, rapidly lose energy to synchrotron emission. The maximum electron energies will usually correspond to the energies where acceleration and cooling balance. In SNR environments inverse Compton losses will usually not be much of an issue. In cosmological settings, where the magnetic fields are probably much weaker, inverse Compton losses need to be folded in, and may actually dominate.

Shocks in young SNRs, being strong and probably CR-modified are also likely to amplify ambient magnetic fields overtaken by the shocks as a feedback aspect of efficient particle acceleration. There is good evidence that such processes may be able to amplify field strengths by an order of magnitude or more.

Cosmological shocks in and around clusters have several distinct properties that need to be remembered when comparisons to young SNR shocks are made. While both have roughly similar speeds, u_s, the cosmological shocks mostly propagate through much hotter media, so are weaker, leading to steeper CR spectra. This makes them much less efficient particle accelerators. This means, as well, that they are likely to be much less modified by CR back reaction than the stronger SNR shocks. The magnetic fields associated with cluster related shocks are poorly understood at this point, but the ambient fields are probably weaker than those encountered by most SNRs. At the same time, the fact that the cosmological shocks are not likely to develop strong CR-produced modifications means that the kinds of magnetic field amplification discussed in §5 will not take place.

Acknowledgements

I am very grateful to the organizers for the invitation and opportunity to participate in this excellent workshop. This work is supported by NASA grant NNX09AH78G and US NSF grant AST0908668. Figure 1 is an adaptation of a figure originally prepared by Hyesung Kang.

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