Cosmic plasmas: their research frontiers

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Abstract

Nonthermal particle acceleration is one of the most important topics in the physics of cosmic plasmas. Owing to rapid advances in in situ and remote-sensing observations of these plasmas, much details on acceleration processes are now being accumulated. I will review several topics from these recent observational and theoretical accomplishments and discuss their physical significance. © 2001 Published by Elsevier Science Ltd.

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1. Introduction

The universe is thought to be ‘magnetized’, fully or partially, and filled with various plasma phenomena. The cosmic magnetic field, the origin of which could be dated back even to the primordial time just after the big bang [1,2], is being generated via dynamo processes in planets, stars, and galaxies. While the cosmic plasma phenomena sometimes exhibit exotic properties like extremely strong magnetic fields beyond quantum limit (e.g. Ref. [3]), there often exist close similarities with those occurring in artificial plasmas. One of the examples is magnetic reconnection, in which magnetic energies are efficiently converted to particle energies, namely plasma thermal/kinetic energies as well as those of nonthermal particles. Extensive collaborations among theoretical, fusion, solar, and magnetospheric plasma physicists have given fruitful results in establishing the importance of reconnection in various environments [4–6]. In star forming regions, plasmas coexist with neutral gases and dust particles (e.g. Ref. [7]), where interactions among them play important dynamical roles in the angular momentum transfer process from the center to outer regions, which is thought to be an essential process for stars to be created from molecular clouds. Elementary physical/chemical interactions necessary to understand such interstellar processes are not far away from those discussed in terms of laboratory/industrial experiments on plasma–neutral beam interactions as well as plasma–dust interactions.

The last decade of the 20th century was the “golden age” of cosmic plasma research (K. Shibata, in this symposium). Its first milestone was the launch of the YOHKOH satellite, a solar-astrophysical mission jointly supported by Institute of Space and Astronautical Science (ISAS), Japan and NASA. An example of YOHKOH soft-X ray images of the sun is shown in Fig. 1a in which the bright regions in the solar corona are seen above the sunspot groups (Fig. 1b, courtesy of K. Shibata). From the high time cadence of these YOHKOH X ray observations, a new dynamical picture of the solar corona has been established. Recently, the Solar Heliospheric Observatory (SOHO) satellite has provided us wider-view images of the solar corona, and details of coronal mass ejections (CME), which are quite rich in magnetohydrodynamic (MHD) features, such as streamers, loops and helical structures (Fig. 1c from the SOHO Web site.1 See Ref. [8] for details).

In the large portion of cosmic plasmas, the usual two-body Coulomb collisions are relatively unimportant and collective interactions through long-range electromagnetic forces become dominant. This means that the physical status of cosmic plasmas can depart far away from their thermodynamic equilibrium. In fact, it is getting clearer that through collective plasma processes particles accelerated well beyond thermal energies play energetically significant roles in a variety of dynamical events in the universe, such as magnetotail substorms, planetary bow shocks, solar flares/CME and relating shocks, the solar wind terminating shock, supernova shocks, pulsars, gamma ray bursts, active galactic nuclei, and even in inter-galactic spaces within clusters of galaxies.

Fig. 2 shows the energy spectrum of nonthermal particles

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in the cosmic plasma around us, namely cosmic ray particles, which have an impressively huge energy coverage from 100 MeV (10^8 eV) to slightly less than 1 ZeV (10^{21} eV). An important aspect, at least from the plasma physics point of view, is that these particles actually interact with background thermal plasmas, although there is the difference of many decades between the energies of cosmic ray particles and thermal plasma particles.

The first feature of the interaction between cosmic ray particles and the background plasma is seen as a falloff in the lowest energy part (\( E < \sim \text{several GeV} \)) in the figure.

This is now known as the solar wind modulation effect, namely the effect of the solar wind flow on the cosmic ray transport in the heliosphere (Section 2). A power law distribution with a unique spectral index (\( \sim E^{-2.7} \)) seen in the range from \( \sim 10 \text{ GeV} \) up to \( \sim 1 \text{ PeV} (10^{15} \text{ eV}) \) is now successfully described by a diffusive shock acceleration model at supernova remnants (Section 3), which is one of the biggest triumphs in cosmic plasma physics during the last quarter of the 20th century.

Beyond 1 PeV, in Fig. 2, we see a slightly steeper slope (\( \sim E^{-3.3} \)), which is thought to reflect the effect of particle escape from our Galaxy. One of the biggest puzzles in the contemporary cosmic ray physics is the existence of ultrahigh energy cosmic ray (UHECR) particles beyond 100 EeV (10^{20} eV), which can neither be trapped in our Galaxy with its magnetic field (\( \sim \text{several } \mu \text{G} \)), nor travel beyond a cosmologically tiny distance (several tens of Mpc \( \sim 10^{26} \text{ cm} \)). The latter limitation is due to the inelastic collisions between UHECRs and background 2.7 K photons which result in rapid decelerations (so-called ‘GZK’ effect). While some authors are proposing that these UHECRs are created from some exotic entities such as cosmic strings, topological defects, or magnetic monopoles (top-down models), many people believe that they are due to acceleration processes in some extreme environments (bottom-up models) [10–12]. The origin of the UHECRs thus remains to be a challenging problem for cosmic plasma physicists in the 21st century.

The above is the most significant nonthermal acceleration phenomenon. As it should occur in the most remote regions, we cannot resolve their spatial/temporal structures in detail with current observational techniques. To single out the physical process(es) responsible for such a phenomenon, we should use theoretical techniques as much as possible. Since physics is an experimental science, we need a ‘laboratory’ by which we can calibrate our tools, which here is our theory. In the following sections, I will show several examples in which heliospheric in situ observations have played important roles of a laboratory for cosmic plasma physics.
Fig. 3. Intensity of cosmic rays on the ground (neutron monitor counting rate: a curve with dots) and sunspot numbers (bars). The numbers on the left-hand side are for cosmic rays, whereas those on the right-hand side are for sunspots.

2. Solar modulation of galactic cosmic rays

The solar modulation effect seen as a falloff in the lowest energy part \( (E \sim \text{several GeV}) \) of Fig. 2 shows the dependence of the solar activity. In Fig. 3, the time variations \((1951-2000)\) of the ground level intensity of cosmic rays, namely counting rate of the neutron monitor in Climax and sunspot numbers, are shown. These neutrons are the product of nuclear interactions between incident cosmic ray protons (or heavier nuclei) of \( \sim > \) several GeV/n and atmospheric nuclei. There is a clear anti-correlation between these variations.

As stated in Section 1, the solar wind is rich in MHD phenomena of various scales, and filled with MHD waves or discontinuities including wide wavelength ranges from the thermal ion cyclotron radius \((\text{of the order of} \sim 10^{-1} \text{ km})\) to the macroscopic scale \((1-10^5 \text{ astronomical unit}) \) \((\text{AU} = 1.5 \times 10^{11} \text{ m})\). Cosmic ray particles in the heliosphere interact with these MHD waves (of the frequency \( \omega \) and wavenumber \( k \)) through the cyclotron resonance condition:

\[
\omega - k\tilde{\omega} = \pm \Omega
\]

where \( \tilde{\omega} \) is the velocity of particles, \( \Omega \) is their cyclotron frequency (with relativistic mass correction), and the sign should be chosen according to the helicity matching condition between the particle spiral orbits and the MHD waves. For these waves, we roughly have \( \omega \sim kV_A \), where \( k = \frac{\omega}{\Omega} \) and \( V_A \) is the Alfvén velocity \((\sim 10^{-10} \text{ km/s for the nominal solar wind condition at 1 AU})\). For cosmic ray particles (of the velocity \( \sim c \)), the inequality \( |\tilde{\omega}| \gg V_A \) holds. Therefore, Eq. (1) can be approximated as

\[
k \sim \frac{\Omega}{|\tilde{\omega}|}
\]

or for the wavelength \( \lambda = 2\pi/k \),

\[
\lambda \sim \frac{2\pi|\tilde{\omega}|}{\Omega}
\]

For example, for protons of \( 5 \text{ GeV}, \quad |\tilde{\omega}| \sim c = 3 \times 10^5 \text{ km/s and } \Omega \sim 0.1 \text{ s}^{-1} \) in the magnetic field of \( 5 \text{ nT} \) (a typical value at 1 AU), so that MHD waves resonating with \( 5 \text{ GeV} \) protons have a wavelength of \( \sim 2 \times 10^{12} \) cm \( \sim 0.1 \text{ AU} \).

Transport of cosmic rays thus interacting with the solar wind is described by the diffusion–convection equation (developed by Parker, Jokipii, Axford, Gleeson and others, in the ’60s):

\[
\frac{\partial f}{\partial t} + \tilde{\omega}f = \frac{1}{3} \left( \text{div} \, \tilde{\omega} \right) \frac{\partial f}{\partial \ln p} + \tilde{\nabla}(D\tilde{\nabla} f)
\]

where \( f = f(\tilde{x}, p, t) \) is the phase space distribution function of nonthermal particles, \( \tilde{x} \), the spatial coordinate, \( p \), the magnitude of particle momentum, \( \tilde{\omega} = \tilde{\omega}(\tilde{x}, t) \), the velocity field of the background plasma, \( D = D(\tilde{x}, p, t) \), the spatial diffusion coefficient.

You might think it curious that the ‘convection term’, namely the second term in the left-hand side of Eq. (4) is with the plasma velocity, \( \tilde{\omega} \), not with the particle velocity, \( \tilde{\omega} \). In what follows, let us derive the diffusion–convection equation heuristically. For more rigorous derivation, we refer the readers to Refs. [13–16].

As the master equation, we write the Vlasov equation in the framework of the quasi-linear theory:

\[
\frac{\partial f}{\partial t} + \tilde{\omega}f + \frac{\partial p}{\partial \tilde{\omega}} \frac{\partial f}{\partial p} = \left( \frac{\partial f}{\partial t} \right)_{\text{corr}}
\]

Here, we assume that \( \frac{\partial p}{\partial t} \) in the third term of the left-hand side represents Lorentz force exerted by the averaged electromagnetic field, whereas the effective ‘collision term’ on the right-hand side, \( (df/dt)_{\text{corr}} \), describes the effect of correlations among the fluctuation term of \( f \) and various modes in the turbulent electromagnetic field (see Appendix A). We assume that particles are well pitch-angle scattered by the MHD waves in the turbulence so as to have isotropic pitch angle distribution in the rest frame of the background.
plasma. Hence, we set
\[ \tilde{\omega} \sim \tilde{u} \]  

(6)

Of course, there appears the ‘residues’, \((\tilde{\omega} - \tilde{u})\tilde{\nabla}f\), from this approximation. We take care of them by ‘renormalizing’ the right-hand side, \((df/dt)_{\text{corr}}\). Now, we have
\[ \frac{\partial f}{\partial t} + \tilde{u} \tilde{\nabla}f + \frac{dp}{dr} \frac{df}{\partial p} = \left( \frac{df}{dt} \right)_{\text{corr}} \]  

(7)

Under the condition that turbulence scatters particles isotropically, energies of the particles are subject to adiabatic deceleration/acceleration along with the divergence/convergence of the background plasma motion:

\[ \text{div} \, \tilde{u} > 0 \rightarrow \text{deceleration} \frac{dp}{dr} < 0 \]

\[ \text{div} \, \tilde{u} < 0 \rightarrow \text{acceleration} \frac{dp}{dr} > 0 \]

Similar to the elementary problem of a molecular gas confined in an externally driven piston, we have
\[ \frac{dp}{dr} = -\frac{P}{3} (\text{div} \, \tilde{u}) \]  

(8)

where the factor 1/3 results from the averaging over the pitch angle of particles. With this representation of the particle momentum change, we now have
\[ \frac{\partial f}{\partial t} + \tilde{u} \tilde{\nabla}f - \frac{1}{3} (\text{div} \, \tilde{u}) \frac{df}{\partial \ln p} = \left( \frac{df}{dt} \right)_{\text{corr}} \]  

(9)

The right-hand side of Eq. (9), \((df/dt)_{\text{corr}}\), turns out to be the spatial diffusion term (e.g. Refs. [14,15]):
\[ \left( \frac{df}{dt} \right)_{\text{corr}} = \tilde{\nabla} (D \tilde{\nabla} f) \]  

(10)

where the diffusion coefficient \(D\) generally has functional dependence on the particle species, on the position \(x\) and on the particle momentum \(p\) (and also on the time \(t\) for the nonsteady problem). With Eqs. (9) and (10), we get the diffusion–convection equation (4). It is noted that in the above \(V_A/\nu \ll 1\) is assumed and that the diffusion term in the momentum space, which is the order of \((V_A/\nu)^2\), is negligibly small.

In the late of 1970s, Jokipii and his colleagues [17] pointed out that the magnetic drift motions (gradient-\(B\) and curvature) of cosmic ray particles in the heliospheric magnetic field have significant effects on the transportation process. With the drift velocity \(\tilde{w}_D\), they modified the diffusion–convection equation (4) to
\[ \frac{\partial f}{\partial t} + \tilde{u} + \tilde{w}_D \tilde{\nabla}f = \frac{1}{3} (\text{div} \, \tilde{u}) \frac{df}{\partial \ln p} + \tilde{\nabla} (D \tilde{\nabla} f) \]  

(11)

Along with the variation of the solar activity, parameters in Eq. (11), \(\tilde{u}\), \(\tilde{w}_D\), and \(D\), are expected to be modified and then to result in the modulation of cosmic rays in the heliosphere.

3. Diffusive shock acceleration

In 1977–1978, several authors [18–21] have published almost simultaneously their ideas on the shock acceleration, now called in terms of diffusive shock acceleration. Here, we follow the discussion by Blandford and Ostriker [21]. Their starting point is the diffusion convection Eq. (4) initially developed to explain the solar modulation of cosmic ray particles. What is new in their idea is to apply the momentum change term \((\text{div} \, \tilde{u} / 3) (\partial f / \partial \ln p)\) in Eq. (4) to the shock front environment. In the shock rest frame,\(^2\) the plasma velocity \(u_1(=|\tilde{u}_1|)\) in the upstream region \((x < 0)\) is supersonic, whereas \(u_2(=|\tilde{u}_2|)\) in the downstream region \((x > 0)\) is subsonic. Therefore, the flow field around the shock front \((x = 0)\) has convergence:
\[ \text{div} \, \tilde{u} = -(u_1 - u_2) \delta(x) \]  

(12)

where \(\delta(x)\) is the Dirac’s delta function.

We have now the governing equation of the diffusive shock acceleration process in one-dimensional representation:
\[ \frac{\partial f}{\partial t} + u_i \frac{\partial f}{\partial x} = -\left\{ \frac{1}{2} (u_1 - u_2) \frac{\partial f}{\partial \ln p} \right\} \delta(x) + \frac{\partial}{\partial x} \left( D \frac{\partial f}{\partial x} \right) \]  

(13)

with \(i = 1\) and 2, corresponding to upstream and downstream regions, respectively. The acceleration effect accompanying with the flow convergence at the shock has a simple physical interpretation as described in Fig. 4: in this cartoon illustration, the upstream and downstream scatterers (namely, MHD waves) are playing in this ‘tennis’ game. While the attacking upstream players sit on the supersonic ‘rockets’ moving with the velocity \(u_1\), the defending downstream players stand on the ‘carts’ moving slowly with the subsonic velocity \(u_2\). The ‘tennis ball’, namely the particle being accelerated, bouncing between the players gets a net momentum gain \(2(u_1 - u_2)\) per one ‘rally’. After repeating this rally for many times, significant acceleration can occur.

Assuming the steady state, we get a simple analytic

\(^2\) For simplicity, they consider the case where the averaged magnetic field is parallel to the shock normal direction (parallel shock configuration). Note that the the magnetic field can be taken to be parallel to the shock normal direction both in the upstream and downstream regions in this case. (Switch-on parallel shocks, in which the upstream magnetic field is parallel to the shock normal direction but the downstream magnetic field is oblique to the shock normal direction, are limited to the low-\(B\) and low Mach number regime and have a minor astrophysical significance.) For oblique shocks, where both upstream and downstream magnetic fields make oblique angles to the shock normal direction, accelerated particles are not only being scattered by turbulence, but also make gradient-\(B\) and curvature drift motions owing to the abrupt change of the magnetic field at the shock front. These drift motions bring additional energy changes to the cosmic ray particles (mainly acceleration — so-called shock drift acceleration effect). It has been shown [22] that even with such additional energy gain, the resultant spectral index of accelerated particles, \(\Gamma\) or \(\gamma\), remains the same as in the parallel shock case.
solution of Eq. (13):

\[
f(x, p) = \begin{cases} 
  f_{-\infty}(p) + (f_+(p) - f_{-\infty}(p)) \exp \left( u_1 \int_0^x \frac{d\chi}{D(x', p)} \right) & (x < 0) \\
  f_+(p) & (x < 0) 
\end{cases}
\]

with

\[
f_+(p) = \Gamma p^{\Gamma - 1} \int_0^p f_-(p') p'^{\gamma - 1} \, dp'
\]

where \( f_{-\infty}(p) \) is the particle distribution function at \( x = -\infty \) given as the boundary condition (namely, particles externally injected into the acceleration region). If the ‘injection’ is limited to the lowest energy and given by \( f_{-\infty}(p) \propto \delta(p - p_0) \), we have the power law spectrum, \( f_+(p) \propto p^{-\Gamma} \) above \( p_0 \). For the omnidirectional differential number density of cosmic ray particles, \( N(x, p) \) \( (N(x, p) \, dp = 4\pi f(x, p)p^2 \, dp) \), we have \( N(x, p) \propto p^{-\gamma} \) with

\[
\gamma = \Gamma - 2 = \frac{u_1 + 2u_2}{u_1 - u_2} = \frac{r + 2}{r - 1}
\]

where \( r = u_1/u_2 \) is the shock compression ratio. It is remarkable that while the dependence of \( f(x, p) \) on \( x \) is determined through the detailed functional form of the diffusion coefficient \( D(x, p) \) for \( x \) and \( p \) as seen in Eq. (14), the resultant spectral index \( \Gamma \) or \( \gamma \) depends only on the shock compression ratio \( r \) but not on such details. Furthermore, in the strong shock limit (\( r \to 4 \)), we have \( \gamma \to 2 \), which is close to the observed spectral index (2.7) of cosmic ray particles between 10 GeV and 10^{15} eV (Fig. 2). If we assume that the acceleration of cosmic ray particles occurs in the strong supernova shocks, the difference of 0.7 is attributable to the escape effect from the galactic disk during their propagation. This is the accepted scenario of the supernova origin of galactic cosmic rays.

When the diffusive shock acceleration theory was established from late seventies to early eighties, there were extensive collaborations between astrophysical theorists and space experimentalists working on data from heliospheric shock environments [23–25]. Fig. 5 is from the recent observation [26,27] of the interplanetary shock which propagated from the sun to 1 AU in 32 h after its ejection from an explosive event on the sun. From this timing of shock arrival, we got the average propagation velocity\(^3\) of \( \sim 1300 \) km/s. The arrival of an interplanetary shock is seen as the jumps in \( \mathbf{B} \) and \( V_\infty \) (the top and fourth panels). In the bottom panel, the protons of the energy between several tens of keV and several MeV show the characteristic behavior in the diffusive shock acceleration process, namely, an exponential increase toward the shock front, and the flat-top shape in the downstream region. The energy spectra of accelerated protons and electrons (not shown here) were also found to be consistent with the theory [26].

4. Cosmic ray modified shocks

A recent topic relating to the shock acceleration process is its intrinsic nonlinearity: since acceleration at strong shocks is expected so efficient that energy densities of accelerated particles can become comparable or even exceed the thermal/magnetic energy densities in the shock upstream region. We expect therefore that the back-reaction of

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\(^3\) From the Rankine–Hugoniot condition, the local shock speed is estimated to be \( \sim 920 \) km/s. (The difference of 380 km/s between the average and local velocities is known to be attributable to the deceleration of the shock during its propagation to 1 AU.) Since the upstream solar wind plasma was flowing with the velocity of 550 km/s, the shock velocity relative to the upstream plasma was \( \sim 370 \) km/s, corresponding to the Alfvén Mach number \( (u_1/V_A) \) of \( \sim 5.8 \).
accelerated particles modifies the shock structure itself. This modification has been discussed theoretically in terms of ‘cosmic-ray-modified’ shocks (CRMS).

Drury and Völk [28] has treated the problem as an interaction between the plasma as a background gas (flow velocity \( u \), mass density \( \rho \), pressure \( P_g \), ratio of the specific heats \( \gamma_e \)) and a ‘cosmic ray gas’ (pressure \( P_c \), ratio of specific heats \( \gamma_c \)). It is assumed that while \( P_c \) is not negligible, the mass density of cosmic ray particles is much smaller than \( \rho \) and thus can be neglected. In a steady-state two-fluid system, behaviors of the background and cosmic ray gases are described by the equations:

\[
\rho u \frac{\partial}{\partial x} u + \frac{\partial}{\partial x} (P_c + P_g) = 0
\]

\[
u \frac{\partial}{\partial x} \rho + \rho \frac{\partial u}{\partial x} = 0
\]

\[
u \frac{\partial}{\partial x} P_g + \gamma_c P_g \frac{\partial u}{\partial x} = 0
\]

where \( \tilde{D} \) is a spatial diffusion coefficient for cosmic ray particles averaged over their momentum. (For simplicity, we neglect the effect of magnetic field in the plasma other than its role in the scattering of cosmic ray particles.)

From the above equations, we obtain three integration constants, \( A, B, C \), for which we have

\[ho u = A = \rho_1 u_1
\]

\[
Au + P_c + P_g = B = \rho_1 u_1^2 + P_{c1} + P_{g1}
\]

\[
\frac{1}{2} Au^2 + \frac{\gamma_e}{\gamma_e - 1} u P_g + \frac{\gamma_c}{\gamma_c - 1} u P_c = C + \tilde{D} \frac{\partial P_c}{\partial x}
\]

\[
(\rho_1, u_1, P_{c1}, P_{g1} \) are values of the corresponding quantities at the upstream boundary.)

Note that the ram pressure of the gas, \( \rho u^2 = Au \propto u \).

From Eqs. (18)–(20), we can obtain a Rankine–Hugoniot-like condition for a CRMS, from which it is shown [28] that a
CRMS consists of either (1) a gradual part and an abrupt (subshock) part (Fig. 6), or (2) a gradual part alone (without subshock). It is noted that the second case occurs when the partial pressure of low energy cosmic ray particles given as the upstream boundary condition exceeds some critical value.

For a while, the observational evidence of CRMS has been limited to the relatively weak effect seen at the Earth’s bow shock. Only recently a clear observational evidence of this effect is obtained at the passage of the strong interplanetary shock shown already in Fig. 5. Fig. 7 shows an enlarged time profile over 4 h around the shock arrival. About 1 h before the shock arrival, partial pressures of accelerated particles as well as thermal and magnetic pressures showed gradual increases (bottom panel). Corresponding to these pressure increases, the flow velocity of the solar wind plasma showed a gradual increase of \( \sim 50 \text{ km/s} \) ahead of the main velocity jump of \( \sim 400 \text{ km/s} \) (the panel above the bottom). Note that these velocity increases were in the observer’s rest frame. In the shock rest frame, the gradual velocity change corresponds to the gradual deceleration of the upstream plasma flow. Along with it, the ram pressure of the plasma flow (in the shock rest frame) was decreased by \( \sim 1.3 \times 10^{-10} \text{ Pa} \). On the other hand, summing up the partial pressure increases of thermal protons, magnetic field, energetic protons and electrons in the upstream region, we have gotten \( \sim 1.0 \times 10^{-10} \text{ Pa} \), which was fairly close to the ram pressure change. From this approximate pressure balance, we conclude that this interplanetary shock belonged to the category of CRMS.

almost all astrophysically significant shocks are expected to become CRMS, in situ observational evidence as stated in the above is indispensable for the detailed study.

For supernova shocks, Ellison et al. [29] have developed an extensive theoretical model. Interestingly, a recent observation shows that the measured electron temperature in the downstream region of a shock of a supernova 1E0102.2-7219 is significantly lower than the expectation based on the classical MHD shock theory [30]. This temperature discrepancy could be reconciled, as these authors of the report have argued, assuming that a significant fraction of the shock energy has gone into generating cosmic rays rather than contributing to the heating of the postshock electrons and ions. Such a situation is expected if this shock is a CRMS in which the subshock disappears. Further theoretical and

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4. Here, the ‘cosmic rays’ are suprathermal ions and electrons (several tens keV \( \sim \) several tens MeV), accelerated by this interplanetary shock.

5. There were several unobserved components: we did not measure the main thermal component of the solar wind electrons owing to the limited energy coverage of the plasma detector. Further, contribution of energetic electrons above 40 keV was not measured. Considering the uncertainty coming from the above limitations, we accept remaining \( \sim 30\% \) difference in the pressure balance as an allowable error.
observational works seem to be needed to get definite conclusions about the nature of this interesting supernova shock.

5. Neutral beam injection

The next topic is a ‘neutral beam injection’ into the heliosphere. It is now known that the local interstellar matter (LISM) has a relative velocity of ~25 km/s with respect to our heliosphere (Fig. 8). Although the ionized components of the LISM are deflected at the heliospheric bow shock as well as at the heliopause and cannot freely penetrate into the heliosphere, neutral components such as hydrogen and helium atoms can cross these borders. While the solar wind protons and UV photons prevent H atoms from coming into the region nearer than several AU from the sun, He atoms, which are affected only weakly by these effects, can come into the region well inside the earth’s orbit at 1 AU, where they are finally to be ionized\(^6\) and picked up by the solar wind flow with the rate \(\sim 10^{-7} \text{ s}^{-1}\) or \(~1\%\)/day. Since these He\(^+\) ions have much higher rigidity than usual solar wind ions, they can be easily accelerated at various shock environments. People have believed that this explains the origin of the anomalous cosmic ray (ACR) component in the energy range of several MeV/n to several tens of MeV/n.

In Fig. 9, which covers the energy range between 10 and 1000 MeV/n, namely, the lowest energy part of Fig. 2 and further below, the observation of ACRs is shown [31]: upper curves are for protons, and lower curves for helium nuclei. Solid curves with label A and B show the expectations from a simple modulation model for protons (A) and helium nuclei (B), respectively. Dashed curves are from observations in 1965 for both protons and helium nuclei.

It is seen in Fig. 9 that in the energy range below \(~30 \text{ MeV/n}\), the helium flux (filled circles) was anomalously high and exceeded that of protons. Similar anomalies were also found for nitrogen and oxygen fluxes [32]. It was soon after these discoveries of the ACR components that Fisk et al. [33] proposed that the pickup interstellar ions as He\(^+\), N\(^+\) and O\(^+\), can be accelerated up to the ACR energy at the terminating shock of the solar wind, which is expected to be formed at \(~70–100 \text{ AU}\) from the sun. Since then their proposal has been confirmed in two ways: (1) The measured charge state of the oxygen ACRs is mainly +1, as consistent with the prediction.\(^7\) (2) Pickup He\(^+\) ions directly after their ionization have been identified by the plasma detectors aboard of solar-wind cruising spacecraft [35–38]. The final proof of the acceleration process should come from the in situ observation of the terminating shock. Since the Voyager spacecraft will reach the terminating shock in the early part of the 21st century, people are now awaiting these new and exciting observations. It is interesting to note [39] that the physical condition at the terminating shock would be similar to that of the CRMs, which is discussed in Section 4: outside 50 AU, the partial pressure of pickup protons\(^8\) is expected to become greater than both the magnetic pressure (by a factor of >10) and the thermal

\(^6\) More than \(~90\%\) of the ionization is by the solar UV photons, and remaining \(< ~10\%\) is accomplished by the impact of suprathermal electrons in the solar wind.

\(^7\) Quite recently [34], it is further found that while the charge status of oxygen ACRs is \(+1\) below \(~20 \text{ MeV/n}\), it changes to \(+2\) (or more) above this energy. Since the ‘lifetime’ of O\(^-\) for further ionization to O\(^-2\) is several years within the inner heliosphere, this observation is interpreted as indicating the age of these oxygen ACRs after their pickup epoch as well as the acceleration rate (\(~10 \text{ MeV year}^{-1}\)).

\(^8\) Interstellar neutral hydrogen atoms cannot penetrate into 1 AU owing to the solar radiation pressure. However, they become the dominant species of pickup ions beyond several AU.
pressure of solar wind protons (by a factor of \(>100\)). Modification of the shock structure owing to the existence of these nonthermal pickup protons relates self-consistently to the acceleration efficiency for themselves.

6. Highest energy of particles accelerated in the heliosphere

In the literatures discussing the origin of UHECRs, we often see the argument of ‘Hillas limit’. This is what Hillas proposed in Ref. [10] giving a theoretical limit of the energy of particles attainable in acceleration processes: In the acceleration region having characteristic velocity \(\beta = u/c\) (normalized by the light speed \(c\)), the magnitude of the magnetic field \(B\), and the characteristic spatial scale \(L\), particles of the charge \(Ze\) can be accelerated up to the maximum energy

\[
E_{\text{max}} = Ze \beta BL
\]  

(21)

If the characteristic velocity \(u\) is the velocity of the background plasma flow, the convection electric field \(E\) is given by \(\beta B\) so that the Eq. (21) is equivalent to the simple relation, \(E_{\text{max}} = Ze EL\). It is of interest to see whether this simple assertion gets support or not from the heliospheric observations.

Note that the Hillas limit for the diffusive shock acceleration process can be obtained from the following discussion. The solution by Blandford and Ostriker in Eq. (13) gives the characteristic scale length \(\lambda\) in the upstream region as

\[
\lambda = \frac{D_1}{u_1}
\]  

(22)

where \(D_1\) and \(u_1\) are the upstream diffusion coefficient and flow speed, respectively. Let us define \(\eta = D_1/D_B\), where \(D_B\) is the Bohm diffusion coefficient \((1/3)\rho_e w_c\) (\(\rho_e\): cyclotron radius, \(w_c\): particle speed). In the strong scattering limit, we expect that \(\eta \to 1\), or equivalently \(D_1 \to D_B\). If the scattering is quite weak, on the other hand, \(\eta \gg 1\) or \(D_1 \gg D_B\). Using \(\eta\), we can write

\[
D_1 = \frac{2\eta E_c}{3Ze B} \cdots \text{NR} \quad \text{or} \quad D_1 = \frac{\eta E_c}{3Ze B} \cdots \text{R} \quad (23)
\]

where cases with NR and R correspond to nonrelativistic \((u/c = \beta \ll 1)\) and ultra-relativistic \((\beta \to 1)\) cases, respectively. Note that we have utilized the relation that \(\rho_e w_c = 2E_c/(Ze B_NR)\) or \(E_c/(Ze B\ R)\), with the particle energy \(E\). For a given \(\lambda\), Eqs. (22) and (23) lead to the energy-scale relation:

\[
E = \frac{3}{2\eta} Ze B\lambda\beta \cdots \text{NR} \quad \text{or} \quad E = \frac{3}{\eta} Ze B\lambda\beta \cdots \text{R} \quad (24)
\]

For the efficient acceleration to occur, the upstream region of the size \(\lambda\) should be embedded\(^{10}\) in the acceleration region of the total size \(L\), so that we require \(\lambda < L\). Therefore, from Eq. (24), we have

\[
E < E_{\text{max}} = \frac{3}{2\eta} Ze B\lambda\beta \cdots \text{NR} \quad \text{or} \quad E < E_{\text{max}} = \frac{3}{\eta} Ze B\lambda\beta \cdots \text{R} \quad (25)
\]

In the limit of \(\eta \to 1\) (Eq. 25) coincides with the original Hillas limit within a numerical factor 3/2 or 3 for the cases of NR or R, respectively. (In the above, we have modified the Drury’s argument from Ref. [40]. His definition of \(L\) slightly differs from ours.)

For various acceleration phenomena in the heliosphere, the corresponding Hillas limits are estimated in Table 1. Interestingly, for many cases, these simple estimations coincided with the observed maximum energies in the order of magnitude. Note, however, that the observed maximum energies at heliospheric shocks might exceed the ‘theoretical Hillas limit’ since the estimation of \(\eta\) based on the in situ observations of the MHD turbulence usually gives \(\eta\) of the order of 10–100 (i.e. moderately strong turbulence), while the agreement between the Hillas limit and observed maximum energies in Table 1 requires \(\eta \to 1\) (i.e. extremely strong turbulence). To see whether this ‘discrepancy’ is serious or not, we need further observational and theoretical studies.

7. Concluding remarks

We have briefly given an overview for selected topics

\(^{9}\) The ‘astrophysical convention’ of the Bohm diffusion coefficient differs from that of the fusion community by a numerical factor, 16/3.

\(^{10}\) Otherwise, the particle losses due to the escape beyond upstream and downstream boundaries are important, and the acceleration efficiency is reduced significantly.
Table 1
Theoretical Hillas vs. observations

| $B_1 = B/1G$ | $L_{10} = L/10^{10}$ cm | $\beta = v/c$ | $E_{\text{max}}$ (Hillas limit) $\simeq 3\beta B_1 L_{10} \text{ TeV}$ | $E_{\text{max,obs}}$ | Remarks |
|--------------|------------------------|---------------|---------------------------------|---------------------|---------|
| Solar flare Van Allen Belt Earth CRAND* + stochastic betatron acceleration | $10^2$ | 0.001–0.01 | 300 GeV–3 TeV | < ~30 GeV* | CME shocks |
| Jupiter Heliosphere | 0.1 | 0.1 (1.4$R_\oplus$) | 0.03$^b$ | 1 GeV | 0.6 GeV |
| Bow shock + self-generated MHD turbulence | 1 | 1 (1.4$R_\oplus$) | 0.03$^b$ | 100 GeV | ? |
| Bow shock + IMF kink | $6 \times 10^{-5}$ @ 1 AU | $10^3$ (1 AU) | 0.001–0.002 | 0.2–0.4 GeV | ESP < ~0.1 GeV |
| | $6 \times 10^{-5}$ | 3 (50$R_\oplus$) | 0.001–0.002 | 0.5–1 MeV | ACR < ~0.1 GeV |
| | $6 \times 10^{-5}$ | 3 (50$R_\oplus$) | 0.001–0.002 | 0.5–1 MeV | $p \sim 0.1–0.2$ MeV |
| | | | | | CME shocks |
| | | | | | Terminating shock |
| | | | | | Shock |
| Magnetotail | $10^{-4}$ | 1 (15$R_\oplus$) | 0.003–0.005 | 0.9–1.5 MeV | $e, p \sim 1$ MeV |
| | | | | | Multiple but a few reflections at the shock$^d$ |
| | | | | | Reconnection |

* Ref. [41].
$^b$ For the case of Van Allen belt, $\beta$ is not the velocity of the background plasma, but a coefficient (~0.03) giving the stably trapping condition for particles in the planetary dipole field.
$^*$ CRAND is cosmic ray albedo neutron decay. Some of the neutrons, which are produced through the nuclear interactions between cosmic ray particles and atmospheric nuclei, escape upward from the atmosphere and beta decay in the region of the Van Allen belt. Part of the decay products, protons, are then trapped in the planetary dipole field. It is known that the observed energy spectrum of these CRAND protons can be explicable only if we take into account of the stochastic acceleration by turbulent electromagnetic fields after their trapping [42].
$^d$ Ref. [43].

from recent progresses in the cosmic plasma research. Of course, there is my personal bias in the selection of the topics. I am afraid that some of important topics cannot be included because they are beyond my personal capability. Anyway, what I want to say is that we are facing at a vast variety of interesting (sometimes quite puzzling) plasma processes in the universe. Importance of interdisciplinary collaborations cannot be overemphasized.

Appendix A. Quasi-linear theory

Let us derive the quasi-linear equation for the pitch-angle scattering process from the Vlasov equation:

$$\frac{\partial f}{\partial t} + \frac{\partial f}{\partial \vartheta} + \frac{q}{M} \left[ \vec{E} + \frac{1}{c} \vec{w} \times \vec{B} \right] \frac{\partial f}{\partial \vec{w}} = 0 \quad (A1)$$

We assume that the electromagnetic field can be decomposed into the averages and perturbations:

$$\vec{E} = \vec{E}_0 + \vec{E}_1$$

$$\vec{B} = \vec{B}_0 + \vec{B}_1$$

where $\vec{E}_0$ is the frozen-in electric field, $-\vec{w} \times \vec{B}/c$ (\(\vec{w}\) : the average bulk velocity of the plasma), and can be eliminated from the formalism by taking a suitable coordinate transformation. The following discussion is ordered by the smallness parameter $\epsilon = |B_1/B_0|$. Let $f_0$ and $f_1$ be the zeroth and first order components of $f$. Retaining the terms up to the first order of $\epsilon$, we have

$$\frac{Df_1}{Dt} = \frac{\partial f_1}{\partial t} + \frac{\partial f_1}{\partial \vartheta} + \frac{q}{M} \left[ \vec{E}_1 + \frac{1}{c} \vec{w} \times \vec{B}_1 \right] \frac{\partial f_1}{\partial \vec{w}} = 0$$

where the symbol $\langle \rangle$ means the correlation. We limit the discussion to the case of MHD waves ($\vec{B}_1, \vec{E}_1$) propagating parallel to the average magnetic field direction ($\vec{x}$) of $\vec{B}_0$, which can be decomposed into circularly polarized Fourier modes:

$$B_{ly} = \frac{1}{2} \sum_{m,k} \left[ B_{mk} e^{ikx - \omega t} + c.c. \right]$$

$$B_{lz} = \frac{1}{2} \sum_{m,k} \left[ -i B_{mk} e^{ikx - \omega t} + c.c. \right]$$

$$E_{ly} = \frac{1}{2} \sum_{m,k} \left( \frac{\omega}{kc} \right) \left[ -i B_{mk} e^{ikx - \omega t} + c.c. \right]$$
\[ E_{1z} = - \frac{1}{2} \sum_{\alpha, k} \left( \frac{\omega}{k} \right) B_{\alpha k} e^{i(kx-wt)} + \text{c.c.} \]  

(A5)

where \( B_{\alpha k} \) are complex Fourier amplitudes for \((\omega, k) \) \((\omega = \pm kV_A)\). We also represent the \(yz\) components of particle velocity \( \vec{w} = (w_x, w_y, w_z) \) as \( w_x = w_{\perp} \cos(\Omega t - \alpha) \)

(A6)

where \( \alpha \) is the cyclotron phase angle, and \( \Omega \) the cyclotron frequency, \( qB_{\perp}/Mc \) (\( \Omega \) becomes positive or negative according to the sign of the charge \( q \)). We also define \( w = \sqrt{w_x^2 + w_y^2 + w_z^2} \) and \( \mu = \frac{w_x}{w} \). After some manipulation, we have

\[
\frac{DF_{1,\text{m.k}}}{Dt} = -\frac{1}{2} \left\{ \frac{qB_{\alpha k}}{M} \frac{\omega}{k} \left( \hat{G}_{\alpha k, f_0} e^{i(\omega - kw_x - \Omega \tau + i \alpha)} + \text{c.c.} \right) \right\}
\]

(7)

where the operator \( \hat{G}_{\alpha k} \) is given by

\[
\hat{G}_{\alpha k} = \frac{w_{\perp}}{w} \left\{ \frac{\partial}{\partial w} + \left[ \frac{k}{\omega} - \frac{\mu}{w} \right] \frac{\partial}{\partial \mu} \right\}
\]

(A8)

By integrating Eq. (A7) along the particle trajectory, we have

\[
f_{1,\text{m.k}} = -\frac{1}{2} \int_{-\infty}^{\infty} \times dt \left\{ \frac{qB_{\alpha k}}{M} \frac{\omega}{k} \left( \hat{G}_{\alpha k, f_0} e^{i(\omega - kw_x - \Omega \tau + i \alpha)} + \text{c.c.} \right) \right\}
\]

\[
= \frac{1}{2} \left\{ \frac{qB_{\alpha k}}{M} \frac{\omega}{k} \left[ \frac{1}{\omega - kw_x - \Omega} \hat{G}_{\alpha k, f_0} e^{i(\omega - kw_x - \Omega \tau + i \alpha)} + \text{c.c.} \right] \right\}
\]

(A9)

Substituting Eq. (A9) into Eq. (A4), we have

\[
\frac{DF_{f_0}}{Dt} = \sum_{\alpha, k} \left\{ \frac{qB_{\alpha k}}{M} \frac{\omega}{k} \left[ e^{-i(\omega - kw_x - \Omega \tau + i \alpha)} \right] \right\}
\]

\[
\times \left[ 1 - \frac{k}{\omega} \mu w \right] \frac{\partial}{\partial \alpha} + \text{c.c.} \right\} f_1
\]

\[
= \sum_{\alpha, k} \frac{q^2}{2M^2c^2} \left( \frac{\omega}{k} \right)^2 B_{\alpha k} B_{\alpha k} \left\{ -i \hat{G}_{\alpha k} + \frac{1}{w_{\perp}} \right\}
\]

\[
\times \left[ 1 - \frac{k}{\omega} \mu w \right]\left[ \frac{1}{\omega - kw_x - \Omega} \hat{G}_{\alpha k, f_0} + \text{c.c.} \right] \right\}
\]

\[
= \frac{\pi}{2} \frac{q^2}{M^2c^2} \sum_{\alpha, k} B_{\alpha k}^2 \left( \frac{\omega}{k} \right)^2 \frac{1}{w_{\perp}} \hat{G}_{\alpha k} \left[ w_{\perp} \delta(\omega - kw_x - \Omega) \right]
\]

(A10)

where the combination of operators, \( \hat{G}_{\alpha k} \cdots \hat{G}_{\alpha k} \) includes the second order derivatives (with respect to \( w \) and \( \mu \)) representing phase space diffusion. For particles with sufficiently high velocity (\( w \gg V_A \)), the pitch angle diffusion term (\( \propto \nabla^2/\mu^2 \)) dominates over the energy diffusion term (\( \propto \nabla^2/\omega^2 \)) and the cross-terms (\( \propto \nabla^2/\omega \partial \omega / \partial \mu \)).

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