Effects of Anomalous Cosmic Rays on the Structure of the Outer Heliosphere

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Abstract

Based on Voyager 1 observations, some anomalous cosmic rays (ACRs) may have crossed the heliopause and escaped into the interstellar medium, providing a mechanism of energy transfer between the inner and outer heliosheaths that is not included in conventional magnetohydrodynamics (MHD) models. In this paper, we study the effect of energetic particles’ escape through the heliopause on the size and shape of the heliosphere using a simple model that includes diffusive transport of cosmic rays. We show that the presence of ACRs significantly changes the heliosphere structure, including the location of the heliopause and termination shock. It was found that the heliopause would contract for certain values of the ACR diffusion coefficients when the diffusive particles’ pressure is comparable to the pressure of the plasma background. The difference in Voyager 1 and 2 observations of energetic particles during their respective termination shock crossings is interpreted here as due to the differences in diffusion environments during the different phases of the solar cycle. The shorter period of enhanced ACR intensities upstream of the shock measured by Voyager 2 may have been caused by weaker radial diffusive transport compared with the time of Voyager 1 crossing. We conclude that ACR diffusive effects could be prominent and should be included in MHD models of the heliosphere.

Key words: cosmic rays – diffusion – magnetohydrodynamics (MHD) – methods: numerical – plasmas – Sun: heliosphere

1. Introduction

The heliopause (HP) is the boundary of the heliosphere, separating the solar wind from the moving local interstellar medium (LISM). It was reported that Voyager 1 (V1) had crossed the boundary in 2012 August (Webber & McDonald 2013) and entered interstellar space. Several remarkable changes occurred in charged particle intensities upon the HP crossing. The suprathermal heliospheric ions, including the anomalous cosmic rays (ACRs), were observed past the HP, but disappeared completely over a period of less than a month (Stone et al. 2013; Florinski et al. 2015), indicating rapid escape along the nearly laminar interstellar magnetic field lines (Burlaga et al. 2015). On the contrary, the intensity of galactic cosmic rays (GCRs) increased significantly at the heliopause and remained at a steady (unmodulated) level in the ISM (Cummings et al. 2016), a scenario consistent with results from global modeling (Jokipii 2001; Guo & Florinski 2014; Kota & Jokipii 2014). Because both kinds of cosmic rays behave very differently from the bulk plasma across the HP, they ought not to be treated as a single plasma component from a macroscopic point of view. The purpose of this paper is to evaluate the effect of separate treatment of the diffusive ion species on the structure of the heliosphere.

Generally, the HP is thought to be a tangential discontinuity where the combined plasma and magnetic field pressures on the two sides are balanced. This can be written as \( P_{\text{HIS}} = P_{\text{OHS}} \), where \( P = \{ P_B + P_T + P_{\text{GCR}} + P_{\text{ACR}} \} \), \( P_B \) being the magnetic field pressure, \( P_T \) the background (“thermal”) plasma pressure, and \( P_{\text{GCR}} \) and \( P_{\text{ACR}} \) are the contributions from the respective cosmic-ray species. As mentioned above, the measured intensity of low-energy energetic charged particles dropped to background levels when V1 entered the interstellar region, indicating that \( P_{\text{ACR}} \) had a sharp decrease. Although the HP may be rippled by waves and turbulences caused by potential instabilities (Florinski et al. 2005; Borovikov et al. 2008; Florinski 2015) and magnetic reconnection events (Opher et al. 2011; Swisdak et al. 2013), we treat it as an ideal discontinuity. Consequently, the heliosphere will react to compensate the pressure loss of the ACRs near the HP.

The effects of cosmic rays on the heliosphere have been considered previously. In a two-component model (background “thermal” plasma and GCRs), Myasnikov et al. (2000a) found that GCRs could significantly modify the shape and structure of the termination shock (TS) and the bow shock (if it exists) and change the width of the heliosheaths. However, the more sophisticated three-component model (background, H atoms, and GCRs) by the same authors (Myasnikov et al. 2000b) showed a different outcome, namely, that the GCR influence on the plasma flow is almost negligible compared with the influence of H atoms, the only exception being the location of the bow shock. The pressure of ACRs leads to a formation of a smooth precursor upstream of the TS and a change in the radial distance to the subshock (Fahr et al. 2000; Alexashov et al. 2004; Florinski et al. 2004). The recent HP crossing by V1 provided observational evidence that ACR intensities decreased rapidly beyond the boundary. This transfer of energy from the inner to the outer heliosheath, not considered previously in computational work, should affect the pressure balance at the HP. Therefore, it is necessary to reevaluate the effect of ACRs on the outer heliosphere using a more accurate model.

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In this paper, we first present observational results based on \textit{VII} data and then perform numerical simulations to investigate the effects of ACRs on the heliosphere using an MHD–neutral–cosmic-ray coupled model, where the cosmic rays are diffusive and are governed by a transport equation. The main difference between the model of Alexashov et al. (2004) and this model is the inclusion of both the interplanetary and interstellar magnetic fields and the decoupling of ACRs from the background plasma in the outer heliosheath (see Section 3). We also compare model-derived geometries of the TS and the HP and the thickness of the heliosheath with what could be deduced from \textit{Voyager} observations.

2. Observations

First, we analyze the spectra of the energetic charged particles measured by \textit{VII} before and after its encounter with the HP, years 2012.4 and 2012.9, respectively (Figure 1). Similar spectra were originally reported by the \textit{Voyager} observational teams (Krimigis et al. 2013; Webber & McDonald 2013). The data used herein are from the Low Energy Charged Particle (LECP) instrument that measures differential fluxes of ions with energies 0.04–40 MeV/nucleon, and proton fluxes in three energy channels, 0.57–1.78 MeV, 3.4–17.6 MeV, and 22–31 MeV. The data from Cosmic Ray Subsystem (CRS) were used to plot the spectra of lower-energy GCRs for comparison.

The “before” and “after” spectra are shown in Figure 1 with solid and dashed lines, respectively. The former is a power law with a slope close to $-1.5$ at the lowest energies, followed by a rollover at higher energy (Krimigis et al. 2013). The “after” spectra are remarkably flat over the lower energy range, indicating that the detector has entered the so-called “depletion region” where the ACRs escape and become undetectable at all energies. It was reported that the magnitude of the magnetic field increased from $\sim 0.24$ to 0.43 nT to maintain the pressure-balanced structure, which is assumed to be the heliopause (Burlaga et al. 2013). Keep in mind that the ACR particles diffuse across the boundary, and their pressure is the same on each side of the heliopause. This rapid decrease of heliospheric particles beyond the HP indicates that they travel essentially scatter-free in the interstellar region; therefore, the pressure balance condition should exclude the contribution of ACRs (or at least the component of the pressure tensor parallel to the magnetic field). Note that in this work the term ACRs encompasses all heliospheric energetic ions with energies down to a few KeV.

Figure 2 shows the observed partial pressures of energetic particles in different energy ranges along the \textit{VII} direction near the HP. Because the low-energy particle intensities from the LECP instruments are below the background after the HP crossing, the corresponding pressure is assumed to be zero. Only CRS data for 133–242 MeV protons are plotted beyond HP, showing that the GCR pressure does not change too much after the crossing compared to those of ACRs because of their lower intensities. Here, we assume that the increase in GCR pressure across the HP can be ignored in comparison with the contribution from the particles with lower energies. From the figure, the partial pressure of ions with energies above 5 KeV is estimated to be $\sim 0.5$ pdyne cm$^{-2}$ before HP crossing. Because \textit{Voyagers} do not measure particle energies down to 5 KeV, here we assume that the intensities of lower energetic particles obey the same power law as the particles at higher energies. The pressure of the proton plasma in the inner heliosheath is 1.9–2.1 pdyne cm$^{-2}$ based on the \textit{IBEX} ribbon observations (Schwadron et al. 2011; Livadiotis et al. 2013), which allows us to make a rough estimate that the upper limit on the fraction of the pressure lost across the HP is $\sim 25\%$ of the total plasma pressure. This fraction will drop to 10% if we consider only energetic particles within the energy range 40–4000 KeV, which contribute about 0.2 pdyne cm$^{-2}$, as can be seen from Figure 2.

If one assumes that all energetic particles are diffusive and are capable of crossing into the outer heliosheath, then it would be expected that their escape will result in a change of momentum and energy density in the inner heliosheath. In the next section, we perform global numerical simulations of the ACR mediated heliospheric interface to quantify this effect.

3. Numerical Modeling

We use a two-fluid hydrodynamic approach to model the coupled system of plasma and cosmic rays, in which the cosmic rays are treated as a single fluid with negligible mass. This is a common approach when the details of the ACR energy distribution are not of interest (e.g., Myasnikov et al. 2000a;
Fahr et al. 2000). The equations to be solved can be written as

\[
\frac{\partial U}{\partial t} + \nabla \cdot F(U) = S + Q,
\]

(1)

where the nine-component vectors \( U, F(U), \) and \( S \) are

\[
U = \begin{pmatrix} \rho \\ \rho \mathbf{u} \\ B \\ E \\ \rho P_e \end{pmatrix}, \quad F(U) = \begin{pmatrix} \rho \mathbf{u} \\ \rho \mathbf{u} u_p + P_T I - BB \\ uB - Bu \\ (E + P_T)u - (u \cdot B)B \\ \rho P_e u \end{pmatrix}, \quad S = \begin{pmatrix} 0 \\ -\nabla \cdot (P_e I) \\ 0 \\ -u \cdot \nabla P_e + \alpha \nabla \cdot u \\ \rho \{\nabla \cdot (\kappa \nabla P_e) + [(1 - \gamma_e)\alpha - \gamma_e P_e] \nabla \cdot u\} \end{pmatrix}.
\]

(2)

(3)

Here, \( \rho, \mathbf{u}, \) and \( B \) represent the plasma density, velocity, and magnetic field, respectively. The energy density is \( E = P_g/(\gamma - 1) + \rho u^2/2 + B^2/2 \) and the total pressure \( P_T = P_g + B^2/2 \), where \( P_g \) is the pressure of the thermal plasma. \( P_e \) denotes the “cosmic-ray” (that includes all heliospheric energetic particles) pressure, and \( \kappa \) is the isotropic energy-averaged diffusion coefficient. This is clearly an oversimplification of the problem driven by the numerical constraints on the time step. The general diffusion tensor is \( \kappa_{ij} = \kappa_{ii}\delta_{ij} + (\kappa_i - \kappa_{ii})B_i B_j/B^2 \), assuming axial symmetry about the direction of \( B \). The isotropic approximation for the diffusion tensor is clearly invalid in the LISM because the nearly laminar interstellar magnetic field ensures that \( \kappa_i \gg \kappa_{ii} \); the problem is circumvented by switching off the coupling between plasma and ACRs in that region (see Equation (6)). \( \gamma \) and \( \gamma_e \) are the adiabatic indices of the thermal plasma and ACRs, both of which are assumed to be 5/3. The source term \( Q \) describes charge exchange between the plasma and neutral atoms (Pauls et al. 1995). This model features four neutral populations that originate from the undisturbed interstellar medium, the outer heliosheath, the inner heliosheath, and the supersonic solar wind, respectively. To model the neutrals, the values of \( B \) and \( P_e \) are set to zero in the above equations. The quantity \( \alpha \) is the local injection rate of ACR particles from pickup ions (PUIs), which is simply assumed to be proportional to the thermal pressure, \( \alpha = \alpha'P \), where \( \alpha' \) is a free parameter measuring the injection efficiency (Zank et al. 1993; Fahr et al. 2000; Alexashov et al. 2004) from thermal plasma to ACRs by means of the shock acceleration mechanism near the TS. There are six fluids in the coupled system in total (plasma, cosmic rays, and four neutral populations). The system was solved numerically on a level 5 three-dimensional spherical geodesic mesh (Florinski et al. 2013) with 2562 hexagonal cells on the sphere and 432 concentric layers in the radial direction. The radial resolution of the mesh near TS and HP was about 0.7 and 1.0 au, respectively.

The inner boundary was a 10 au sphere centering at the Sun, and the outer boundary was set at 800 au. The solar wind conditions were calculated from the Parker solution (Parker 1958) based on typical observations at 1 au during a solar minimum (Ebert et al. 2009). The plasma parameters at 1 au were as follows: number density of protons \( n_p = 5.5 \text{ cm}^{-3} \), flow speed \( u_p = 388 \text{ km s}^{-1} \), temperature \( T_p = 64,000 \text{ K} \), and radial component of interplanetary magnetic field \( B_r = 2.45 \text{ nT} \). For the unperturbed LISM, we used the proton number density \( n_{\text{LISM}} = 0.05 \text{ cm}^{-3} \), neutral hydrogen density \( n_{\text{LISM},H} = 0.195 \text{ cm}^{-3} \), temperature \( T_{\text{LISM}} = 6200 \text{ K} \), and magnetic field \( |B_{\text{LISM}}| = 0.3 \text{ nT} \). The LISM plasma flows with speed \( |V_{\text{LISM}}| = 23.2 \text{ km s}^{-1} \), and its direction is given by the neutral He flow with \( (\lambda, \beta)_{\text{He}} = (79^\circ, -4^\circ) \), but \( \lambda \) is the ecliptic longitude and \( \beta \) is the ecliptic latitude. The LISM hydrogen flow is deflected from its original direction (which is same as the He flow) by the interaction with the ionized component near the heliosphere and is observed at \( (\lambda, \beta)_{\text{H}} = (72^\circ, 39^\circ) \) (Funsten et al. 2013), and its projection on the HDP. The cosmic-ray pressure \( P_c \) includes only ACRs; a zero gradient boundary condition is used at the inner boundary and their pressure is set to zero at the outer boundary.

The cosmic rays diffuse very fast along the magnetic field lines in the interstellar medium because of the nearly laminar magnetic field on short spatial scales (below a few au; Burlaga et al. 2013, 2015). Here, we simply follow the empirical expression of diffusion coefficient that has been previously used in stochastic models of cosmic-ray transport in the outer heliosphere (Strauss et al. 2013; Guo & Florinski 2014):

\[
\kappa_{ii} = \lambda_0 \frac{w_B B_0}{3 |B|} \left( \frac{p}{1 \text{ GeV c}^{-1}} \right) \delta.
\]

(4)

where \( \lambda_0 = 0.0375 \text{ au} \), \( B_0 = 5 \text{ nT} \) in the heliosphere and \( \lambda_0 = 1.0 \times 10^5 \text{ au} \), \( B_0 = |B| \) beyond the heliopause. \( \delta = 1 \) for \( p > 1 \text{ GeV c}^{-1} \) and \( \delta = 1/3 \) otherwise. The values of \( \kappa_{ii} \) for energetic particles with mean energy 707 KeV n^{-1} is plotted as the dark green curve in Figure 3, featuring an expected sudden jump in the diffusion coefficient across the HP. Here, the mean energy is calculated as \( \sqrt{T_L T_H} \), where \( T_L = 5 \text{ KeV} \) and...
\(T_H = 100\,\text{MeV}\) are the low and high limits of the particle energies. In reality, the effective diffusion coefficient \(\kappa\) is a ratio of the integrated diffusive term from the Parker transport equation to the diffusive term in the fluid model (Drury & Voelk 1981), for example,

\[
\kappa_{\text{ct}} = \int_0^\infty p^3 v \kappa_{\text{ct}}(p) \frac{\partial f}{\partial x} dp \left( \int_0^\infty p^3 v \frac{\partial f}{\partial x} dp \right)^{-1},
\]

where \(p\) is the particle’s momentum, and \(f(p)\) and \(\kappa(p)\) are the distribution function and the diffusion coefficient in the Parker model. It shows that \(\kappa\) depends on the gradient of \(f\) (or \(P_k\)). For simplicity we choose three constant energy-averaged diffusion coefficient values for the ACRs: \(\kappa_1 = 2.5 \times 10^{20} \, \text{cm}^2 \, \text{s}^{-1}\), \(\kappa_2 = 2.5 \times 10^{21} \, \text{cm}^2 \, \text{s}^{-1}\), and \(\kappa_3 = 2.5 \times 10^{22} \, \text{cm}^2 \, \text{s}^{-1}\) that are plotted in Figure 3. These values approximate the mean diffusion coefficient of ACR particles with energies above 5 \(\text{KeV} \, \text{n}^{-1}\) in the heliosphere. The parallel diffusion coefficient in the LISM is extremely large, where as the perpendicular diffusion coefficient is expected to be small because of weak turbulence. Since the magnetic field is tangential to the heliopause, the perpendicular diffusion coefficient would appear to be more relevant to the problem. However, particle escape from the interaction region is due to streaming along \(B\), which is a very fast process that cannot be treated properly in the present numerical framework. To overcome this difficulty, the coupling between plasma and ACRs in the interstellar region was switched off \((S = 0)\), and the diffusion was set to zero, so that the ACR pressure equation read

\[
\frac{\partial (P_k)}{\partial t} + \nabla \cdot \left( P_k \mathbf{u} \right) = 0.
\]

This simplification implies that the ACR pressure is passively convected outside of the heliosphere. Because \(P_k = 0\) at the outer boundary and the plasma flows toward the HP, the ACR pressure will be vanishingly small in the LISM in this approximation, so that the ACR diffusion at HP from the LISM to inner heliosheath is expected to be negligible. If diffusion is re-enabled beyond the HP, the ACRs will diffuse far into the interstellar medium against the weak flow of the LISM plasma, and the pressure \(P_k\) will decrease gradually with increased radial distance. As a result, the effect of \(\nabla P_k\) on the HP dynamics will be less pronounced than the case of turning off source term \(S\), as we will show below.

4. Model Results and Discussion

4.1. Precursor of Termination Shock

In order to compare the effects of ACRs on the structure of the heliosphere, we performed a control simulation without the ACRs and three further simulations with different diffusion coefficients \((\kappa_{1-3})\). The injection efficiency \(\alpha\) was set to 0.1 initially. Figure 4 shows the results of the radial speed of the plasma flow along the \(V_1\) direction. In the figure, the four cases are labeled with different colors as shown. The vertical dotted lines indicate the positions of the HP and the TS. For simplicity, here we only plot the TS position for the case of no ACRs. As the diffusion coefficient increases from \(\kappa_1\) to \(\kappa_3\), the TS moves outward from roughly 70 to 75 au along the \(V_1\) direction. The ACRs couple with the plasma strongly for the case of \(\kappa_1\) and \(\kappa_2\), so a diffusive precursor of the TS forms upstream followed by a plasma subshock. This structure is pronounced for the case of \(\kappa_2\), with a scale length of \(\sim 10\) au. For the case with weaker diffusion, \(\kappa_1\), the precursor size decreases to a much shorter scale, \(\sim 1-2\) au as shown in the figure. As for the higher diffusion \(\kappa_3\), the precursor is nearly invisible, indicating the diffusion is too fast and the ACRs are weakly coupled to the TS background.

Theoretically, the precursor is caused by the upstream scattering of ACRs that originated near the TS or in the inner heliosheath, leading to a slowdown of the plasma flow. For example, \(V_2\) had observed a clearly identifiable TS precursor with a length of \(\sim 0.7\) au before the TS crossing in 2007 (Richardson et al. 2008). Because the plasma instrument is not functioning on board Voyager 1, it is not known whether \(V_1\) had experienced a similar solar wind slowdown or not during its shock encounter. Here, we assume that \(V_1\) had a precursor length \(a\), which is larger than that of \(V_2\), but shorter than the size of energetic ion precursor that lasted from 2002.58 to 2004.96 at \(V_1\) (Decker et al. 2005), which translates to a width of 9.7 au, assuming a stationary shock (it should be pointed out that some models featured inward motion of the shock during that period; see, e.g., Izmodenov et al. 2008; Washimi et al. 2012). Here, we assume that \(0.7\) au \(< a < 9.7\) au. Another difference between the two probes is the distances of TS crossing, 94 au for \(V_1\) (Decker et al. 2005) and 85 au for \(V_2\) (Stone et al. 2008). Based on the obtained simulation results and assuming that TS is nearly symmetric between \(V_1\) and \(V_2\) for the similar interplanetary conditions, one may argue that the \(V_1\) crossing could be described as the case of \(\kappa_2\) and the \(V_2\) crossing as the case of \(\kappa_1\).

From the right-hand side of Equations (1)–(3), the coupling between plasma and cosmic rays is mainly determined by the gradient of ACR pressure and the PUI injection efficiency at plasma flow compressions. One can use the diffusive length \(L_d = \kappa_k/\nu_p\) to estimate the size of the precursor (Florinski et al. 2005), where \(\kappa_k\) is the radial diffusion coefficient (in this case, simply equal to the isotropic energy-averaged diffusion coefficient \(\kappa\)). In the model, \(\nu_p \sim 350 \, \text{km} \, \text{s}^{-1}\) upstream of the TS, which yields diffusive lengths \(L_{d1} = 0.48\) au, \(L_{d2} = 4.8\) au, and \(L_{d3} = 48\) au. These values give rough estimates of the extent of the precursor shown in Figure 4. Due to the limit of grid resolution, for \(\kappa_1\), the size of precursor \(L_{d1}\) is less distinct in the TS upstream from the red curve, compared to that of \(\kappa_2\) colored as green. There is no apparent sign of the precursor for \(\kappa_3\) because the coupling between plasma and ACRs becomes very weak due to an extremely low intensity of cosmic rays (see panel (D) of Figure 6). Only the intermediate value of the diffusion coefficients, e.g., \(\kappa_2\), produces clearly identifiable precursors in the plasma flow ahead of the TS.

Based on \(V_2\) plasma observations one can make an estimate of the radial diffusion coefficient based on the width of the shock precursor. For \(L_d = 0.7\) au, the corresponding diffusion coefficients will be \(\sim 3.6 \times 10^{20} \, \text{cm}^2 \, \text{s}^{-1}\), which is close to \(\kappa_1\). Likewise, we deduce that \(\kappa_2\) would be the appropriate value of the diffusion coefficient for the \(V_1\) shock encounter. Note that the calculated distance to the TS is smaller by \(\sim 5\) au using the smaller diffusion coefficient. \textit{Voyager 1} crossed the shock 9 au farther than \(V_2\). Here, we argue that this difference could be, in part, attributed to the different energetic particle transport environments during the two \textit{Voyager} shock encounters. Other contributing factors, considered previously, include the
The thermal pressure exists everywhere, while the ACR pressure drops to zero beyond the heliopause as expected. It was assumed that pickup ions are accelerated to ACR energies near the TS and are convected into the inner heliosheath. For a larger diffusion coefficient, the energetic particles are accelerated less efficiently because of a shorter residence time at the acceleration region near TS, with an acceleration rate being proportional to $u_c^2/\kappa$ (Drury 1983). Moreover, larger diffusion leads to a more efficient transport of accelerated particles to the other regions from TS, resulting in a lower ACR pressure as shown in the panel D of Figure 6. This result is similar to the previous work by Alexashov et al. (2004). Simultaneously, the thermal pressure changes in a different manner, finally approaching that of the no ACRs case for the largest diffusion $\kappa_3$.

As mentioned above, we interpret $V_1$ observation as the case of $\kappa_2$ and $V_2$ as that of $\kappa_1$. In the pressure plots shown in panels B and D of Figure 6, the ACR pressure for $\kappa_2$ is significantly smaller than that for $\kappa_1$ in the inner heliosheath. The situation changes in the unshocked solar wind before TS crossing, where the ACR pressure for $\kappa_2$ is much larger than that for $\kappa_1$. This pressure reversal of the two cases is found to be similar to the behavior of intensities of energetic particles with low energies from $V_1$ and $V_2$ observations. The intensities of energetic particles below $\sim 10$ MeV measured by $V_2$ in the heliosheath was about twice that of $V_1$, which is shown in Figure 2 of Stone et al. (2008). Note that the intensities above 10 MeV were also larger during the $V_2$ crossing, but not at the same time. Conversely, upstream of the TS, the intensities of energetic particles between 0.14 and 0.22 MeV observed by $V_1$ in the assumed precursor region were higher than those measured by $V_2$ (see Figure 2 of Decker et al. 2008).

The solar activity declined from the 2004 to 2007, and as a result the turbulent fluctuation level in solar wind was expected to decrease during that period. The relevant quantity is $A = (\delta B^2)^{1/2}/B$, where $\delta B$ is the amplitude of the magnetic fluctuations. Here, we assume that the parallel diffusion coefficient of cosmic rays $\kappa_{||} \sim B^{3/2}/\delta B^2$ based on the quasilinear theory (QLT; Zank et al. 1998). Observations at 1 au shows that $\kappa_{||}$ was globally increasing from 2004 to 2010 (Mewaldt et al. 2010). In the outer heliosphere, however, perpendicular diffusion of cosmic rays described by $\kappa_{\perp}$ may become more important than the parallel diffusion because of

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**4.2. Heliosheath and Heliopause**

We now turn our attention to the HP position and the thickness of the inner heliosheath. We use a passive tracer approach to label the solar wind and interstellar medium, which have different origins, so that the interface between the two fluids is identified as the HP. The positions of the HP for the four simulations discussed here are shown in Figure 4 with vertical dotted lines of the same color as the radial profile of the physical quantity plotted. The HP is located between 120 au and 126 au. While this is similar to the $V_1$ heliopause crossing distance of 122 au, it should be noted that the distances to the TS in all four models are under 80 au, which is not consistent with Voyager observations.

The HP becomes narrower as the diffusion coefficient increases from $\kappa_1$ to $\kappa_2$, and then expands again for $\kappa_3$, where it is located at the same distance as in the model without the energetic particles. One can see that the thickness of the inner heliosheath for $\kappa_2$ is the smallest among all the cases. This can be seen in greater detail in Figure 5, which shows the shapes of the HP and the TS positions for the four cases in the meridional plane (left panel) and the equatorial place (right). Comparing with the larger diffusion cases ($\kappa_2$ and $\kappa_3$), the TS has a smaller radius for the smaller diffusion coefficients ($\kappa_1$). This result is similar to the previous work by Alexashov et al. (2004) and indicates that in the weak diffusion case cosmic rays accumulate in the inner heliosheath, and hence compress the TS with larger cosmic-ray pressure (see Figure 7). Note that in the downstream direction the TS is located closer to the Sun by a distance of $\sim 20$ au for $\kappa_1$ compared with the strong diffusion cases.

Because the ACRs are treated as a massless fluid, they are completely described by their energy density (i.e., pressure). In Figure 6, the radial profiles of the dynamic pressure $P_d = \rho V^2/2$, the thermal pressure $P_T$, the magnetic pressure $P_B = B^2/8\pi$, and the ACR pressure $P_{\text{ACR}}$ along the $V_1$ direction are shown in the four panels, respectively. Note that the magnetic pressure can be usually neglected near the TS and the dynamic pressure can be always neglected at the HP. The sum of thermal and ACR pressures is expected to determine the positions of the two discontinuities because the other two pressures do not dramatically change for different diffusion coefficients as shown in the figure.

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The solar activity declined from the 2004 to 2007, and as a result the turbulent fluctuation level in solar wind was expected to decrease during that period. The relevant quantity is $A = (\delta B^2)^{1/2}/B$, where $\delta B$ is the amplitude of the magnetic fluctuations. Here, we assume that the parallel diffusion coefficient of cosmic rays $\kappa_{||} \sim B^{3/2}/\delta B^2$ based on the quasilinear theory (QLT; Zank et al. 1998). Observations at 1 au shows that $\kappa_{||}$ was globally increasing from 2004 to 2010 (Mewaldt et al. 2010). In the outer heliosphere, however, perpendicular diffusion of cosmic rays described by $\kappa_{\perp}$ may become more important than the parallel diffusion because of...
the nearly azimuthal magnetic field in the inner heliosheath. Assuming that the perpendicular diffusion coefficient is proportional to the turbulence amplitude $A$, e.g., $\kappa_\perp = 0.5 A^2 \kappa_{||}$ (Le Roux et al. 1999), one obtains $\kappa_\perp \propto B^{-1/3}$. It takes around 1 year for the solar wind to reach the TS, thus one might expect that the perpendicular diffusion coefficient during the $V1$ TS crossing was larger than that for $V2$, which would justify our assumption that the radial diffusion coefficient during the $V1$ shock crossing was larger than during the $V2$ shock encounter.

However, the intensities of GCRs measured by $V2$ are larger than those by $V1$ near the TS in the heliosheath (Stone et al. 2008), indicating that the averaged diffusion coefficients for GCRs are less for $V1$ than $V2$. This result seems to be the opposite of the above assumption for ACRs. We argue that the diffusion nature for ACRs can be different from that of GCRs because of their dependence on energies or rigidities, the diffusion coefficients for $V1$ may increase with energies at a more gradual manner than those for $V2$, and finally, the latter will exceed at relatively high energies (Burger & Hattingh 1998; Qin & Schalchi 2012).

The observations of 48 s averaged magnetic field magnitude in the post-TS region by the two probes appear to be more complicated than the expected global scenario of turbulence from declining solar activity. For $V1$, the fluctuations of magnetic field was $\delta B = 0.033$ nT, and the average magnetic field was $B = 0.123$ nT, thus $\delta B/B \sim 0.28$ in the 20 day interval after the HP crossing (Burlaga et al. 2006). The turbulence amplitude increases to 0.35 ($\delta B = 0.054$ nT and $B = 0.154$ nT) in the unipolar magnetic field region, and $V2$ observations yielded an even greater value of 0.49 ($\delta B = 0.045$ nT and $B = 0.091$ nT; Burlaga & Ness 2009). This increase in turbulent fluctuations over time appears to be inconsistent with what could be anticipated based on declining levels of solar activity. Based on the measured $B$, we may conclude that $\kappa_\perp$ was larger for the $V1$ case compared with $V2$ by a factor of 2 in the unipolar region, but was smaller than $V2$ by a ratio of 0.4 in the post-TS region.

The left panel of Figure 7 presents the model-derived profiles of the total pressure $P_t = P_g + P_r$, and the ACR pressure, $P_r$, along the $V1$ direction. In the weak diffusion scenario of $\kappa_1$, the energetic particles are effectively accelerated near the TS and are accumulated in the inner heliosheath, resulting in the increased total pressure compared with the case without the ACRs. For the case of $\kappa_2$, the effects of diffusion and convection are comparable, and the total pressure is higher than without the ACRs near the TS, but becomes lower near the HP. We may also infer that the decreased total pressure for $\kappa_2$ near the HP is the main reason for a narrower HP among all the cases. As before, the total pressure for the highest diffusion case finally vanishes in the interstellar medium.

The right panel of Figure 7 shows the corresponding ratio of $P_r/P_t$. This has a maximum at the TS where the ACRs are injected. Here, we focus on the results for the inner heliosheath since it is the main region where ACRs are convected after their birth from PUIs near the TS. For the low diffusion case of $\kappa_1$, the pressure ratio is 0.45–0.55 in the sheath, which is larger than the observational estimate of 0.25, as mentioned above.
Figure 6. Model-derived radial profiles of (A) plasma dynamic pressure, (B) thermal pressure, (C) magnetic pressure, and (D) ACR pressure along the V1 direction, with units of pdyne cm$^{-2}$.

Figure 7. Left: radial profiles of the total pressure $P_t = P_g + P_c$ (solid lines) and ACR pressure $P_c$ (dashed) along the V1 direction for the four diffusion cases. Right: the corresponding ratio of the ACR pressure to the total pressure.
The case of $\kappa_2$ yields a ratio of $\sim 0.15$–$0.22$, which appears to be feasible.

### 4.3. PUI Injection Rates

We now investigate the case of $\kappa_2$ by varying the PUI injection rates. We ran two additional simulations with injection efficiency $\alpha'$ taken to be 0.3 and 0.5. As the left panel of Figure 8 shows, increasing $\alpha'$ leads to enhanced ACR pressure and a more pronounced precursor ahead of the TS. The thermal pressure $P_t$ decreases from $\sim 1.3$ to 1.0 pdyne cm$^{-2}$, as more thermal particles are converted to ACRs. Notice that the HP shrinks as $\alpha'$ increases. For the very strong injection case the HP moves inward by 6 au relative to the case without the ACRs. Note that the solar wind boundary conditions used here are appropriate for the past cycle 23/24 solar minimum; the solar wind speed of 388 km s$^{-1}$ and number density of 5.5 cm$^{-3}$ are relatively small compared to those used in the simulations performed by other groups, where the HP is usually located at larger distances, 130–157 au (Linde et al. 1998; Pogorelov et al. 2004; Izmodenov et al. 2005; Opher et al. 2006).

The above results demonstrate that the separation of ACRs from the plasma background could result in a narrower width of the heliosheath, which could help explain the unexpectedly early crossing of the HP by $V_1$. The TS moves outward by $\sim 2$–3 au for the high injection case that can be seen from Figure 8, showing that the mediation of ACRs creates noticeable changes in the geometry of the heliosphere. The difference in the shock distance between $\kappa_2$ to $\kappa_1$ increases to 7–8 au.

The corresponding ratio of the ACR pressure ($P_{\alpha}$) to the total pressure ($P_t$) is plotted in the right panel of Figure 8. It can be seen that higher injection efficiency corresponds to a larger ratio in the sheath, where the ratio gradually decreases in the radial direction. The ratio has a minimum of $\sim 0.1$–$0.3$ before the HP crossing for the three cases. This number is comparable to the ratio of $\sim 0.25$ deduced from observations and coincides with the percentage of the partial pressure over the total pressure in the energy range beyond 6 KeV in the inner heliosheath, based on IBEX ENA spectra (Livadiotis et al. 2013). The ratio increases rapidly upstream of the TS where the ACR gradient is large. Faster diffusion also causes a gradual decrease of ACR pressure with heliocentric distance in the sheath, which can be seen in the left panel. Such a decrease is not consistent with the observation of high-energy ACR fluxes in the sheath (Stone et al. 2008). It is possible that a more sophisticated time-dependent treatment of diffusion in the heliosheath could reproduce the positive gradient of the ACR pressure caused by the decreasing solar activity in the sheath, in order to compare the numerical results with the observations in some detail (Senanayake et al. 2012).

The above simulation results are based on the assumption that ACRs and plasma are decoupled in the interstellar medium. The ACR pressure dropped to a very low level once $V_1$ crosses the HP, as shown in Figure 7. Enabling diffusion in the LISM and the source term $S$ in Equation (1) beyond HP will make ACRs and plasma coupled everywhere. Figure 9 shows the numerical results for the cases when plasma and ACRs are coupled beyond the HP. In these cases, ACRs diffuse uniformly at three constant diffusion coefficients $\kappa_{1-3}$. The left panel shows the radial profiles of the total pressure $P_t = P_{\alpha} + P_{e}$ (solid lines) and ACR pressure $P_{\alpha}$ (dashed) along the $V_1$ direction. Different from the previous cases, the ACRs penetrate a great distance ($\sim 180$ au for $\kappa_1$ and $\sim 300$ au for $\kappa_2$) into the LISM. The gradual decrease of $P_{\alpha}$ does not match the observation shown in Figure 2 and indicates that the contributions from $\nabla P_{\alpha}$ in Equation (3) is much less than this model predicts.

The inward short of the HP becomes $\sim 1$ au for $\kappa_2$ and $\kappa_3$, which is much less than in the uncoupled simulations. For $\kappa_1$, the HP moves out by $\sim 3$ au. The radial plasma flow speed along the $V_1$ direction is presented in the right panel of Figure 9. The change in the position of the TS becomes less pronounced as well ($\sim 3$ au or less). The above numerical results indicate that the effects of cosmic rays on the structure of the outer heliosphere is weaker for the uniform diffusion cases. This is the consequence of momentum transfer from the ACRs to the LISM plasma that occurs beyond the HP.
on these numerical experiments, we propose that the ACRs–plasma decoupled model should be used in the LISM, where ACRs escape rapidly along the magnetic field lines and eventually deposit their momentum far from the heliospheric interaction region.

5. Summary

In this work, we carried out global numerical simulations of the heliosphere to investigate the effects of accelerated ions on the shape and size of the outer heliosphere. A simple diffusion model of cosmic rays including an isotropic energy-averaged diffusion coefficient was incorporated into an existing global MHD–neutral model of the heliosphere and the surrounding LISM. Energetic particles were introduced near the TS for a range of injection rates and were followed through the heliosphere using a range of values for the diffusion coefficient. The following statements summarize our findings.

1. Compared to the background plasma-only case, the thickness of the inner heliosheath increased for the relatively small diffusion coefficient of $\kappa_1 = 2.5 \times 10^{20}$ cm$^2$/s$^{-1}$, and decreased for the larger diffusion coefficient of $\kappa_2 = 2.5 \times 10^{21}$ cm$^2$/s$^{-1}$. For the even faster diffusion case, $\kappa_3 = 2.5 \times 10^{22}$ cm$^2$/s$^{-1}$, the ACR effects became minor and the inner heliosheath returned to the state obtained in the no-ACR simulation. The largest reduction in the width of the HP was $\sim 6$ au, which implies that the ACR effect possibly contributed to the unexpectedly early HP encounter by V1 at 122 au from the Sun.

2. The location of the TS varied for different diffusion environments, which may be used to interpret the difference in TS crossings by V1 and V2. Some observations show that V1 may have observed a significantly longer plasma precursor before the TS crossing compared to V2. Intensities of energetic particles with energies of a few hundred keV detected by V1 before the TS crossing were much larger than those measured by V2, but the former became less than the latter after the TS crossing. Based on the simulations, the above observations could be explained if the V1 TS crossing occurred when the ACR radial diffusion coefficient in the heliosphere was larger ($\kappa_2$) than during the V2 crossing ($\kappa_1$). The distance to the TS for $\kappa_3$ was larger than that of $\kappa_1$ by $\sim 5$ au. We conclude that diffusive effects of ACRs might play an important role in the timing and distance of the termination shock and heliopause encounters by the two Voyager probes.

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