The Astrophysical Journal, 744:127 (9pp), 2012 January 10 doi:10.1088/0004-637X/744/2/127

ACCELERATION OF GALACTIC COSMIC RAYS IN THE INTERSTELLAR MEDIUM

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Received 2011 September 28; accepted 2011 November 22; published 2011 December 20

ABSTRACT

Challenges have arisen to diffusive shock acceleration as the primary means to accelerate galactic cosmic rays (GCRs) in the interstellar medium. Diffusive shock acceleration is also under challenge in the heliosphere, where at least the simple application of diffusive shock acceleration cannot account for observations. In the heliosphere, a new acceleration mechanism has been invented—a pump mechanism, driven by ambient turbulence, in which particles are pumped up in energy out of a low-energy core particle population through a series of adiabatic compressions and expansions—that can account for observations not only at shocks but in quiet conditions in the solar wind and throughout the heliosheath. In this paper, the pump mechanism is applied to the acceleration of GCRs in the interstellar medium. With relatively straightforward assumptions about the magnetic field in the interstellar medium, and how GCRs propagate in this field, the pump mechanism yields (1) the overall shape of the GCR spectrum, a power law in particle kinetic energy, with a break at the so-called knee in the GCR spectrum to a slightly steeper power-law spectrum. (2) The rigidity dependence of the H/He ratio observed from the *PAMELA* satellite instrument.

Key words: acceleration of particles – cosmic rays – ISM: magnetic fields – ISM: supernova remnants – Sun: heliosphere

1. INTRODUCTION

The acceleration of galactic cosmic rays (GCRs) in the interstellar medium has been a subject of interest and extensive research for many decades. Prevailing theories for the processes by which the GCRs are accelerated usually involve some form of diffusive shock acceleration at shocks generated by supernovae. This theory came into prominence in the late 1970s, with the seminal work of Axford et al. (1977), Krymsky (1977), Bell (1978), and Blandford & Ostriker (1978), who showed that diffusive shock acceleration naturally yields power-law spectra similar to the observed GCR spectrum. Supernovae shocks are also among the most energetic processes in the interstellar medium and certainly a possible source of energy for the GCRs.

In recent years, however, challenges have arisen to diffusive shock acceleration, in its simplest form, as the mechanism to produce GCRs (e.g., Butt 2009). Isolated, large supernovae remnants (SNRs) may be too rare, introduce anisotropies not observed, and may not be sufficiently large nor have sufficient energy to accelerate very high energy GCRs that have gyroradii larger than the supernovae shock (e.g., Lagange & Cesarsky 1983). Moreover, recent observations from the *PAMELA* satellite instrument have revealed structure in the GCR spectrum in the magnetic rigidity range between 5 and 1000 GV that appears to be inconsistent with the expected spectra from diffusive shock acceleration (Adriani et al. 2011). The ratio of H to He is found to vary with rigidity, R, falling off as $R^{-0.1}$, whereas the expectation from diffusive shock acceleration in SNRs is that the ratio would be independent of rigidity (Schwarzschild 2011).

A similar disquiet with regard to diffusive shock acceleration has arisen in heliospheric physics, where diffusive shock acceleration is routinely applied to acceleration at the many shocks that are observed in the solar wind. There is, for example, relatively little compelling observational evidence that diffusive shock acceleration, at least in its simplest form, is the dominant acceleration mechanism at shocks. In simple diffusive shock acceleration, the spectral index of the accelerated particles depends upon the compression ratio of the shock. Yet the correlation between the Mach number of the shock, and thus the compression ratio, and the spectral index of the accelerated particles is generally not seen. For example, van Ness et al. (1984) found a correlation between the spectral index and the compression ratio for 75% of the shocks observed, but only when fairly generous brackets were put upon the compression ratio. A more recent study by Desai et al. (2004) showed little or no correlation for the spectral index of oxygen accelerated at shocks.

The dissatisfaction with diffusive shock acceleration in the solar wind became acute with the recent discovery by Gloeckler & Fisk (2011) that downstream from the shocks observed by the *ACE* spacecraft at 1 AU in 2001, where acceleration is occurring, the spectral index of accelerated particles is usually a power law with a spectral index of −5, when expressed as a distribution function, irrespective of the compression ratio of the shock. The −5 spectrum is known as the common spectral shape and is observed in many different plasma conditions in the heliosphere (Fisk & Gloeckler 2008, 2009; Fisk et al. 2010; Gloeckler & Fisk 2006; Gloeckler et al. 2008). For example, it is the dominant spectral shape of particles accelerated in the heliosheath, currently being explored by the *Voyager* spacecraft (Decker et al. 2006; Gloeckler et al. 2008).

The common spectral shape cannot be the result of diffusive shock acceleration. Not only is the common spectral shape inconsistent with the standard prediction of diffusive shock acceleration that the spectral index is correlated with the shock compression ratio, but also the common spectral shape is observed in the quiet solar wind, far from any shocks. The common spectral shape can also not result from traditional stochastic acceleration in which particles gain energy by diffusing in velocity space. The spectra that result from traditional stochastic acceleration are not normally power laws and certainly not power laws with a specific spectral index of −5.

The inability of diffusive shock acceleration and traditional stochastic acceleration to account for the common spectral shape led Fisk & Gloeckler (2006, 2007, 2008, 2011a, 2011b; Fisk et al. 2010) to develop a new acceleration mechanism—a pump mechanism, driven by plasma turbulence, in which particles are pumped up in energy through a series of adiabatic compressions and expansions, in which the particles can escape from a compression region, or flow into an expansion region by spatial diffusion. The mechanism is a redistribution mechanism in which the energy in a low-energy, but hot (suprathermal) core particle population is redistributed to higher energies, without the damping of turbulence. The mechanism is shown to yield naturally a −5 spectrum independent of the plasma conditions, and it contains a first-order acceleration that makes the mechanism particularly efficient and able to explain the observations of particles accelerated at shocks in the solar wind (Gloeckler & Fisk 2011).

It is thus worthwhile to explore whether the pump acceleration mechanism of Fisk & Gloeckler, which is successful in the solar wind, might also be applicable to the acceleration of GCRs in the interstellar medium, and as with the solar wind, relieve the concerns with diffusive shock acceleration as the primary mechanism for accelerating GCRs. The conditions required for the pump acceleration mechanism of Fisk & Gloeckler are fairly easy to satisfy. There needs to be a suprathermal core distribution of particles, which contains sufficient energy to be redistributed and account for the energy in the GCRs. The hot thermal plasma in superbubbles could be an adequate core particle population. There needs also to be large-scale compressions and expansions of the plasma, which the subsonic interstellar medium might readily contain.

In this paper, we apply the pump mechanism of Fisk & Gloeckler to the acceleration of GCRs in the interstellar medium for the purpose of explaining two distinct observations: (1) the overall shape of the GCR spectrum, a power law in particle kinetic energy, with a break at the so-called knee in the GCR spectrum, to a slightly steeper power-law spectrum and (2) the rigidity dependence of the H/He ratio observed from *PAMELA*. We need also to satisfy certain constraints such as the observed lifetime of GCRs.

We begin by summarizing the observations that we will explain. We then specify the model for the interstellar magnetic field and the propagation of GCRs in this field that we will use. We describe the pump acceleration mechanism of Fisk & Gloeckler and the equation that describes the acceleration of GCRs in the interstellar medium (the relativistic form of the governing acceleration equation of Fisk & Gloeckler is derived in the Appendix). We point out that the pump mechanism is an energy distribution mechanism, and thus the energy that is imparted to GCRs originates in the hot, tenuous plasma in superbubbles, not from the damping of turbulence. We solve the governing acceleration equation and show that the resulting spectra can account for the observed GCR spectrum, and it can accomplish the acceleration within a reasonable lifetime for GCRs. We then consider relatively low-rigidity GCRs (∼5–100 GV) and show that the pump mechanism can yield the rigidity dependence of the H/He ratio observed from *PAMELA*. Finally, we discuss, in general terms, how the pump acceleration mechanism of Fisk & Gloeckler could yield results that are consistent with other observations of GCRs such as the composition and apparent spatial variations, revealed by enhancements in the gamma rays that GCRs produce, as well as small anisotropies.

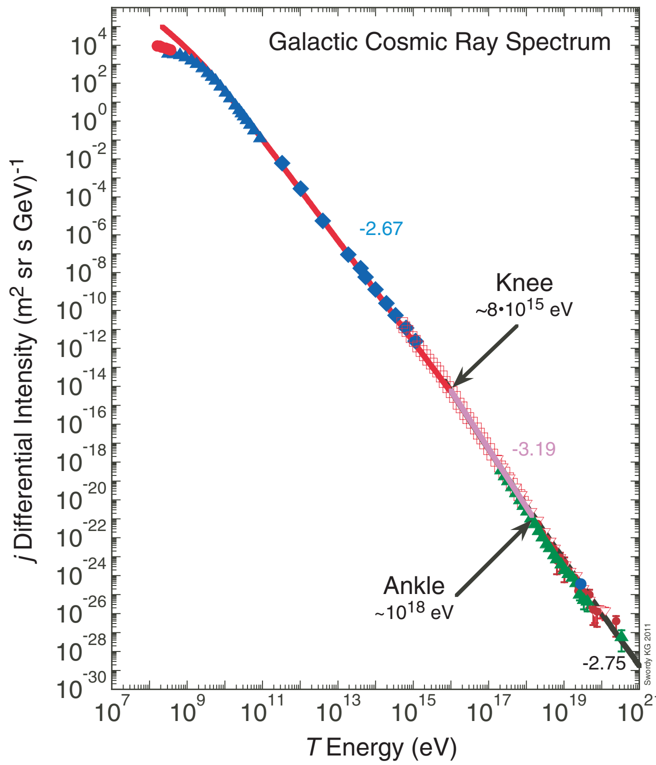

Figure 1. Differential intensity, j, vs. kinetic energy, T, of galactic cosmic ray protons (symbols) and fits to the spectrum computed from Equation (5) for $T < 8 \times 10^{15}$ eV and Equation (8) for $8 \times 10^{15} \leqslant T < 10^{18}$ eV (red curve). The relation $j = p^2 f$ was used to convert phase-space density, f, to j. The spectrum above the ankle ($T \geqslant 10^{18}$ eV), assumed to be of extragalactic origin, was fit using a power law with a spectral index of −2.75. Filled circles are *Voyager 1* measurements (Cummings et al. 2008), while data represented by all other symbols are a compilation of various measurements by Swordy (2001).

2. THE OBSERVATIONS

There are two distinct observations of GCRs that we intend to explain.

1. Shown in Figure 1 is a composite spectrum for the differential intensity of GCRs (number of particles crossing unit area in unit time per unit of particle kinetic energy; Swordy 2001; Cronin et al. 1997). At energies below $\sim 5 \times 10^9$ eV the spectrum is altered by heliospheric effects, the modulation of GCRs by the solar wind. At energies above $\sim 5 \times 10^9$ eV the spectrum is a power law with spectral index ∼−2.7, extending up to the so-called knee in the cosmic-ray spectrum at $\sim 8 \times 10^{15}$ eV, at which the spectrum steepens to another power law with spectral index ∼−3.15 that extends up to the so-called ankle at $\sim 10^{18}$ eV. We assume that particles with energies above $\sim 10^{18}$ eV are of extragalactic origin and concentrate on explaining the spectrum between 5×10^9 eV and 10^{18} eV.
2. The observations from *PAMELA* reveal that in the magnetic rigidity range from $R \sim 5$ GV to ∼200 GV the H to He ratio varies as $R^{-0.1}$. There are also indications that there are breaks in the H and He spectra at ∼150 GV, where the spectra become noticeably harder.

We also need to impose the constraint that the choices of parameters in our model yield the escape lifetime of GCRs of ∼15 million years, as inferred from radioactive isotopes

(Mewaldt et al. 2001). The observed radioactive isotopes have energies of several hundred MeV nucleon^{-1}. Since these particles suffer substantial adiabatic deceleration in the solar wind, this is the escape lifetime of particles with energies ∼1 GeV.

3. A MODEL FOR THE INTERSTELLAR MEDIUM

The pump acceleration mechanism requires a low-energy core particle population, containing energy in excess of the energy in GCRs. The pump mechanism extracts part of the energy from the core to create the GCRs, without damping of the turbulence. Thus, the pump mechanism does not have the weakness often associated with traditional stochastic acceleration where the energy placed in the accelerated particles is extracted from the energy in turbulence, which is insufficient to account for the energy in the GCRs.

The core particle population that we invoke is the hot thermal plasma in superbubbles: densities of <0.01 cm^{-3} and temperatures of $>10^6$ K (e.g., Chu 2007). Superbubbles appear to be expanding and thus have a pressure in excess of the average pressure in the interstellar medium, or a pressure in excess of the ∼1 eV cm^{-3} in GCRs. The thermal speeds of the hot plasma should readily allow the particles to be injected into the pump mechanism. The low plasma density will not result in significant ionization losses.

We assume that particles accelerated by the pump mechanism in superbubbles then spread into surrounding denser regions. In these denser regions low-energy particles suffer ionization losses. Thus, only particles with energies above several hundred MeV nucleon^{-1}, which should suffer negligible ionization losses (Gloeckler & Jokipii 1967), can be expected to spread from the superbubbles into the surrounding Galaxy. At these higher energies the particles should spread roughly uniformly throughout the Galaxy and then continue to be accelerated to higher energies throughout the entire Galaxy.

The interstellar medium contains a mean magnetic field, which lies in the plane of the Galaxy. We assume that there are compressions and expansions of the interstellar medium, which are aligned with the mean magnetic field, and have a cross-sectional diameter of l and a coherence length along the mean magnetic field of R_l. We take the compression and expansion regions to be highly elongated, with $R_l \gg l$.

Magnetic field lines in the turbulent interstellar plasma are expected to random walk. We assume that field lines can diffuse to the edge of a compression/expansion region, i.e., diffuse a distance $l/2$ within a distance R_l. The mean magnetic field in contrast must be coherent to a higher order than the random field. The random field diffuses by a distance $l/2$ over a length scale R_l. The corresponding coherence length of the mean field along the mean field direction, i.e., the scale length over which the dispersion of the mean field has grown to $l/2$, must thus be $\Delta z = R_l^2/l$.

We assume that the GCRs are transported in the interstellar medium by cross-field diffusion due to particles following random walking magnetic field lines (Jokipii & Parker 1969). Cross-field diffusion is the means by which particles can escape from a compression region or flow into an expansion region. Cross-field diffusion is the means by which particles can escape from the Galaxy.

For GCRs with gyroradii $r_g < l$, the cross-field diffusion coefficient for escape from a compression or flow into an expansion, κ_l, and for escape from the Galaxy, κ_g, should be the same and given by

$$\kappa_l = \kappa_g = \frac{1}{4\sqrt{3}}\frac{l^2}{R_l}v = \frac{1}{4\sqrt{3}}\frac{l^2}{R_l}\frac{pc}{\left(m_o^2c^2 + p^2\right)^{1/2}} \quad \text{for} \quad r_g < l. \tag{1}$$

Here, v is particle speed and we assume that the particle distribution is quasi-isotropic with the speed of the particle along the magnetic field, $v/\sqrt{3}$. On the right-hand side of Equation (1) we have expressed particle speed, v, in terms of particle momentum, p, with m_o being the rest mass of a particle and c the speed of light.

GCRs with gyroradii $r_g > l$, particularly particles with gyroradii very much larger than l, can escape from a compression region or flow into an expansion region simply by crossing the compression/expansion region in a straight line. The effective diameter of the compression region is $l/2$, in which case the average distance across this diameter in any direction is l/π, and the equivalent cross-field diffusion coefficient for escape from a compression or flow into an expansion is

$$\kappa_l = \frac{\pi l c}{4\sqrt{3}} \quad \text{for} \quad \mathrm{r}_g > l, \tag{2}$$

where we have assumed that particles with $r_g > l$ are highly relativistic with $v = c$.

GCRs with gyroradii $r_g > l$ should average over the small-scale fluctuations in the magnetic field and be able to sense only the mean magnetic field. Thus, GCRs with gyroradii $r_g > l$ can escape from the Galaxy only by following the random walk of the mean magnetic field, in which case for these particles the cross-field diffusion coefficient for escape from the Galaxy becomes

$$\kappa_g = \frac{c}{4\sqrt{3}}\frac{l^3}{R_l} \quad \text{for} \quad \mathrm{r}_g > l. \tag{3}$$

Note, in this model for the propagation of GCRs in the interstellar medium it is harder for a particle with gyroradii $r_g > l$ to escape from the Galaxy than for a particle with gyroradii $r_g < l$. The latter particles can follow the fine-scale random walking field, which provides avenues for escape. Particles with gyroradii $r_g > l$ average out and cannot follow the fine structure magnetic field; rather they can only follow the more coherent mean magnetic field, which lies in the plane of the Galaxy.

3.1. The Energy Source for GCRs in the Pump Acceleration Mechanism

There is a very important feature of the pump acceleration mechanism that needs to be emphasized. As can be seen in the derivation of Equation (4) in the Appendix, the pump mechanism is an energy redistribution mechanism; the energy in the core particles, in this case in the hot tenuous plasma of a superbubble, is redistributed to form the accelerated GCRs. The pump mechanism is an irreversible process, in which entropy increases, but the energy in the core particles plus the GCRs is constant.

Thus, unlike traditional stochastic acceleration processes, in which particles are accelerated by diffusing in momentum space, the GCRs accelerated by the pump mechanism do not acquire their energy by the damping of turbulence. Rather, the source of the energy in the GCRs is the energy in the core particles. For this reason, superbubbles are an ideal location in which to apply the pump acceleration mechanism. The hot tenuous plasma in superbubbles could well contain sufficient energy to provide the energy required by the accelerated GCRs. Indeed, *Fermi* observations of the Large Magellanic Cloud appear to show enhancements in gamma rays, and thus in GCRs, correlated with superbubbles (Abdo et al. 2010).

The pump acceleration mechanism has been applied by Fisk & Gloeckler (2009) to the acceleration of anomalous cosmic rays (ACRs) in the heliosheath. The ACRs are accelerated by pumping energy out of the core particle population of interstellar pickup ions. To fit the observed spectra of ACRs it is necessary to extract $\sim 1/3$ of the pickup ion energy to form the ACRs. There are similarities between the heliosheath and our model for the interstellar medium; each is a subsonic gas, with compressions and expansions, and the presence of a hot, tenuous core particle population. It might be reasonable to assume then that to account for the observed GCR spectra by the pump mechanism it is necessary to extract $\sim 1/3$ of the energy in hot tenuous gas in superbubbles.

4. THE PUMP ACCELERATION MECHANISM OF FISK AND GLOECKLER

The situation in which the pump mechanism of Fisk & Gloeckler accelerates particles is as follows. There is a volume of plasma containing random compressions and expansions. Three particle populations are present in the volume: (1) the thermal plasma, e.g., the thermal particles in the interstellar medium, which contains the mass and is responsible for the random compressions and expansions; (2) a particle population with particle momentum greater than the thermal momentum of the bulk plasma and less than an upper threshold momentum, $p \leqslant p_{\rm th}$, which we refer to as the core particle population. The core particles also undergo random compressions and expansions, but are not sufficiently mobile and do not readily escape by spatial diffusion from a compression/expansion region, e.g., a kappa distribution on the thermal plasma could serve this role; (3) the suprathermal tail, which has particle momentum above $p_{\rm th}$ and is being pumped up out of the core particle population (i.e., the tail gains particles and energy from the core). The distinction between the core and the tail is that the more energetic tail particles can spatially diffuse and escape from a compression or diffuse into an expansion region.

The physical principles of the pump mechanism are as follows. (1) In compression regions, the core particles are compressed adiabatically and energy and particles flow across the threshold boundary from the core into the tail. The tail particles are also compressed adiabatically, and raised in energy. (2) The opposite occurs in the expansion regions surrounding the compression region. In the expansion regions, particles and energy flow from the tail back into the core and the energy of the tail particles is reduced. (3) As a result of the energy being gained in compression regions, and lost in the surrounding expansion regions, large spatial gradients will result at higher particle energies between the compression and the surrounding expansion regions. (4) Tail particles are able to spatially diffuse, and so at higher particle energies, particles will diffuse in response to these gradients out of the compression region into the surrounding expansion regions. (5) Subsequently, the compression region will become an expansion region, and the process will be reversed. Particles and energy will flow back into the core from the tail. However, since particles have escaped from the tail by diffusion when it was a compression region, there are fewer particles and less energy to return to the core.

If the process of compressions and expansions is repeated sequentially, then a suprathermal tail will form. The particles in the tail and the energy they contain will systematically increase in time. This is a classic pump mechanism. The combination of adiabatic compressions and expansions, and spatial diffusion of the tail particles, will pump particles out of the core to form a suprathermal tail.

In the Appendix, we derive the relativistic form of the pump acceleration equation of Fisk & Gloeckler. We assume that the acceleration is balanced by escape from the Galaxy, with a characteristic escape distance of $R_g/2$, where R_g is the thickness of the galactic disk. Thus, in a steady state the governing equation for the distribution function $f(p)$ of GCRs accelerated in the interstellar medium is

$$\frac{1}{vp^3}\frac{\partial}{\partial p}\left(\frac{\delta u^2}{9\kappa_l}p\frac{\partial}{\partial p}(vp^4 f)\right) = \frac{4\kappa_g}{R_g^2}f. \quad (4)$$

Here, δu^2 is the mean square speed of the compressions and expansions; κ_l and κ_g are given in Equations (1)–(3); and recall that v in terms of p is $v = pc/(m_o^2c^2 + p^2)^{1/2}$.

5. THE GCR SPECTRUM

Equation (4) has a ready solution for particles with gyroradii $r_g < l$, and thus for the κ_l and κ_g given in Equation (1):

$$f = f_o\left(\frac{p}{p_o}\right)^{-5}\frac{\left(m_o^2c^2+p^2\right)^{1/2}}{m_oc} \times \exp\left[-\beta\int_{p_o}^{p}\frac{dp}{\left(m_o^2c^2+p^2\right)^{1/2}}\right] \quad \text{for} \quad r_g < l, \quad (5)$$

where

$$\beta = \frac{\sqrt{3}l^2c}{2\delta u R_l R_g}. \quad (6)$$

Here, f_o is the value of f, and p_o $(= p_{\rm th})$ is the particle momentum where particles are injected into the acceleration mechanism; we take $p_o \ll m_oc$.

Note that for relativistic particles, the integral in the exponential in Equation (5) results in the log of p, and thus the spectrum of f is a power law in particle momentum, p. We can express f in terms of differential intensity $j = p^2f$, as a function of particle kinetic energy $T = cp$, or Equation (5) becomes

$$j \propto T^{-(2+\beta)} \quad \text{for} \quad r_g < l. \quad (7)$$

For particles with gyroradii $r_g > l$, and thus with κ_l and κ_g given in Equations (2) and (3), respectively, the differential intensity becomes

$$j \propto T^{-(2+\beta')} \quad \text{for} \quad r_g > l \quad (8)$$

where

$$\beta' = \frac{\sqrt{3\pi}}{2}\frac{c}{\delta u}\frac{l^2}{R_l R_g}. \quad (9)$$

Note that β' differs from β by the $\sqrt{\pi}$.

There is evidence from radio observations that the nature of interstellar turbulence changes at a spatial scale of 3.5 pc (Minter & Spangler 1996; Minter 1999). We thus assume that this is an appropriate scale for l, the characteristic diameter of our compression and expansion regions. With this choice for l, and an average magnetic field strength in the interstellar medium of 2 μG, the break in the spectrum between j in Equation (7) and j in Equation (8) coincides with the location of the knee in the GCR spectrum at $\sim 8 \times 10^{15}$ eV.

The observed escape lifetime for mildly relativistic particles is $\tau_{\rm esc} \sim 15$ million years (4.5×10^{14} s) (Mewaldt et al. 2001). Using κ_g in Equation (1), we can then express β as $\beta = (3R_g)/(2\delta u \tau_{\rm esc})$, or with R_g in units of parsecs, and δu in units of km s^{-1}, $\beta = 0.1 R_g/\delta u$. We take R_g equal to 300 pc, and δu equal to 45 km s^{-1}, which may be appropriate if much of the acceleration occurs in superbubbles with their large thermal speeds. The resulting value of β is 0.67, and for β', 1.19. Clearly, other combinations of R_g and δu are possible.

The GCR spectrum predicted by Equation (4), with the κ_I and κ_g from Equations (1)–(3), and with the parameter choices just described, is shown in Figure 1, including the nonrelativistic solution to Equation (4). The nonrelativistic spectrum is not shown below a few hundred MeV since there should be a low-energy cutoff where energy losses from ionization exceed the energy gains from the acceleration. Equation (4) is not valid for particles with gyroradii larger than R_g, which corresponds to energies greater than $\sim 10^{18}$ eV. Particles with energies in excess of $\sim 10^{18}$ eV, e.g., above the ankle of the GCR spectrum, are assumed to be of extragalactic origin.

The predicted spectrum in Figure 1 provides a good fit to the observations.

5.1. *The Efficiency of the Pump Acceleration Mechanism*

It should be noted that the pump acceleration mechanism is quite efficient. It contains a first-order acceleration, which Fisk & Gloeckler (2011b) demonstrate is comparable to or greater than the first-order acceleration in diffusive shock acceleration. There is thus no difficulty in accelerating the GCRs within their lifetime in the galaxy. In Equation (4), we balance the acceleration against the escape, i.e., the characteristic acceleration time equals the characteristic escape time. In Figure 1, we chose the parameters governing the acceleration of GCRs so that for particles with energy below the "knee" in the spectrum (energies less than 8000 TeV), the escape time and thus the acceleration time, which is independent of energy, is 15 million years. It should be noted, however, for particles with energies above the "knee," the acceleration time (determined by Equation (2)) is much longer than 15 million years, but this is compensated for by a much longer escape time from the galaxy (determined by Equation (3)).

Finally, we note the differences between the pump mechanism as it is applied in the heliosphere versus as it is applied in the interstellar medium. In the heliosphere we obtain a power-law spectrum in particle speed with a spectral index of –5, when expressed as a distribution function. In the interstellar medium we attain a power-law spectrum in particle energy with spectra indices of –2.67 and –3.19, when expressed as differential intensity. In the heliosphere we deal with nonrelativistic particles and distribution functions, while in the interstellar medium, we have relativistic particles and differential intensity. Moreover, in the heliosphere, we do not balance acceleration against escape, whereas, for the acceleration of GCRs in the interstellar medium, the spectral shape is determined by balancing acceleration against escape, in Equation (4).

6. THE *PAMELA* OBSERVATIONS

As can be seen from the derivation of Equation (4) in the Appendix, the term $\kappa_I/\delta u^2$ is the average of this term over the multiple compressions and expansions experienced by a particle as it is being accelerated. There is no reason to believe that $\kappa_I/\delta u^2$ is the same in all regions of the Galaxy. It is thus possible that the average of $\kappa_I/\delta u^2$ is different for particles depending upon how long it takes to accelerate the particles and thus how many different compressions and expansions a particle experiences.

We can see from Equation (5) that the time to accelerate a particle to momentum p, $\tau_{\rm acc}$, is proportional to the integral in the exponential in Equation (5), or

$$\tau_{\rm acc} \propto \int_{p_o}^{p} \frac{dp}{\left(m_o^2c^2 + p^2\right)^{1/2}}. \qquad (10)$$

We can express Equation (10) in terms of the particle rigidity, $R = pc/Ze$, where Ze is the electronic charge of the particle. We measure R in units GV, in which case, $m_o c^2/Ze = (A/Z)GV$, where A is mass number. Equation (10) then becomes

$$\tau_{\rm acc} \propto \int_{R_o}^{R} \frac{dR}{((A/Z)^2 + R^2)^{1/2}}. \qquad (11)$$

Note that for a given rigidity, R, the acceleration time is less for He ($A/Z = 2$) than for H ($A/Z = 1$).

We assume that the average of $\kappa_I/\delta u^2$ varies as $(a + b\tau_{\rm acc}/\tau_{\rm max})$, where a and b are constants; $\tau_{\rm max}$ is a sufficiently long acceleration time so that it is reasonable to assume that a particle has experienced a sufficient number of compressions and expansions to be able to experience the average acceleration conditions in the Galaxy. Thus, the average of $\kappa_I/\delta u^2$ appropriate for the Galaxy as a whole is proportional to $(a+b)$. The local value of $\kappa_I/\delta u^2$, appropriate for particles that have experienced a limited number of compressions and expansions, is proportional to a.

It is not necessary that a be positive. Indeed, we assume that there is no core particle population at low energies in local, denser regions of the Galaxy. Ionization losses preclude a core particle population in the local region, which has an average density ~ 0.3 cm^{-3}. GCRs observed locally originate in a nearby superbubble and then propagate into the local interstellar medium. As can be seen from Equation (4), a negative value of a, and thus a negative value of $(a + b\tau_{\rm acc}/\tau_{\rm max})$ at low energies turns the sink term on the right, due to escape from the Galaxy, into a source term, due to particles diffusing in from a nearby superbubble.

A negative value of a also turns the pump acceleration term on the left of Equation (4) into a deceleration mechanism. Recall the basic principles of the pump mechanism: in a compression, particles and energy flow from the core into the tail, and then escape by spatial diffusion. When a compression subsequently becomes an expansion, particles and energy flow back into the core, but since particles and energy have escaped during the compression phase, less energy and fewer particles are returned to the core. Operating sequentially, particles and energy are pumped from the core into the tail. However, if there is an external source of energetic particles, more particles and energy can flow into a compression region than escape by local diffusion. During the pumping sequence, then, the externally supplied particles are pumped down in energy into the core. The steady-state solution, then, for a negative value of a is a balance between inward diffusion from an external source and deceleration in the pump mechanism. When particles are pumped down in energy into an energy range where ionization losses are significant, then there is a low-energy cutoff on the GCR spectrum.

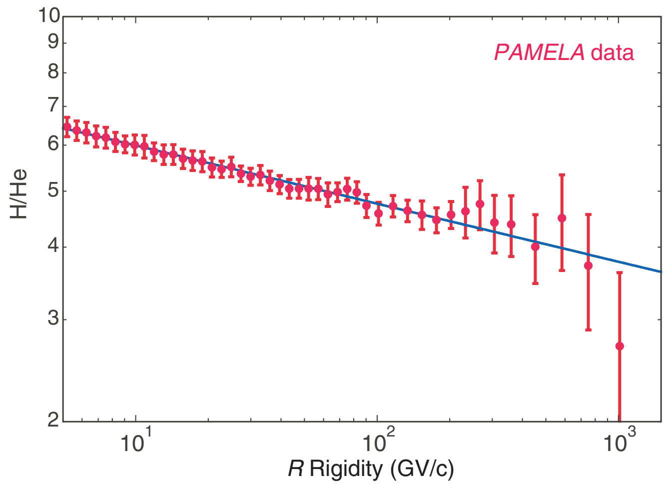

Figure 2. Ratio of protons to helium (H/He) as a function of rigidity, R, measured by the *PAMELA* satellite instrument (Adriani et al. 2011). Error bars are a combination of statistical and systematic errors of the H and He spectra. The blue curve is a fit to the observed ratio using the integral in Equation (12) in the solution to Equation (5) (expressed as a function of rigidity, R), with $a = -0.3$, $b = 1.3$, and $\tau_{\max}$ set equal to the acceleration time of a 200 GV proton.

The integral in Equation (5) thus becomes

$$-\beta \int_{p_o}^{p} \frac{dp}{\left(m_o^2 c^2 + p^2\right)^{1/2}} = -0.67 \int_{R_o}^{R} \frac{(a + b\tau_{\rm acc}/\tau_{\max})dR}{((A/Z)^2 + R^2)^{1/2}}. \quad (12)$$

Here, $a + b = 1$, and the integral on the right-hand side of Equation (12) is valid for $\tau_{\rm acc} < \tau_{\max}$. For values of $\tau_{\rm acc} \geqslant \tau_{\max}$, we assume the particles experience the average value of β and set $a + b\tau_{\rm acc}/\tau_{\max} = 1$.

In Figure 2, we plot the predicted ratio of H to He in the rigidity range from 5 to 200 GV, calculated using the integral in Equation (12) in the solution to Equation (5) (expressed as a function of rigidity, R), with $a = -0.3$, $b = 1.3$, and $\tau_{\max}$ set equal to the acceleration time of a 200 GV proton. Also shown in Figure 2 are the *PAMELA* observations (Adriani et al. 2011). The rigidity dependence of the predicted H/He ratio provides a good fit to the *PAMELA* observations. Note that with this choice for a and b, the crossover between where there is an external source of the local GCRs and where there is escape of the local GCRs occurs at a few GV.

Using the integral in Equation (12) in the solution to Equation (5) modifies the GCR spectrum shown in Figure 1, but only noticeably at energies below ∼5 GeV, where modulation by the solar wind is important. Thus, the modification is only important for determining the level of modulation, e.g., how much modulation still lies beyond the *Voyager* spacecraft.

At rigidities above 200 GV the spectrum of both H and He reverts to the average galactic spectrum with a spectral index of −2.67. This change in the spectra is consistent with the sharp breaks in the spectra at about 200 GV observed by *PAMELA* (Adriani et al. 2011). The breaks are most evident when the differential intensity is multiplied by $R^{2.7}$ (Schwarzschild 2011).

7. OTHER CONSIDERATIONS

The composition of the GCRs accelerated by the pump mechanism will reflect the composition of the core particles and thus should reflect the composition of superbubbles. This is consistent with observations; e.g., Wiedenbeck et al. (2001) find that the composition of GCRs indicates that particles are accelerated from well-mixed interstellar material and do not reflect the elemental anomalies of recent SNRs. Binns et al. (2007) find that isotopic anomalies in GCRs are consistent with particles being preferentially accelerated in superbubbles.

When we deal with acceleration that is distributed throughout the galaxy, the question arises as to whether it is meaningful to discuss source regions of GCRs. In the model developed in this paper, superbubbles are the regions where GCRs are first ejected out of the core of hot tenuous plasma and receive their first acceleration. As the GCRs attain higher energies, sufficient so that they do not suffer ionization loses, they spread throughout the galaxy and continue to be accelerated. However, this acceleration is likely not uniform, anymore than it is uniform in the heliosphere. There are regions where the acceleration is more effective, perhaps the local escape is not as fast as the average escape, and the GCR intensity is enhanced. Since the particles in these enhanced acceleration regions propagate into surrounding regions, one could refer to the enhanced regions as GCR sources. However, since the acceleration occurs throughout the galaxy, they are not isolated sources, but simply part of a continuum of acceleration sites of different strengths. As is observed in the heliosphere, and can be readily modeled, the differences in the local acceleration sites will average out, yielding the average spectrum predicted by Equation (4). For example, propagation by spatial diffusion enters into the equation for the acceleration of GCRs as the divergence of a flux, which when integrated over the volume of the galaxy, yields only the net flow of particles across the outer surfaces of the galaxy. In other words, the effects of spatial diffusion between localized regions of enhanced acceleration integrate to zero inside the galaxy, and the net effect is the average escape from the galaxy.

In these regions of enhanced acceleration, e.g., downstream from shocks, the intensity will be larger than average and the spectrum can deviate from the average spectrum determined by Equation (4). The intensity and the spectral slope are determined by balancing acceleration against escape. The diffusion coefficients that describe the acceleration locally and the local escape are not necessarily the same as those that describe the average acceleration and the average escape. Hence, there could be deviations of local spectra, e.g., inferred from gamma-ray observations, from the average GCR spectrum. We could also expect some anisotropies in GCRs. The pump mechanism is based upon the Parker transport equation, which is valid up to ∼10% anisotropies. In fact, the spectra that yield the H/He ratio observed by *PAMELA* results, in part, from the assumption that there are not local sources of GCRs at low energies, but rather, the GCRs flow into the local region from, e.g., nearby superbubbles.

The pump mechanism for accelerating GCRs should work on electrons equally well as it does on ions. The low-energy (<1–2 GeV) GCR electron spectrum can be determined from the nonthermal radio background (Goldstein et al. 1970) and is often used to estimate the extent to which cosmic rays are modulated by the solar wind (e.g., Webber & Higbie 2008). The inferred low-energy GCR electron spectrum is consistent with $j \propto T^{-2}$. If electrons behave as do ions of the same speed, v, we would expect that the low-energy GCR electron spectrum is $j \propto T^{-2.67}$. However, if electrons, with their much smaller gyroradii, are unable to effectively cross-field diffuse and escape from the galaxy, then the low-energy GCR electron spectrum should be the required $j \propto T^{-2}$.

8. CONCLUDING REMARKS

We have explored whether the pump mechanism of Fisk & Gloeckler, which has been successful in accounting for particle acceleration that is observed in the solar wind, can also account for the acceleration of GCRs in the interstellar medium. We find that with relatively straightforward assumptions concerning the interstellar magnetic field and how GCRs propagate in this field, the pump mechanism can produce the observed GCR spectrum in the range 5×10^9 eV to 10^{18} eV, including the location of the "knee" in the GCR spectrum and the observed change in the power-law slope at the knee. The pump mechanism can also yield a rigidity-dependent H/He ratio consistent with the *PAMELA* observations.

Although the pump acceleration mechanism appears to be sufficient to account for the acceleration of GCRs in the interstellar medium, it is interesting to ask whether it is the only acceleration mechanism. It is certainly likely to be dominant over any traditional stochastic acceleration mechanism, in which particles diffuse in momentum space. The pump mechanism contains a first-order acceleration, and thus is much faster than traditional stochastic acceleration, and it has a more acceptable energy source in the hot tenuous plasmas of superbubbles, as opposed to the damping of turbulence. Diffusive shock acceleration at supernovae shocks could still occur. However, if the heliosphere provides a meaningful analog, the pump acceleration mechanism appears to dominate over diffusive shock acceleration. The advantage in the heliosphere is that particles are accelerated in compressive turbulence that occurs throughout the solar wind, whereas particles gain energy in shocks only by remaining near the shock location (Fisk & Gloeckler 2011b).

There are many observations of GCRs that an acceptable theory for their acceleration needs to be able to explain. In this paper, we have dealt with only two basic observations—the overall shape of the GCR spectrum and the rigidity-dependent H/He ratio. These are the observations that provide the most challenge to the theory in which GCRs are accelerated by diffusive shock acceleration in SNRs. Clearly, to explain the other GCR observations, a more detailed model for GCR acceleration using the pump mechanism needs to be constructed and tested against additional GCR observations.

Finally, we note in Figure 1 that we predict the GCR spectrum in the local interstellar medium in the energy range where cosmic rays are modulated by the solar wind. This spectrum can be compared with *Voyager* observations from the heliosheath to predict the extent to which modulation occurs beyond the *Voyager* spacecraft. In particular, this comparison can provide constraints on how cosmic rays interact with and penetrate through the heliopause.

This work was supported in part by NASA Grant NNX10AF23G and by NSF Grant AGS-1043012. This paper benefited substantially from discussions held at the meetings of the International Team on −5 Tails and ACRs and the Workshop on Particle Acceleration in Cosmic Plasmas held at the International Space Science Institute in Bern, Switzerland.

APPENDIX

In this Appendix, we re-derive for relativistic energies the governing acceleration equation of the pump mechanism of Fisk & Gloeckler, which is described in Section 4. The derivation is modeled after the derivation for nonrelativistic energies presented in Fisk & Gloeckler (2011a, 2011b).

We assume that the particle behavior can be described by the Parker transport equation. The Parker equation is appropriate for describing particle behavior in the mesoscale compressions and expansions required for the pump mechanism, which have large cross-sectional dimensions compared to the gyroradii of the particles that are being accelerated. The distribution function $f(\mathbf{r}, p, t)$ thus behaves as

$$\frac{\partial f}{\partial t} + \delta\mathbf{u} \bullet \nabla f = \frac{\nabla \bullet \delta\mathbf{u}}{3} p \frac{\partial f}{\partial p} + \nabla \bullet (\vec{\kappa} \bullet \nabla f), \qquad \text{(A1)}$$

where $\delta\mathbf{u}$ is the random convective velocity of the turbulence and $\vec{\kappa}$ is the spatial diffusion tensor. Note that an explicit assumption in the Parker transport equation is that f is quasi-isotropic, with only first-order anisotropies allowed.

In an acceleration mechanism based on the Parker equation there are certain restrictions on the source of energy to the tail particles, which are useful in the derivation. We integrate the Parker equation over all tail particle speeds and volume to form an equation for the behavior of the tail pressure

$$P = \frac{4\pi}{3} \int_{p_{\rm th}}^{\infty} v p^3 f \, dp \qquad \text{(A2)}$$

or

$$\int_{\rm vol} d\,({\rm vol}) \left[\frac{\partial P}{\partial t} + \delta\mathbf{u} \bullet \nabla P + 4\pi \frac{(\nabla \bullet \delta\mathbf{u})}{3} \int_{p_t}^{\infty} \gamma v p^3 f \, dp \right]$$
$$= - \int_{\rm vol} d\,({\rm vol}) \frac{\nabla \bullet \delta\mathbf{u}}{3} \frac{4\pi}{3} v p^4 f |_{p=p_{\rm th}}$$
$$+ \int_{\rm sur} d\mathbf{S} \bullet \frac{4\pi}{3} \int_{p_t}^{\infty} v p^3 \vec{\kappa} \bullet \nabla f \, dp. \qquad \text{(A3)}$$

Here,

$$\gamma = \frac{p}{v p^4} \frac{\partial}{\partial p} (v p^4), \qquad \text{(A4)}$$

which is 5 for nonrelativistic particles and 4 for relativistic particles.

The integrand on the left-hand side of Equation (A3) is the time rate of change of the average pressure in the volume plus the time rate of change of the pressure due to the change in the volume of a compression or expansion region. The volume of the entire system is constant. The changes in the pressure due to changes in the volume of a compression or expansion region must thus integrate to zero. The first expression on the right-hand side of Equation (A3) is the net flow of energy across the threshold boundary due to the compressions and expansions. The second term on the right of Equation (A4) is the flow of energy due to spatial diffusion across the outer boundary of the volume containing the compressions and expansions. In the case of the Galaxy, there is no net flow of energy from outside the Galaxy; there is only escape of energy due to spatial diffusion.

Equation (A3) thus requires that the only source of energy to the average pressure of the tail is the net flow of energy from the core due to the compressions and expansions. The loss of energy is due to escape from the volume of the system (Galaxy) due to spatial diffusion. In a steady state, the energy gain from the pump mechanism balances the energy loss from escape. Note that there is no damping of turbulence in this model. There is only redistribution of the energy in the core particles to form the tail.

The spatial diffusion term is difficult to deal with in the Parker equation. It contains divergences of spatial gradients. There is a useful approximation that we can apply to greatly simplify the calculations.

We first divide the distribution function into two terms, $f = f_o + \delta f$. Here, f_o is the portion of the distribution function that results from multiple compressions and expansions. The pump mechanism systematically pumps the tail up through a series of compressions and expansions, and f_o is a result of this history. A given location in space should have approximately the same time history as any other location in our volume, and thus f_o is not a strong function of position, nor does it depend upon the local value of $\nabla \bullet \delta \mathbf{u}$.

The deviation of f from f_o due to the local value of $\nabla \bullet \delta \mathbf{u}$ is δf, which clearly has strong spatial gradients that are responsible for particles spatially diffusing. We can then approximate the spatial diffusion term in the Parker equation as a gain or loss term, or $-\delta f/\tau$, where τ is the characteristic time for escape from a compression or diffusion into an expansion.

The timescale τ for escape from a compression or diffusion into an expansion region should be short compared with the characteristic escape time from the Galaxy. Thus, the Parker equation for particles in individual compressions and expansions is

$$\frac{\partial (f_o + \delta f)}{\partial t} + \delta \mathbf{u} \bullet \nabla (f_o + \delta f) + \frac{\gamma}{3} (\nabla \bullet \delta \mathbf{u}) (f_o + \delta f) = \frac{\nabla \bullet \delta \mathbf{u}}{3 v p^3} \frac{\partial}{\partial p} (v p^4 \delta f) + \frac{\nabla \bullet \delta \mathbf{u}}{3 v p^3} \frac{\partial}{\partial p} (v p^4 f_o) - \frac{\delta f}{\tau}. \quad \text{(A5)}$$

For the next step in the derivation, we need to relate δf to f_o. We saw in Equation (A3) that the only source of energy to the average tail is the core. Thus,

$$\frac{\partial f_o}{\partial t} = \frac{1}{3 v p^3} \frac{\partial}{\partial p} (v p^4 \langle \delta f \nabla \bullet \delta \mathbf{u} \rangle), \quad \text{(A6)}$$

where the angular brackets denote an average over multiple compressions and expansions. If we integrate Equation (A5) over particle momentum to find the time rate of change of the pressure associated with f_o, the only source of the pressure over multiple cycles is the second-order flow of energy from the core, as required.

We then subtract Equation (A6) from (A5), ignoring the difference between $\delta f \nabla \bullet \delta \mathbf{u}$ and its average (a standard quasi-linear approximation), or

$$\frac{\partial \delta f}{\partial t} + \delta \mathbf{u} \bullet \nabla \delta f + \frac{\gamma}{3} (\nabla \bullet \delta \mathbf{u}) (f_o + \delta f) = \frac{1}{3 v p^3} \frac{\partial}{\partial p} (v p^4 (\nabla \bullet \delta u) f_o) - \frac{\delta f}{\tau}, \quad \text{(A7)}$$

and integrate Equation (A7) to form an equation for the pressure P_o due to f_o and the pressure δP due to δf, or

$$\frac{\partial \delta P}{\partial t} + \delta \mathbf{u} \bullet \nabla \delta P + \frac{(\nabla \bullet \delta \mathbf{u})}{3} \frac{4\pi}{3} \int_{p_{\rm th}}^{\infty} \gamma v p^3 (f_o + \delta f)\, dp = \frac{4\pi}{3} \int_{p_{\rm th}}^{\infty} dp \left[\frac{1}{3} \frac{\partial}{\partial p} (v p^4 (\nabla \bullet \delta u) f_o) - v p^3 \frac{\delta f}{\tau} \right] = \frac{4\pi}{3} \left[\frac{1}{3} v p^4 (\nabla \bullet \delta \mathbf{u}) f_o \bigg|_{p_{\rm th}} - \int_{p_{\rm th}}^{\infty} v p^3 \frac{\delta f}{\tau} dp \right]. \quad \text{(A8)}$$

The left-hand side of Equation (A8) is the change in pressure of the tail particles as a result of a compression or expansion. It is the work done on the tail in a compression, or by the tail in an expansion. The right-hand side of Equation (A8) is the flow of energy from the core plus the flow of energy from surrounding regions by spatial diffusion.

We saw in Equation (A3) that if we average over multiple expansions and compressions the tail particles cannot do net work. This is possible only if there is no flow of energy into or out of an expansion/compression, i.e., that the right-hand side of Equation (A8) is zero. A flow of energy into or out of an expansion/compression will alter the pressure, the pressure gradients, and the work done by the expansion/compression. These are unacceptable second-order correlations that do not average to zero over multiple compressions and expansions.

The requirement that the flow of energy into or out of an expansion/compression, averaged over multiple expansions/compressions, cannot result in work being done is consistent with the second law of thermodynamics. We have a cyclic pump mechanism with the purpose of extracting energy from the core to create the tail. According to the second law of thermodynamics, it is impossible to devise a system which, working in a cycle, shall produce no effect other than the extraction of heat from a reservoir and the performance of an equal amount of mechanical work.

For the right-hand side of Equation (A8) to be zero in general, we require that

$$\frac{\nabla \bullet \delta \mathbf{u}}{3} \frac{\partial}{\partial p} (v p^4 f_o) = \frac{v p^3 \delta f}{\tau}. \quad \text{(A9)}$$

Finally, substituting Equation (A9) into (A6) yields the final equation for the time evolution of f_o:

$$\frac{\partial f_o}{\partial t} = \frac{1}{v p^3} \frac{\partial}{\partial p} \left(\frac{\langle (\nabla \bullet \delta \mathbf{u})^2 \tau \rangle}{9} p \frac{\partial}{\partial p} (v p^4 f_o) \right). \quad \text{(A10)}$$

It is convenient to express τ in terms of a spatial diffusion coefficient, κ, or, $\langle (\nabla \bullet \delta \mathbf{u})^2 \tau \rangle \equiv (\delta u^2 \lambda^2)/(\lambda^2 \kappa) \equiv \delta u^2/\kappa$. Here, λ is a characteristic size of a compression or expansion region, in a direction normal to the mean magnetic field. Note that κ is the cross-field diffusion coefficient.

Finally, if we balance the acceleration in individual compressions and expansions against escape from the Galaxy, we derive Equation (4), the governing equation for the acceleration of GCRs in the interstellar medium:

$$\frac{1}{v p^3} \frac{\partial}{\partial p} \left(\frac{\delta u^2}{9 \kappa_l} p \frac{\partial}{\partial p} (v p^4 f_o) \right) = \frac{4 \kappa_g}{R_g^2} f. \quad \text{(4)}$$

It should be noted that Jokipii & Lee (2010) argue that density is not properly conserved in Equation (A10). They integrate (A10) over all tail particle momentum to find an equation for the time rate of change of the density n_o associated with f_o. If we use the form of Equation (A10) in Equation (A6), this integration yields

$$\frac{dn_o}{dt} = -\frac{4\pi p^3}{3} \langle \delta f \nabla \bullet \delta \mathbf{u} \rangle |_{p_{\rm th}} + \frac{4\pi}{3} \int_{p_{\rm th}}^{\infty} \frac{p^2}{v} \left(\frac{\partial}{\partial p} (vp) \right) dp \, \langle \delta f \nabla \bullet \delta \mathbf{u} \rangle. \quad \text{(A11)}$$

The first term on the right-hand side of Equation (A11) represents a flow of particles across the threshold boundary from the core. Jokipii & Lee (2010) state that the second term on the right-hand side is a spurious source term that appears to be creating particles.

The actual requirement for the conservation of density is not that the density n_o is conserved, but rather that the total density, $n_o + \delta n$, is conserved, where δn is the density associated with δf. We showed in Equations (A7) and (A8) that in order for the tail particles to do no net work when averaged over multiple compressions and expansions (as is required by the Parker equation),

$$\frac{\partial \delta f}{\partial t} + \delta \mathbf{u} \bullet \nabla \delta f + \frac{\gamma}{3} (\nabla \bullet \delta \mathbf{u}) (f_o + \delta f) = 0, \tag{A12}$$

where γ is given in Equation (A4). If we convert Equation (A12) into an equation for the density, and average over the volume, we find that

$$\frac{d\delta n}{dt} = -\frac{4\pi}{3} \int_{p_{\rm th}}^{\infty} \frac{p^2}{v} \left(\frac{\partial}{\partial p} (vp) \right) dp \, \langle \delta f \nabla \bullet \delta \mathbf{u} \rangle. \tag{A13}$$

Thus, adding Equations (A11) and (A12), we see that the only source of density to the tail is the core, and density is properly accounted for in the pump mechanism.

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