Magnetic Pumping as a Source of Particle Heating and Power-law Distributions in the Solar Wind

E. Lichko¹, J. Egedal¹, W. Daughton², and J. Kasper³

¹Department of Physics, University of Wisconsin–Madison, Madison, WI 53706, USA
²Los Alamos National Laboratory, Los Alamos, NM 87545, USA
³University of Michigan, Ann Arbor, MI 48109, USA

Received 2017 August 11; revised 2017 November 10; accepted 2017 November 12; published 2017 November 27

Abstract

Based on the rate of expansion of the solar wind, the plasma should cool rapidly as a function of distance to the Sun. Observations show this is not the case. In this work, a magnetic pumping model is developed as a possible explanation for the heating and the generation of power-law distribution functions observed in the solar wind plasma. Most previous studies in this area focus on the role that the dissipation of turbulent energy on microscopic kinetic scales plays in the overall generation of the plasma. However, with magnetic pumping, particles are energized by the largest-scale turbulent fluctuations, thus bypassing the energy cascade. In contrast to other models, we include the pressure anisotropy term, providing a channel for the large-scale fluctuations to heat the plasma directly. A complete set of coupled differential equations describing the evolution, and energization, of the distribution function are derived, as well as an approximate closed-form solution. Numerical simulations using the VPIC kinetic code are applied to verify the model’s analytical predictions. The results of the model for realistic solar wind scenario are computed, where thermal streaming of particles are important for generating a phase shift between the magnetic perturbations and the pressure anisotropy. In turn, averaged over a pump cycle, the phase shift permits mechanical work to be converted directly to heat in the plasma. The results of this scenario show that magnetic pumping may account for a significant portion of the solar wind energization.

Key words: acceleration of particles – magnetic fields – plasmas – scattering – solar wind

1. Introduction

Spacecraft measurements of the solar wind indicate the presence of an anomalous source of heating throughout the heliosphere. The radial temperature distribution observed not only exceeds the values expected given the rapid expansion of the solar wind within the Parker spiral (Kasper et al. 2002; Richardson & Smith 2003; Cranmer et al. 2009; Stverak et al. 2009), but also the heating that would be expected from steady-state hydrodynamic models (Parker 1965; Hartle & Sturrock 1968; Dursey 1972). Much of the work on this topic has focused on how energy is injected as turbulent fluctuations at large spatial scales, and then propagates through the turbulent cascade and is absorbed by the plasma at the smaller dissipation scale (Howes et al. 2008, 2011; Chandran et al. 2010; Bruno & Carbone 2013; Told et al. 2015). While there exists direct evidence for the turbulent cascade operating in the solar wind (Sahraoui et al. 2009), there are details of the heating problem not readily explained by either the energy cascade or traditional stochastic acceleration mechanisms. One such detail in particular is the observed power law in the solar wind’s particle distributions, $f \propto \nu^{-\gamma}$ where $\gamma = -5$, which is present throughout the solar wind (Fisk & Gloeckler 2006, 2012).

In this Letter, we propose a heating mechanism based on magnetic pumping, a process by which a series of either periodic or random magnetic perturbations heats a plasma. While it in no way precludes energy transfer through the turbulent cascade, this model allows energy transfer to the particles directly from plasma perturbations on the largest scales, bypassing the turbulent cascade. It also accounts for many of the observations seen in the solar wind, such as the power-law distributions in velocity space.

The most important physical aspect of the model, and the aspect in which it differs most from other models, is the role of pressure anisotropy in relation to the fluctuating magnetic field. To elucidate this mechanism, one may consider a simple magnetic flux tube. In the absence of heat fluxes and other isotropizing effects, contractions of the tube lead to pressure anisotropy in phase with the density and magnetic perturbations (Chew et al. 1956), such that there is no net energization. If, however, an isotropizing effect is present there will be a net energization of the plasma, directly related to a finite phase difference between the pressure anisotropy and the plasma compressions, yielding positive work by the term $P_\perp \nabla_\perp \cdot \nu$ when averaged over a pump cycle.

For the one-dimensional flux tube described above, the spatially isotropizing effects can be introduced through pitch-angle scattering. However, in a two-dimensional scenario, where the plasma perturbations include temporal as well as spatial variations, pressure anisotropy is reduced by parallel thermal heat fluxes of particles streaming into and out of the region. Thus, anisotropy in $f(\nu)$ with spatial variations of $I_{\perp}$ decays at the rate $\nu_{\text{eff}} \sim I_{\perp}/\nu$. Below, when applying the model to the solar wind for realistic $dB/B(\omega)$ spectra, we find that $\nu_{\text{eff}} \sim I_{\perp}/\nu$ dominates the ion isotropization, permitting significant ion energization and the associated formation of power-law distributions.

2. Comparison to Other Models

Many models address the heating and formation of power-law distributions in the solar wind (Fisk & Gloeckler 2006, 2012; Chandran et al. 2010), but do not adequately explain the observed levels of energization (Lynn et al. 2013). In contrast to that body of work, the present analysis considers heating
channeled to the plasma by pressure anisotropy. We note that the mathematical framework is similar to that applied by Drake et al. (2013) for Fermi acceleration during magnetic island coalescence, where magnetic or density perturbations yield velocity changes \( \Delta v \propto v \). The analysis then brings about a velocity diffusion equation of the form

\[
\frac{df}{dt} = K(\tau_{pump}, \nu_{eff}) \frac{1}{v^2} \frac{d}{dv} \left( v^2 D \frac{df}{dv} \right),
\]

where \( D \propto (\Delta v)^2/\tau_{pump} \propto v^2/\tau_{pump} \) is the diffusion coefficient and \( K \) is a function with a maximum at \( 1/\tau_{pump} \propto \nu_{eff} \), providing a measure of the level of anisotropy displayed by \( f \).

While the model has similarities with the form of transit time heating studied by Berger et al. (1958), it should not be confused with the transit time damping of linear propagating waves (Stix 1992). In our model the heating efficiency is derived by considering a standing wave, for which wave-particle resonances are unimportant. Thus, the resultant heating does not involve a specific resonant velocity, as in Landau or transit-time damping where only particles with \( v_0 \approx \omega/K_l \) are involved in the heating. With the standing wave being comprised of two oppositely propagating compression waves, the heating can be interpreted as a nonlinear interaction between these two waves. While such interactions are fundamental to the development of plasma turbulence (Chandran & Hollweg 2013), the described heating channel is typically neglected with the assumption of isotropic velocity distributions. It is also worth noting that magnetic pumping has been investigated in the context of fusion research where Coulomb collisions provide pitch-angle scattering, but also relax the heated distributions into Maxwellians (Laroussi & Roth 1989; Borovoy & Hansen 1990). However, in the solar wind, particles may pitch-angle scatter collisionless off of instabilities and fluctuations (Bale et al. 2009; Verscharen et al. 2016). This occurs with negligible energy diffusion (Kulsrud & Pearce 1968) such that Boltzmann’s H-theorem does not apply (Boltzmann 1872), and Maxwellian distributed particles are generally not observed. We show below that in this limit of collisionless scattering, magnetic pumping yields power-law solutions with spectral indices consistent with values observed in situ in the solar wind.

3. Kinetic Simulations

Our initial setup is a one-dimensional flux tube, as shown in Figure 1(a). The domain is doubly periodic and an external, periodically driven current is applied along two infinite current sheets each located halfway between the mid-line and the top and bottom edge of the simulation space. The oppositely directed current sheets cause flux tube expansions and contractions as the current oscillates (as shown in Figure 1(b)). The background distribution is given by a Maxwellian with uniform temperature, \( T = T_0 = T_0 \), with the mass ratio given by \( m_i/m_e = 100 \). The simulations used a non-relativistic thermal speed \( v_{thc}/c = 0.0707 \) and \( \omega_p/\Omega_{ci} = 1 \). Spatial scales are normalized by \( d^\star \) and our pumping frequency, \( \nu_{pump} \), is normalized by \( \omega_p \). In the simulations presented below, we use \( \omega = \omega_{pump} = 0.1 \omega_p \), where \( \omega_{pump} \) is referred to hereafter as \( \omega \). Our density and magnetic field fluctuations are normalized by the initial density, \( n_0 = 1 \), and background magnetic field, \( B_0 = B_0 \delta = 1 \), respectively. For the purposes of this initial analysis, the scattering frequency, \( \nu \), is velocity-independent and is implemented using the Takizuka and Abe Monte Carlo method employed to calculate Coulomb collisions in VPIC (Takizuka & Abe 1977; Daughton et al. 2009). Only electron–ion collisions are included, so the energy diffusion is minimal given the mass ratio. The 1D flux tube simulations were carried out in VPIC for a variety of scattering frequencies. From Figure 1(d) it is clear that magnetic pumping is increasing the temperature of the plasma. Furthermore, both the energization and the phase difference between pressure anisotropy and magnetic field show a dependence on scattering frequency, \( \nu \).

4. Analytical Model Derivation

We next proceed to derive an analytic model to explain the energization and demonstrate that it matches the results from the kinetic simulations. To that end, we consider a periodic flux tube with length \( l \) and radius \( r \). The plasma within the flux tube is assumed to remain uniform, while \( r \) and \( l \) change slowly in time such that the magnetic moment \( \mu = m v^2 / (2B) \) of the particles is conserved. Furthermore, given the periodic boundary conditions, the action integral \( I = \oint \nu \, dl \) is also an adiabatic invariant provided that the length of the tube is not changed significantly during a typical particle transit.

For this system, our first aim is to obtain a reduced drift kinetic equation, \( df/dt = 0 \), where

\[
\frac{d}{dt} = \frac{\partial}{\partial t} + \frac{\partial v}{\partial t} \frac{\partial}{\partial v} + \frac{\partial v}{\partial t} \frac{\partial}{\partial v},
\]

is the total time derivative along the particle trajectories. Note that given the assumption of a uniform plasma the convective spatial derivative term \( (v \cdot \nabla) \) vanishes.

Assuming \( \mu \propto v^2 / B \), \( J \propto l v_0 \), and \( N = \pi r^2 n_l \), we use the conservation of magnetic moment, action, and particle number to rewrite the above time derivative. For the simple geometry

![Figure 1](image-url)
considered, the drift kinetic equation $df/ dt = 0$ simplifies to

$$\frac{\partial f}{\partial t} + \frac{B}{B} v^3 \frac{\partial f}{\partial v} - \left( \frac{n - \frac{B}{B}}{n} \right) v \frac{\partial f}{\partial v} = 0. \quad (2)$$

In this limit without scattering, we note that evolution equations for $p$ and $p_\perp$ are readily derived by calculating the $v^2$ and $v_\perp^2$ moments of of Equation (2), which yield the CGL double adiabatic scaling laws

$$p \propto n^3/B^2, \quad p_\perp \propto nB. \quad (3)$$

In the following, we explore changes induced in $f$ due to uniform perturbations of the flux tube in conjunction with steady pitch-angle diffusion limiting the development of pressure anisotropy. Thus, similar to Drake et al. (2013), we generalize our kinetic equation to include additional physical effects

$$\frac{df}{dt} = \nu LF - c_1 f + c_2 f_{ext}, \quad (4)$$

where $L = \partial/\partial \zeta (1 - \zeta^2) \partial/\partial \zeta$ is the Lorentz scattering operator, $\zeta = v_i/v$ is the cosine of the pitch angle, and $\nu$ is a typical frequency for the scattering processes. The constants $c_1$ and $c_2$ specify the rate of plasma losses and rate of incoming (external $f_{ext}$) plasma, respectively.

For the analysis below it is convenient to change variables from $(v_i, v_\perp)$ to $(v, \zeta)$. Equation (4) then takes the form

$$\frac{df}{dt} + R \left( P_2(\zeta) \frac{\partial f}{\partial \zeta} + \frac{3}{2} (1 - \zeta^2) \frac{\partial f}{\partial \zeta} \right) + \frac{n}{3n} \frac{\partial f}{\partial v} = \nu LF - c_1 f + c_2 f_{ext}, \quad (5)$$

where $P_2(\zeta)$ is the second-order Legendre polynomial and $R = (2/3)n - B/B$. We note that $d/dt \log(p_i/p_\perp) = d/dt \log(n^2/B^2) = 3R$, showing that $R^3$ is proportional to the rate at which the pressure anisotropy builds in the CGL system (the system with $\nu = c_1 = c_2 = 0$).

To evaluate the efficiency by which the plasma is energized in the above framework, we next consider periodic perturbations for the magnetic field and density. An approximate solution to Equation (5) can be obtained by expanding $f$ in a series of Legendre polynomials $f(v, \zeta, t) = \sum \bar{f}_j(\zeta) f_j(v, t)$ where $P_j$ is the $j$-th order Legendre polynomial. The approach provides a set of coupled differential equations, which we solve numerically, and a first-order approximation to the results. These two solutions will then be compared to the results of the kinetic simulations.

We still consider the uniform and periodic flux tube but now with imposed sinusoidal temporal variations in density and magnetic field:

$$\frac{\dot{n}}{n} = \frac{\delta n}{n} = i \omega \delta, \quad \frac{\dot{B}}{B} = \frac{\delta B}{B} = i \omega \delta \phi, \quad R = \delta \omega \delta R \delta, \quad \delta R = \frac{2}{3} \frac{\delta n}{n} - \frac{\delta B}{B} e^{i \phi}, \quad (6)$$

Our aim is again to obtain a solution to Equation (5). Given the periodic variations of the drive, contrary to the analysis in Drake et al. (2013), we do not need to impose an ordering involving $\nu$, but only require that $\delta R \ll 1$. By inserting the above expansion in pitch angle into Equation (5) and integrating over $\int P_2(\zeta)d\zeta$, we can obtain a set of coupled differential equations for an arbitrary order of Legendre polynomial. These equations can be solved numerically to arbitrary precision. In the following comparisons, all numerical solutions were taken to second order, truncating at the $f_2$ equation, as further terms resulted in negligible improvements in accuracy.

Additionally, we obtain an approximate solution by assuming that each $f_n$ is comprised of a slowly varying component and a rapidly varying component, denoted hereafter as

$$f_n = f_n^0(v, t) + f_n^4(v) e^{i \nu t}. \quad (7)$$

Inserting this approximation into Equation (5) and collecting terms proportional to $P_2(\zeta) e^{i \nu t}$, we obtain the relation

$$f_2^0 = K \nu \frac{\partial f_0^0}{\partial \nu}, \quad K = -\frac{\omega}{\nu} \frac{\delta R}{\omega^2 + 36 \nu^2} \quad \nu \neq 0. \quad (8)$$

Equation (8) shows how the $P_2$-perturbation of the distribution develops and will, for finite $\nu$, be offset in phase from the drive oscillation in $R$, by the angle $\theta = \arctan(6 \nu/\omega)$. This phase shift is important when solving for $f_0$ we obtain non-vanishing time averages from the terms involving $f_2$. Since $E = 5/2 \int \nu f_0(v) d\nu$, these non-vanishing terms become the source of the energization. Using Equation (8), an equation is obtained for the slowly varying “background” distribution

$$\frac{\partial f_0^s}{\partial t} = -\frac{3}{5} \frac{\nu (\omega \delta R)^2}{\omega^2 + 36 \nu^2} \frac{1}{\nu^2} \frac{\partial^2 f_0^s}{\partial \nu^2} = -c_1 f_0 + c_2 f_{ext}. \quad (9)$$

For velocities sufficiently large that the cold source is negligible, the solutions to Equation (9) then take the form

$$f_0^s \propto v^\gamma, \quad \gamma = -\frac{3}{2} - \sqrt{\frac{9}{4} + \frac{c_1}{G}}, \quad G = \frac{3}{5} - \frac{\nu (\omega \delta R)^2}{\omega^2 + 36 \nu^2} \quad (10)$$

In the limit of no net losses (i.e., $c_1 = 0$), the exponent, $\gamma$, approaches $-3$. Thus, the heating mechanism is more than adequate to account for the observations of $f \propto v^{-3}$ distributions typically observed in the solar wind (Fisk & Gloeckler 2006, 2012).

From Equation (9) we can directly obtain the heating range

$$\frac{\partial E}{\partial t} = \frac{3}{2} G \int_0^\infty v^2 \frac{1}{v^2} \frac{\partial}{\partial v} \left( \frac{v^4 \partial f_0^s}{\partial v} \right) 4\pi v^2 d\nu = 10 G E, \quad (11)$$

which gives us

$$E = \frac{6 \nu (\omega \delta R)^2}{e^{\omega \delta R} - 1}. \quad (12)$$

When we combine the above expression with the analytic solution for the $\nu = 0$ case, $E_0(t)$, which is essentially sinusoidal, to obtain a solution for the energy system for arbitrary $\nu$ that agrees with the simulation

$$\frac{E(t)}{E_0(t)} = \frac{e^{\nu \delta \omega \delta R^2}}{e^{\nu \delta \omega \delta R^2} - 1}. \quad (13)$$

5. Discussion and Application to the Solar Wind

Given the above results, we next compare the predictions of our analytic model with the results from the kinetic simulations. The relationship between the relative energy evolution ($E/E(t = 0))$ and the scattering frequency is shown in Figure 2(a). Again, there is good agreement between the VPI simulations and both the exact numerical results of our analytic...
model as well as the results from the first-order approximation. Based on the form of Equation (13), the scattering frequency that will maximize the energization is obtained, as shown in Figure 2(a). The analytic solutions and VPIC results all peak at this most efficient frequency, further lending credence to the agreement between the models. Similarly, for the phase difference between \( P_e/P_i \) and \( B \) there is a good agreement between the two analytic solutions and the VPIC simulations, as in Figure 2(b).

For the solar wind, scattering is infrequent and the main isotropizing effect for the pressure anisotropy is thermal streaming, for which we estimate \( \nu_{\text{eff}} \sim \nu_{\text{pert}}/\nu \). To verify this we set up a VPIC simulation using the same domain as the 1D simulations described above, but with no magnetic fluctuations. We initialized the domain with a spatially dependent heating rate. The results of this for both electrons and ions are shown in Figure 3(d), where the decay rate matches our expectations of \( \nu_{\text{eff}} \sim \nu_{\text{pert}}/\nu \). A similar phenomenon has been used in other models, such as in the fluid closure by Hammett & Perkins (1990). The isotropization caused by thermal streaming is much larger than that induced by pitch-angle scattering off waves and Coulomb collisions. To estimate \( \nu_{\text{eff}} \) for the solar wind we need to take into account the spatial anisotropy of the fluctuations, i.e., that \( k_\perp \gg k_\parallel \) (Chen 2016). Because the particle streaming is restricted to be along magnetic field lines, only the field-aligned parts of the perturbations are important, such that \( \nu_{\text{eff}} \sim \nu _{\text{pert}}/k_\parallel 2\pi /v_{\text{pert}}(k_\parallel) \). Because the magnetic moment is conserved during streaming, the energization will now be in the parallel direction. The parallel heating can be transferred to the perpendicular directions through standard scattering mechanisms (Achterburg 1981; Lynn et al. 2013). However, in cases where scattering is sufficiently slow, the large scale and slowly varying background distribution may develop a significant anisotropy \( f_{\parallel} \gg f_{\perp} \). To be investigated elsewhere, this will then introduce additional terms in Equation (9) that will help increase the effectiveness of the heating process beyond the examples considered here.

The role of parallel streaming can be tested directly in 2D kinetic simulations, but given the numerical cost only for a limited number of pump cycles. As shown in Figure 3, the setup is similar to that in Figure 1(a), with the domain extended in the \( x \) direction, and the current sheets providing the oscillating magnetic perturbations only cover a portion of the simulation domain. In this case, the heating is no longer dependent on the scattering frequency, \( \nu \), as it was in the 1D simulations shown in Figure 3(b). The thermal streaming of electrons in and out of the pumping region acts as an effective scattering process, dominating the applied scattering rates and leading to the expected increase in heating for low values of \( \nu \).

To address heating over hundreds of pump cycles, we evolve Equation (5) using parameters relevant to ions in the solar wind. To generate the results in Figure 4 a selection of frequencies were randomly generated and their corresponding \( dB/B(\omega) \) were taken from the spectra in Leamon et al. (1998). After every cycle a new randomly generated frequency and corresponding \( dB/B(\omega) \) was chosen, so that every cycle the plasma would experience a new frequency, consistent with the observation that fluctuations in MHD turbulence decohere after a single cycle. From Chen (2016, Figure 7) we used \( k_\parallel/k_\perp \sim 8 \)
when generating Figure 4. Assuming a solar wind speed of 800 km s$^{-1}$, we obtain an estimate of the total transit time from the Sun to the Earth. We further assume that only 10% of the perturbations that the plasma experiences will be compressional. As shown in Figure 4, we obtain a power-law distribution out to two orders of magnitude in velocity. In the model, the streaming rapidly dissipates the anisotropy at higher velocities, causing the drop-off in the power law in Figure 4 for $v \gg v_{\text{wind}}(k_i/k_j)$. While we observe a power law extending about four orders of magnitude in $E$, work in progress suggests that the streaming can be reduced by trapping effects, extending the power-law part of the distributions to even higher velocities, $v \gg v_{\text{wind}}(k_i/k_j)$.

6. Conclusions

In summary, we have explored magnetic pumping as a possible heating mechanism of the solar wind. Energy associated with large-scale magnetic and density fluctuations heats the plasma directly. The energization is related to the phase of the pressure anisotropy being shifted from the turbulent drive through the inclusion of collisionless pitch-angle mixing. This phase difference between the pressure anisotropy and the driving magnetic perturbations serves as the source of the heating, and an important point of distinction between this model and previous models of energization in the solar wind. Bypassing the turbulent cascade, the model provides an efficient scheme for dissipating energy directly with the largest-scale fluctuations and generating power-law particle distributions extending several orders of magnitude in energy.

This work was supported by NASA grant No. NNX15AJ73G. Contributions from E.L. were conducted with Government support under and awarded by the DoD, Air Force Office of Scientific Research, National Defense Science and Engineering Graduate (NDSEG) Fellowship, 32CFR168a. Simulations were performed with LANL institutional computing resources. We thank Dr. K Klein and the referee for comments and information that helped improve the manuscript.

ORCID iDs

W. Daughton @ https://orcid.org/0000-0003-1051-7559
J. Kasper @ https://orcid.org/0000-0002-7077-930X

References

Achterburg, A. 1981, A&A, 97, 259
Bale, S. D., Kasper, J. C., Howes, G. G., et al. 2009, PhRvL, 103, 211101
Berger, J., Newcomb, W., Dawson, J., et al. 1958, PhFl, 1, 301
Boltzmann, L. 1872, Sitzungsberichte Akad. Wiss., 66, 275, https://babel.hathitrust.org/cgi/pt?id=umn.31951p11139293w;view=1up;seq=301
Borovsky, J. E., & Hansen, P. J. 1990, PhFl, 2, 1114
Bruno, R., & Carbone, V. 2013, LRSP, 20, 2
Chandran, B., Li, B., Rogers, B., Quataert, E., & Germaschewski, K. 2010, ApJ, 720, 503
Chandran, B. D. G., & Hollweg, J. V. 2013, ApJ, 707, 1659
Chen, C. H. K. 2016, JPPb, 82, 535820602
Chew, G. F., Goldberger, M. L., & Low, F. E. 1956, RSPSA, 236, 112
Cranmer, S. R., Matthaeus, W. H., Breech, B. A., & Kasper, J. C. 2009, ApJ, 702, 1604
Daughton, W., Roytershteyn, V., Albright, B., et al. 2009, PhFl, 16, 072117
Drake, J. F., Swisdak, M., & Fermo, R. 2013, ApJL, 763, L5
Durney, B. R. 1972, JGR, 77, 4042
Fisk, L. A., & Gloeckler, G. 2006, ApJL, 640, L79
Fisk, L. A., & Gloeckler, G. 2012, SSRv, 173, 433
Hammett, G. W., & Perkins, F. W. 1990, PhRvL, 64, 3019
Hartle, P. A., & Sturrock, R. E. 1968, ApJ, 151, 1155
Howes, G. G., Dorland, W., Cowley, S. C., et al. 2008, PhRvL, 100, 065004
Howes, G. G., TenBarge, J. M., Dorland, W., et al. 2011, PhRvL, 107, 035004
Kasper, J., Lazarus, A., & Gary, S. 2002, GeoRL, 29, 20
Kulsrud, R., & Pearce, W. P. 1968, ApJ, 156, 445
Laroussi, M., & Roth, J. R. 1989, PhFl, 1, 1034
Leamon, R. J., Smith, C. W., Ness, N. F., Matthaeus, W. H., & Wong, H. K. 1998, JGR, 103, 4775
Lynn, J., Quataert, E., Chandran, B., & Parrish, I. 2013, ApJ, 777, 128
Parker, E. 1965, SSRv, 4, 666
Richardson, J., & Smith, C. 2003, GeoRL, 30, 1206
Sahraoui, F., Goldstein, M. L., Robert, P., & Khotyaintsev, Y. V. 2009, PhRvL, 102, 231102
Stix, T. H. 1992, Waves in Plasmas (New York: AIP)
Stverak, S., Maksimovic, M., Travnicek, P. M., et al. 2009, JGR, 114
Takizuka, T., & Abe, H. 1977, JCoPh, 25, 205
Told, D., Jenko, F., TenBarge, J., Howes, G., & Hammett, G. 2015, PhRvL, 115, 025003
Verscharen, D., Chandran, B., Klein, K., & Quataert, E. 2016, ApJ, 831, 128
Erratum: “Magnetic Pumping as a Source of Particle Heating and Power-law Distributions in the Solar Wind” (2017, ApJL, 850, L28)

E. Lichko¹, J. Egedal¹, W. Daughton², and J. Kasper³

¹Department of Physics, University of Wisconsin–Madison, Madison, WI 53706, USA
²Los Alamos National Laboratory, Los Alamos, NM 87545, USA
³University of Michigan, Ann Arbor, MI 48109, USA

Received 2018 January 23; published 2018 February 13

In Figure 4 of the published article (Lichko et al. 2017), there was an error in the code that started the distribution function evolution with a slightly different initial condition than the one in red shown in the corrected Figure 4. This error, combined with an excessive amount of smoothing, caused the overall density of the distribution function to decline significantly. We have fixed these errors in the corrected version of Figure 4 included here.

These errors do not alter the published article’s conclusion that magnetic pumping is an efficient mechanism for heating ions in the solar wind. The updated figure still shows an evolution toward a power-law distribution, $f \propto v^\gamma$, $\gamma = -3$ within the same conservative estimate of the transit time of a particle from the Sun to the Earth. Additionally, aside from the initial condition and the smoothing, all other aspects of the computation remained the same. Not only does the figure still demonstrate that it approaches the expected power law in roughly the same amount of time, but it also documents that the method works for changing fluctuations, not just coherent waves.

This work was supported by NASA grant No. NNX15AJ73G. Contributions from E.L. were conducted with Government support under and awarded by the DoD, Air Force Office of Scientific Research, National Defense Science and Engineering Graduate (NDSEG) Fellowship, 32CFR168a. Simulations were performed with LANL institutional computing resources. We thank Dr. K. Klein for comments and information that helped improve the manuscript.

Figure 4. Numerical solution for the distribution function after many oscillations. Note that the slope approaches the value commonly observed in the solar wind, $\gamma = -3$. In this plot, $v_{th} = 90 \text{ km s}^{-1}$.

ORCID iDs

W. Daughton © https://orcid.org/0000-0003-1051-7559
J. Kasper © https://orcid.org/0000-0002-7077-930X

Reference

Lichko, E., Egedal, J., Daughton, W., & Kasper, J. 2017, ApJL, 850, L28