Comment on “Electron demagnetization and heating in quasi-perpendicular shocks” by Mozer and Sundkvist

Steven J. Schwartz

1 Blackett Laboratory, Imperial College London, London, UK

Abstract In their paper, Mozer and Sundkvist (2013) present important observations of large-amplitude high-frequency electric field fluctuations in the vicinity of the ramp region of the Earth’s bow shock. In common with other recent work, they emphasize the role of electron scattering by such fluctuations in the problem of electron heating at collisionless shocks, relegating the DC fields to providing a reservoir of energy which can be tapped apparently to whatever degree is required. While the historical approach of attributing the zeroth-order electron heating to adiabatic motion in those DC fields has always required some unspecified scattering process, key elements of those DC processes have been lost or misinterpreted. In particular, particle motion parallel and perpendicular to the magnetic field are coupled to one another. More importantly, an \( O(1/2) \) fraction of the electron population traverses the shock from downstream to upstream or are trapped downstream. Thus, the concept of irreversible heating proceeding from the solar wind through the shock into the magnetosheath, implicit in Mozer and Sundkvist (2013), needs to be reconsidered. A highly simplified calculation illustrates these points.

1. Introduction

Observations of electrons at the Earth’s collisionless bow shock reveal several key characteristics, including inflated “flat-topped” distributions [Montgomery et al., 1970; Feldman et al., 1983] and rapid nearly isotropic “heating” through the steepest ramp region of the shock [Schwartz et al., 2011]. Investigations by Scudder et al. [1986] showed that the shape of the distribution closely followed that expected from Liouville’s theorem based on adiabatic electron motion in the shock DC electric and magnetic fields. That work, building on the suggested role of the cross-shock potential by Feldman et al. [1983], was further advanced by studies of the full pitch angle distributions through the shock performed by Lefebvre et al. [2007] and others.

It has always been recognized that adiabatic electron behavior could never account for all aspects of electron heating at shocks. Scattering by waves or turbulence is necessary to fill voids in phase space (see below for an illustration), to reduce the overall height of the distribution below the solar wind peak values, and to otherwise contribute to the electron isotropization.

Reports of small-scale, large-amplitude spikes or fluctuations in the electric (and magnetic) field [e.g., Walker et al., 2004; Bale and Mozer, 2007; Wilson, 2010], together with reports of ramp scales sufficiently short to demagnetize the electrons [Bale et al., 2003; Balikhin et al., 1993; Schwartz et al., 2011] have perhaps over-heightened the role of small-scale fluctuations and, in the process, oversimplified the features of the DC behavior.

This comment is designed to rebalance this discussion. In particular, it addresses the following points found in Mozer and Sundkvist [2013]:

1. The incorrect assertion that electron magnetic moment conservation leads to constancy of \( T_\parallel / B \) where \( T_\parallel \) is the electron kinetic temperature perpendicular to the magnetic field of magnitude \( B \).
2. The corresponding, although quantitatively less critical, assertion that under adiabatic assumptions the DC electric field heats purely the parallel electron temperature.
3. The notion that the cross-shock potential provides a “reservoir” of energy and that electron scattering results or can result in electrons acquiring only a fraction of that energy.
4. The implicit notion that the entire electron population starts in the solar wind, proceeds through the shock ramp where it is heated (i.e., its entropy is increased) to finish in the magnetosheath.
Mozer and Sundkvist [2013] also employ a Maxwellian fit to the core of the electron distribution in the magnetosheath in order to determine the downstream electron temperature. By their own calculation this underestimates the total electron temperature by a factor as large as 2. While they argue that the cadence required for full distributions, from which full moments could be calculated, would be much longer, it would be perfectly straightforwards to calculate moments based on the 1-D distributions they have available (Schwartz et al. [2011] call these pseudotemperatures). This point will not be developed further here, other than to note that increasing electron temperatures by a factor \( \sim 2 \) would make the discrepancy in Mozer and Sundkvist [2013, Figure 6] much less pronounced. It is interesting to compare this figure with Schwartz et al. [2011, Figure 3].

2. Adiabatic Electron Behavior in DC Electric and Magnetic Fields

It is illustrative to construct a simple, analytical model for the behavior of electron distributions in the presence of a layer across which the magnetic field increases and which includes an electrostatic potential jump. In terms of the shock application, the calculations are carried out in the de Hoffmann-Teller (dHT) frame [de Hoffmann and Teller, 1950] in which the bulk flow is aligned with the magnetic field and the shock layer is at rest. In this frame, the motional \( -\mathbf{v} \times \mathbf{B} \) electric field is zero. In the Normal Incidence Frame (NIF) employed by Mozer and Sundkvist [2013], electrons can drift along the shock front and lose energy against that motional field. A consequence of this, consistent with the right-handed rotation of the magnetic field out of the coplanarity plane [Thomsen et al., 1987] is that the dHT cross-shock potential is significantly smaller than that in the NIF [Goodrich and Scudder, 1984].

Denoting the two sides of this layer by “\( u \)” and “\( d \)” for upstream and downstream, adiabatic behavior conserves the magnetic moment and total energy of a particle, i.e.,

\[
\frac{mv^2_{d\parallel}}{2B_d} = \frac{mv^2_{u\parallel}}{2B_u}
\]

and

\[
\frac{mv^2_d}{2} = \frac{mv^2_u}{2} + e\phi
\]

where \( \mathbf{v} = v_\parallel \hat{\mathbf{B}} + v_\perp \) is the particle velocity and \( \phi \) the dHT electrostatic potential jump across the layer.

For illustrative purposes, take the upstream distribution to be a spherical “water bag”, i.e., one whose phase space distribution is constant within a sphere of radius normalized to unity and zero outside. The corresponding normalized potential is \( \Phi \equiv 2e\phi/m \). The only other parameter that enters is the field compression ratio \( \alpha \equiv B_d/B_u \). The model that follows is not intended to be a realistic model of a shock, but has been constructed specifically to isolate and illustrate aspects of some key physical processes.

The upstream distribution is shown in Figure 1 (top). While this is a long way from a real solar wind distribution, there is an important feature which is indicated in the figure, and a consequence of the fact that the electron thermal speed is much larger than the solar wind bulk flow. A large fraction (half in this case since the bulk flow is neglected altogether) of the electron population found upstream have “escaped” from the shock. In the application of Liouville’s theorem which follows, the trajectories of these electrons will need to be followed backward in time. They were “heated” not by the local shock layer, but by processes much deeper at the downstream boundary condition. In recent work, Mitchell and Schwartz [2013a, 2013b] show how this downstream population is made up of electrons which first crossed the Earth’s curved bow shock at distant locations. More discussion of the two-way Liouvillemapping of electron pitch angle distributions can be found in Lefebvre et al. [2007].

3. \( \Phi = 0 \) and \( T_\perp / B \)

The simplest mapping is one for which there is no cross-shock potential. Equation (2) then reveals that the particle trajectories remain on circles. Since Liouville’s Theorem implies that the phase space density is constant following a particle trajectory, mapping the trajectories on the boundary of the water bag reveals that the downstream distribution will be a circle of the same unitary radius. This is despite the fact that each mapped particle trajectory will see its perpendicular kinetic energy \( mv^2_{\perp} / 2 \) increase by the field ratio \( \alpha \). Accordingly \( T_{d\perp} = T_{u\perp} \) Contrary to the intuitive expectation from magnetic moment conservation that \( T_\perp / B = \text{constant} \). In this simple example, that expectation does not hold because some electron trajectories
Figure 1. Water bag model distribution function. (top) Upstream distribution, which is constant in height out to the circular boundary and zero beyond. This sketch also represents the downstream distribution in the case where the cross-shock potential $\Phi = 0$. (bottom) The corresponding Liouville-mapped downstream distribution (roughly to scale for $\alpha = 3$ and $\Phi = 9$). The crescent at $v_\parallel > 0$ corresponds to the transmitted electrons, the trajectories of which have been followed forward in time, while that at $v_\parallel < 0$ have been followed backward in time from their endpoint in the upstream to their starting point downstream. The remainder of phase space does not have trajectories which connect from upstream to downstream. The left portion, and some of the right portion, need to be filled by electrons originating downstream but which do not have sufficient energy to overcome the potential barrier, or by electrons scattered from the transmitted population.

which initially have $v_\parallel > 0$ increase $v_\perp$ at the expense of $v_\parallel$ until they mirror. Note that those mirrored trajectories occupy portions of $v_\parallel < 0$ in the upstream region not accessible from downstream due to magnetic focusing, resulting in a full, closed circular contour both upstream and downstream.

Mirroring is not necessary to break the link between magnetic moment conservation and $T_\perp / B$ as revealed below. In general, the parallel and perpendicular degrees of freedom are coupled, and the CGL [Chew et al., 1956] expectation of $T_\perp / B = \text{constant}$ requires closure assumptions about the parallel degree of freedom.

4. $\Phi \neq 0$

The mapping for $\Phi \neq 0$ is shown in Figure 1 (bottom) for a field compression ratio of $\alpha = 3$ and a potential $\Phi = 9$. Note that for values of $\Phi > \alpha - 1$ there is sufficient energy for all trajectories within the initial water bag contour to map into the downstream region; none mirror although in practice some might depending on the detailed electric and magnetic profiles [Lefebvre et al., 2007].

The moments of the upstream and downstream mapped distributions can be calculated analytically in this idealized model (see Appendix A). The perpendicular heating and downstream temperature anisotropies are plotted in Figure 2 as the solid lines. While the model is too idealized to capture all the electron physics, it is straightforward to fill in some minimalistic sense the large void in phase space left by the Liouville mapping. That region is delineated by dashed lines in Figure 1, and the resulting moment properties are displayed as dashed curves in Figure 2. Other infilling strategies might involve something closer to isotropization, but the purpose here is to illustrate what Liouville mapping does, and does not, predict.

It is clear from Figure 2 that perpendicular heating proportional to the magnetic field strength, which here would give $T_\perp / T_\parallel = 3$, only occurs for $\Phi = \alpha - 1 = 2$. Without infilling, weaker perpendicular heating results. Interestingly, filling the void increases the heating above “adiabatic” despite the fact that the void is obviously elongated in the parallel direction.

On the other hand, the mapping does show that the parallel anisotropy increases with increasing $\Phi$ as Mozer and Sundkvist [2013] suggest. The analytical expression for $T_\parallel$ asymptotes to the value of $\Phi$ for large $\Phi$. Infilling the void leads to a cooling of the resulting parallel temperature, which asymptotes to a third of the value of $\Phi$ in this case.

5. The Downstream Boundary Condition

The mapping exercise displayed above starts with an upstream distribution. As discussed earlier, the fact that the solar wind bulk speed is subthermal means that a large fraction of electrons actually observed
upstream have traversed the local shock from the downstream side. Thus, this portion of the distribution is not the corresponding portion of the ambient pristine solar wind, but a mapping of electrons in the magnetosheath which have encountered some other portion of the bow shock [Mitchell and Schwartz, 2013a]. Accordingly, much of the $v_\perp<0$ region in Figure 1 (bottom) should actually be filled by a pre-existing and nonlocally heated electron population. Parts of the void for $v_\parallel>0$ will correspond to downstream electrons with insufficient energy to escape upstream, or to electron trajectories trapped within the shock layer itself [Lefebvre et al., 2007].

This discussion goes to the heart of the question of dissipation at the bow shock. Ordinarily, dissipation is associated with an increase in entropy and therefore an irreversible (in time) process. Scattering of electrons by fluctuations would also seem to be an irreversible process, which therefore cannot capture the physics of electrons escaping from far downstream into the upstream regions. Perhaps the scattering is so strong that it prevents downstream electrons from participating in the ramp processes. However, such strong scattering might be expected to bring the electron distributions not only toward isotropy, which requires only pitch angle scattering, but also toward thermodynamic equilibrium rather than the ubiquitous flat-topped distributions. In brief, electrons arriving from further downstream do not represent heating associated with dissipative processes that increase the local entropy.

This is less a criticism of Mozer and Sundkvist [2013] per se than a question for wider debate. By contrast, the original DC field discussion only required an analysis of the shock energetics, not dissipation. Since scattering must occur, this question of dissipation, or indeed the use of the term heating, remains an open one.

6. NIF Potential

The choice of reference frame is purely one of convenience, and the expected and observed differences in the DC potential between the NIF and dHT frames are well documented [Goodrich and Scudder, 1984; Thomsen et al., 1987]. Working in the dHT frame has the advantage, even in the case in which there is scattering by AC fields, that the transverse drift along the shock does not enter into the DC electron energetics. In their reply, Mozer and Sundkvist [2013] provide a correction to their paper and describe the procedure by which they suggest that their reported fields are actually shown in the dHT frame rather than the NIF one, despite the fact that the downstream perpendicular field is nonzero. While it is not of major concern to the main thrust of this conclusion, the frame “tied to the plasma” is not the dHT frame because it does not render the shock at rest. Of course, the physics of the problem is frame independent. However, the quantitative contributions of different electromagnetic fields is not.

Mozer and Sundkvist [2013] argue that electron demagnetization due to scattering enables them to tap part of the “reservoir” of energy available in the cross-shock NIF potential. Indeed, demagnetization does inhibit, for example, the electron drift along the shock which would otherwise reduce the electron energy gain from the NIF potential to the dHT one [Goodrich and Scudder, 1984] in the quasi-perpendicular case or to the adiabatic limit in the exactly perpendicular case [Balikhin et al., 1993; See et al., 2013]. Thus, in general electron scattering should lead to more rather than less electron heating.

Mozer and Sundkvist [2013] offer no explanation as to how scattering can prevent electrons traversing the shock from gaining the full potential drop, or at least a significant fraction of it if the transverse drift is not
completely eliminated. That energy must reappear somewhere, and since the length and timescales are all associated with electron physics this would naively suggest that the energy can at best be redistributed amongst the electron population. Perhaps some of it can be radiated away [e.g., Sundkvist et al., 2012].

7. Demagnetization

In their reply, the authors present electron demagnetization as the main point of their paper. However, the title, the last sentence of the abstract, and the last sentence of section 1 make it clear that they are advocating a strong link between such demagnetization and the amount of electron heating.

The actual demagnetization evidence consists of two pieces. First, they report evidence for large-amplitude short-scale electric fields, the impact of which “must” result in demagnetization. While such observations are welcome, it is not possible to deduce the degree of that demagnetization nor the quantitative consequences it has on either individual electron trajectories or the bulk electron heating. To do so would not be easy, of course, and would require a detailed understanding of the spatial and temporal behavior of the fields and particles. Second, they report an example together with anecdotal statistics of deviations in $T_{\perp}/B$ from expectations based on their interpretation of adiabatic theory. The purpose of the present comment was to reveal the shortcomings of those expectations. They have chosen not to address those enumerated shortcomings in their reply.

Accordingly, it is not possible to establish a quantitative link between electron demagnetization and heating. The preexisting paradigm includes electron response to macroscopic fields and electron scattering in microscopic fields, including possible electron demagnetization due short-scale DC and AC electromagnetic processes. Their paper contributes to the discussion of the relative importance of these ingredients through its reporting of detailed microscopic fields but does not represent a major shift.

8. Conclusions

Mozer and Sundkvist [2013] present some interesting observations of large-amplitude small-scale fluctuating fields within the ramp of the Earth’s bow shock. While this comment calls into question some of the methodologies and assertions in their paper, it is clear from the discussion here and in previous literature that scattering must play an important role in the electron physics at shocks. This is evident, for example, from the fact that the crescent-shaped features in Figure 1 occupy only a small portion of phase space. However, that scattering operates within the framework of the DC shock fields, and there is ample evidence to suggest that those fields account for many features of the electron population. In particular, the dHT potential sets the overall scale of the spread in velocity. In the model of Mitchell and Schwartz [2013a] the potential actually adjusts to respond to the velocity scale of electrons arriving from deeper in the magnetosheath. Those electrons fill a large portion of velocity space. Total or partial demagnetization of the electrons may cause the behavior to deviate from strictly adiabatic predictions, but the extent of that deviation is not yet understood. Moreover, the multiscale, collisionless nature of the bow shock continues to raise questions about the meaning of heating and dissipation in this context.

Appendix A: Moments of the Idealized Distributions

For completeness, this appendix shows the moment calculations from which the results shown in Figure 2 are drawn. In general, for the $v_\parallel$-symmetric distributions treated here, the density, parallel, and perpendicular temperature moments are calculated as

$$\left\{ \begin{array}{c} \frac{1}{2} k_B T_\parallel^{(max)} \\ \frac{1}{2} k_B T_\parallel^{(min)} \end{array} \right\} = 2\pi \int_{v_\perp=0}^{v_\perp=0} \int_{v_\perp=0}^{v_\perp=0} \left\{ \begin{array}{c} \frac{1}{2} m v_\parallel^2 \\ \frac{1}{2} m v_\perp^2 \end{array} \right\} F dv_\parallel dv_\perp \quad \text{(A1)}$$

where the factor of two inside the integration exploits the $v_\parallel$ symmetry, and the leading factor of two in the $T_\perp$ moment reflects the two degrees of freedom in the perpendicular directions. $F$ is the height of the water bag. Note that since $T = (nk_B T)/n$, the height does not influence the temperatures.
The outer \( v_u = 1 \) boundary maps to the circle \( v^2_{||} + v^2_{\perp} = \Phi + 1 \), while the \( v_\| = 0 \) boundary maps to the ellipse \( v^2_{||} + \beta^2 v^2_{\perp} = \Phi \) where \( \beta^2 = (\alpha - 1)/\alpha \). The crescent-shaped regions in Figure 1 are the intersections between these two shapes. The resulting moments are

\[
n_u = \frac{4\pi}{3} F
\]

\[
n_u \frac{1}{2} k_B T_{\|} = n_u \frac{1}{2} k_B T_{\perp} = \frac{2\pi m F}{15}
\]

\[
n_d = \frac{4\pi}{3} F \left\{ \left[ \Phi^{3/2}_{\|} - \left( \Phi_{\|} - \alpha \right)^{3/2} \right] - \beta \left[ \Phi^{3/2}_{\perp} - \left( \Phi_{\perp} - \alpha \right)^{3/2} \right] \right\}
\]

\[
n_d \frac{1}{2} k_B T_{\perp} = \frac{2\pi m F}{15} \left\{ \left[ \Phi^{3/2}_{\|} - \left( \Phi_{\|} - \alpha \right)^{3/2} \right] - \beta \left[ 2\Phi^{5/2}_{\perp} - \left( 2\Phi_{\perp} + 3\alpha \right) \right] \right\}
\]

where \( \Phi_{\|} \equiv \Phi + 1 \) and \( \Phi_{\perp} \equiv \alpha \Phi / (\alpha - 1) \). The moments of the filled distribution can be written down from the above expressions by omitting the terms in square brackets multiplied by \( \beta \), corresponding to setting \( v_\perp \) (min) = 0 in the moment integration.

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