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Heavy ion reflection and heating by collisionless shocks in polar solar corona

Gaetano Zimbardo

Department of Physics, Universitá della Calabria, I-87036 Arcavacata di Rende, Italy

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ABSTRACT

We propose a new model for explaining the observations of preferential heating of heavy ions in the polar solar corona. We consider that a large number of small scale shock waves can be present in the solar corona, as suggested by recent observations of polar coronal jets by the Hinode and STEREO spacecraft. The heavy ion energization mechanism is, essentially, the ion reflection off supercritical quasi-perpendicular collisionless shocks in the corona and the subsequent acceleration by the motional electric field $\mathbf{E}=-(1/c)\mathbf{V}\times\mathbf{B}$. The acceleration due to $\mathbf{E}$ is perpendicular to the magnetic field, giving rise to large temperature anisotropy with $T_\perp \gg T_\parallel$, which can excite ion cyclotron waves. Also, heating is more than mass proportional with respect to protons, because the heavy ion orbit is mostly upstream of the quasi-perpendicular shock foot. The observed temperature ratios between O^{5+} ions and protons in the polar corona, and between α particles and protons in the solar wind are easily recovered. We also discuss the mechanism of heavy ion reflection, which is based on ion gyration in the magnetic overshoot of the shock.

1. Introduction

Understanding the dynamics of collisionless shocks is one of the most important problems in space and astrophysical plasmas. Collisionless shocks are observed in such diverse environments as solar flare blast waves, planetary bow shocks, the solar wind termination shock, supernova remnant shocks, and the galactic wind termination shock. Here we propose a new model, based on the presence of collisionless shocks, for explaining the preferential heating of heavy ions which is observed both in the polar solar corona and in the solar wind. In particular, in the nearly collisionless corona, beside the high temperatures, Soho/UVCS observations have shown that heavy ions in coronal holes, like O^{5+} and Mg^{9+}, are heated more than protons, and that heavy ion heating is more than mass proportional; further, mild temperature anisotropy is observed for protons, with $T_\perp/T_\parallel \simeq 2-3$, while strong temperature anisotropy is found for O^{5+}, with $T_\perp/T_\parallel \simeq 10-100$ (Kohl et al., 1997, 1998; Cranmer et al., 1999a). Here, $T_\perp$ ($T_\parallel$) stands for the perpendicular (parallel) kinetic temperature. As a consequence of magnetic mirroring in the diverging magnetic field of coronal holes, heavy ions are observed to be faster than protons in the solar wind (Kohl et al., 1998; Marsch and Tu, 2001). In addition, the collisional coupling with protons indicates that the Mg^{9+} heating has to be faster than minutes (Esser et al., 1999). These observations give stringent constraints on the heating mechanism for polar corona. Ion cyclotron heating has been considered since long (e.g., Marsch et al., 1982; Isenberg and Hollweg, 1983; Cranmer et al., 1999b; Hollweg and Isenberg, 2002), but a full understanding is still lacking.

Shock waves are considered to be common in the chromosphere/transition region and in the corona (e.g., Yokoyama and Shibata, 1995, 1996; Ryutova and Tarbell, 2003; Vecchio et al., 2009). Typically, photospheric convection leads to the emergence of small magnetic bipoles, which lead to magnetic reconnection with the large scale unipolar magnetic field of coronal holes; small scale plasma jets are formed in the reconnection regions, and fast shocks can form when jets encounter the ambient plasma (Shibata et al., 1992; Yokoyama and Shibata, 1996; Tsuneta, 1997; Tsuneta and Naito, 1998; Lee and Wu, 2000). Indeed, recent X-ray Hinode and UV STEREO observations have shown that many more plasma jets are present in the polar corona than previously thought (Cirtain et al., 2007; Patsourakos et al., 2008; Nisticó et al., 2009). Therefore, a large number of small scale, nonstationary shocks can form in the polar reconnection regions and propagate toward the high altitude corona. Further, the nonlinear steepening of compressional pulses in the collisionless high corona can lead to the formation of fast shocks, as has been observed by Cluster in the solar wind (Lee et al., 2009). In the high corona large scale shocks, which are detected as type II radio bursts (Nelson and Melrose, 1985), are launched by solar flares and by the emergence of coronal mass ejections (Aurass and Mann, 2004; Aschwanden, 2005; Magdalenic et al., 2008). These shocks can travel all the way from active regions to coronal holes, thus contributing to the high altitude heating of heavy ions. In the

E-mail address: zimbardo@fis.unical.it

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case of solar flares, the associated reconnection outflow termination shocks can be so strong as to accelerate electrons to 100 keV energies in a fraction of a second (Tsuneta and Naito, 1998; Warmuth et al., 2009), and in some cases both the upper and the lower termination shocks are identified in radio observations (Aurass and Mann, 2004). We consider that thanks to magnetic reconnection a similar scenario can be found, on a smaller scale, in coronal holes.

We notice that previously collisionless shock heating of polar corona heavy ions was proposed by Lee and Wu (2000), and that Hollweg and Isenberg (2002) have criticized this idea on the basis of the argument that if many shocks were present in the corona, the associated velocity jump would cause a large broadening of the observed emission lines, due to the Doppler shift, which would not be consistent with the observations. However, the plasma volume occupied by shocks is very small compared to that of the emitting regions, so that the Doppler broadening due to the velocity jump can be difficult to detect. In addition, an estimate of the power dissipated by shocks in the low corona, given in the following, shows that only a small fraction of the polar corona surface has to be covered with shocks, thanks to the high heating efficiency of shocks. Therefore the associated Doppler broadening can be very hard to detect. Nevertheless, current observations by Hinode and STEREO are revealing the presence of fast flows and Doppler broadened lines in the low corona (Kamio et al., 2009; Subramanian et al., 2010).

The shock heating of heavy ions considered by Lee and Wu (2000) was mostly in connection with subcritical shocks, that is for the cases when ions are not reflected by the shock surface. Here, we propose that the preferential heating of heavy ions in polar coronal holes is due to ion reflection at supercritical quasi-perpendicular shocks and to the ion acceleration by the $\mathbf{V} \times \mathbf{B}$ electric field in the shock frame. In this connection, we notice that more than mass proportional heating of α particles and O^{6+} ions has been observed in the solar wind by the Ulysses spacecraft, between 2.7 and 5.1 AU, downstream of interplanetary shocks, and that most of those shocks were supercritical (Berdichevsky et al., 1997).

2. The model

In the low β corona, a shock wave is formed when a superAlfvénic plasma flow having velocity $V_1 > V_A$ collides with the ambient coronal plasma. Here, the plasma β is given by $\beta = 8\pi p/B^2$, where p is the total plasma pressure, B is the magnetic field magnitude, V_1 is the plasma velocity upstream of the shock, and $V_A = B/\sqrt{4\pi\rho}$ is the Alfvén speed, with ρ the mass density. The Alfvénic Mach number is defined as $M_A = V_1/V_A$, and the fast Mach number as $M_f = V_1/V_f$, with V_f the fast magnetosonic speed. The plasma velocity in the reconnection outflow region is of the order of V_A in the inflow region, that is much larger than V_A in the outflow region, and this leads to the formation of shocks (e.g., Yokoyama and Shibata, 1995; Tsuneta and Naito, 1998; Aschwanden, 2005). We notice that although the typical Alfvén velocity in the corona, of the order of 1000 km/s, is larger than the observed jet velocities of 200–800 km/s (Cirtain et al., 2007; Moreno-Insertis et al., 2008), the Alfvén velocity in the reconnection outflow region can be much lower, since B is weaker within the magnetic field reversal. These shocks are effectively collisionless even in the low polar corona, since the shock dynamics evolves on times of the order of a few proton gyroperiods (about 10^{-4} s in a 10 G field) while the collision times are on the order of tens of seconds at least (Esser et al., 1999).

It is well known both from laboratory (Paul et al., 1965; Phillips and Robson, 1972) and from spacecraft experiments (e.g., Gosling and Robson, 1985; Scudder et al., 1986; Bale et al., 2005) that above a critical Mach number M_f^* a fraction of ions, which grows with the Mach number (Phillips and Robson, 1972; Leroy et al., 1982; Sckopke et al., 1983; Quest, 1986), is reflected off the shock, leading to the so-called supercritical shocks. When the angle θ_{Bn} between the shock normal (pointing in the upstream direction) and the upstream magnetic field is larger than about 45°, the reflected ions reenter the shock after gyrating in the upstream magnetic field. Such shocks are termed quasi-perpendicular. Conversely, for $\theta_{Bn} < 45°$, the reflected ions propagate upstream, forming the ion foreshock which characterizes the quasi-parallel shocks. The critical Mach number M_f^* is about 2.7 for low β perpendicular shocks, and can decrease to 1.5 for quasi-perpendicular shocks in a warm plasma (Edmiston and Kennel, 1984). Such moderate Mach numbers should be easily attained in the reconnection processes of the corona, so that ion reflection is a relatively common process. Ion reflection can be considered to be the main dissipation mechanism by which supercritical collisionless shocks convert the directed flow energy into heat, while the electrons are heated less (Gosling and Robson, 1985; Veltri and Zimbardo, 1993a, b).

For the solar corona, we consider a quasi-perpendicular supercritical collisionless shock, and we assume a simple one dimensional shock structure. The upstream quantities are indicated by the subscript 1, and the downstream quantities by the subscript 2. We adopt the normal incidence frame (NIF) of reference, in which the shock is at rest, the upstream plasma velocity is along the x axis and perpendicular to the shock surface, $\mathbf{V}_1 = (V_{x1}, 0, 0)$, the upstream magnetic field lays in the xz plane, $\mathbf{B}_1 = (B_{x1}, 0, B_{z1})$, so that the motional electric field $\mathbf{E} = -\mathbf{V} \times \mathbf{B}/c$ is in the y direction, $E_y = V_{x1}B_{z1}/c$. We further assume that $B_{z1} \gg B_{x1}$ ($\theta_{Bn} \simeq 90°$), in order to simplify the discussion. In this section, we assume that ion reflection is specular, i.e., that the component of the ion velocity perpendicular to the shock surface reverses its sign at the shock ramp. Actually, the motion of the reflected ions is more complicated, depending also on the increase of the magnetic field in the shock foot, ramp, and overshoot. For instance, Sckopke et al. (1983) have shown, by analysing ISEE 1 and ISEE 2 data at the Earth's bow shock, that nonspecularly reflected ions are present in front of the shock, together with ions which *appear* to have been specularly reflected. As shown by the theoretical and numerical analysis of Gedalin (1996), the shock magnetic field has a major role in the reflection process, and the velocity dispersion of the reflected ions is much larger than that corresponding to specular reflection (see Section 3 for more details). Because of mathematical simplicity, we assume here that ion reflection is specular; the effects of deviations from specular reflection are considered in Section 3.

An order-of-magnitude estimate of the energy gained by ions after reflection at the shock can be obtained by approximating the reflected ion trajectory with a circle of radius ρ_i, with ρ_i the ion gyroradius, in the upstream magnetic field. Assuming specular reflection, on average the ion velocity at the reflection point is perpendicular to the shock and along the negative direction of the x axis. Keeping in mind that ion reflection gives rise to a nonadiabatic displacement in the y direction, the work W done by the electric field is

$$W = q_i E_y \Delta y, \tag{1}$$

where $\Delta y \sim 2\rho_i$. For specularly reflected ions, the gyroradius has to be evaluated with the upstream flow speed (neglecting the thermal velocity of the incoming ion distribution), i.e., $v_\perp \simeq |V_{x1}|$, so that

$$W \sim q_i E_y \times 2\rho_i = 2q_i \frac{V_{x1}B_{z1}}{c} \frac{m_i V_{x1} c}{q_i B_{z1}}, \tag{2}$$