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Theory of the formation of a collisionless Weibel shock: pair vs. electron/proton plasmas

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Abstract

Collisionless shocks are shocks in which the mean-free path is much larger than the shock front. They are ubiquitous in astrophysics and the object of much current attention as they are known to be excellent particle accelerators that could be the key to the cosmic rays enigma. While the scenario leading to the formation of a fluid shock is well known, less is known about the formation of a collisionless shock. We present theoretical and numerical results on the formation of such shocks when two relativistic and symmetric plasma shells (pair or electron/proton) collide. As the two shells start to interpenetrate, the overlapping region turns Weibel unstable. A key concept is the one of trapping time τ_p, which is the time when the turbulence in the central region has grown enough to trap the incoming flow. For the pair case, this time is simply the saturation time of the Weibel instability. For the electron/proton case, the filaments resulting from the growth of the electronic and protonic Weibel instabilities, need to grow further for the trapping time to be reached. In either case, the shock formation time is $2\tau_p$ in two-dimensional (2D), and $3\tau_p$ in 3D. Our results are successfully checked by particle-in-cell simulations and may help designing experiments aiming at producing such shocks in the laboratory.

Keywords: Collisionless shocks; Weibel instability

1. INTRODUCTION

Shockwaves constitute one of the most basic concepts in fluid mechanic. Already in 1808, Poisson derived the non-linear equation for the evolution of a large amplitude sound wave (Poisson, 1808). Few years later, in 1848, Stokes understood that the solutions of the equation derived by Poisson necessarily evolve a discontinuity, that is, what we now call a "shockwave" (Stokes, 1848; Salas, 2007).

Shockwaves can also be generated when two media move with respect to each other at a velocity larger than the speed of sound in one of them. Owing to such a large amount of potential, and mundane, generators, shockwaves are ubiquitous in fluid mechanic. Of course, the shock front is not a mathematical discontinuity. In order to slow down at the front, an upstream particle can only collide more often with the others. The shock front is therefore a few mean-free path thick, that is, a discontinuity in the fluid limit. Here, the front size is of the order of the mean-free path (Zel'dovich & Raizer, 2002).

Let us now turn to the Earth bow shock, namely, the shock of the Earth magnetosphere within the solar wind. *In situ* measurements by the 4 "Cluster" satellites indicate a shock front some 100 km thick (Bale *et al.*, 2003; Schwartz *et al.*, 2011). Yet, the proton mean-free path at this location is about 10^8 km (Kasper *et al.*, 2008). Here, the shock front is six orders of magnitude *smaller* than the mean-free path. How can this be?

Starting with the pioneering works of Sagdeev (1966), such shocks have been studied from the 1960s and dubbed "collisionless shocks". A number of review paper are now available, where the reader will find about our current knowledge regarding their characteristics, both macroscopic and

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microscopic (Treumann, 2009; Marcowith *et al.*, 2016). Among the various aspects of these shocks that are currently under investigation, the way they are formed is of significant importance, be it to guide experimental efforts aiming at producing them in the laboratory (Fiuza *et al.*, 2012; Chen *et al.*, 2013; Sarri *et al.*, 2015). Such is the topic of this paper.

Beside their ubiquity in astrophysics, collisionless shocks have been under scrutiny for decades (Bell, 1978; Blandford & Ostriker, 1978), because they could hold the key to two of the most intriguing contemporary enigmas: high-energy cosmic-rays and gamma-ray bursts (Vietri *et al.*, 2003; Piran, 2005).

As previously noted, a fluid shock can be launched by colliding two media at a velocity larger than the speed of sound in one of them. However, if both media are collisionless plasmas shells, there is no "collision", as the two shells will start interpenetrating instead. At low energy, the interaction between them can be mediated by the potential jump originating from the Debye sheets at their borders (if any) (Stockem *et al.*, 2014). But at high energies, each shell seamlessly runs over the Debye sheet of the other, and the region where the two shells overlap features a counter-streaming plasma system. This counter-streaming system is the key to the shock formation.

Counter-streaming plasma systems are notoriously unstable. Consider Figure 1, with symmetric, initially cold shells, heading toward each other with a Lorentz factor γ_0. The dominant instability in the overlapping region is the Weibel one as soon as $\gamma_0 > \sqrt{3/2}$ (Bret *et al.*, 2013). It grows unstable modes with a wave vector normal to the flow.

Although the unstable "history" differs whether one deals with pair or electron/proton plasmas, the outcome is the same in both cases: after a time τ_p that we shall call the "trapping time", the flow which keeps entering the overlapping region is trapped inside. The growth of one, or various, instabilities generated enough turbulence to block the incoming flow. From the macroscopic point of view, we thus have a region of space, the overlapping region at $t = \tau_p$, where the bulk velocity is 0. This is the onset of the shock formation.

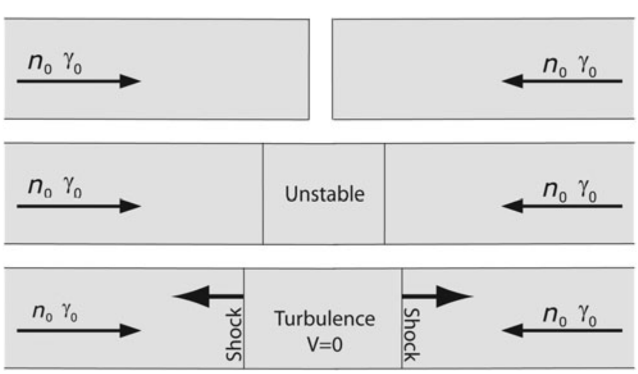

Fig. 1. Setup considered. Two plasma shells are heading toward each other. As they overlap, the overlapping region turns Weibel unstable. After the trapping time τ_p, the turbulence in the central region has become strong enough to trap the incoming flow. The shock formation follows. In Section 2, two pair plasma shells are considered. In Section 3, these are two electron/proton plasmas.

As already hinted, the saturation time is not reached the same way whether the system consists in pair, or electron/proton plasmas. This is why the rest of the paper is divided into two main parts. In Section 2, the case of pair plasmas is examined. We also explain here how the system evolves from the trapping time to the shock formation time. Then in Section 3, we turn to the slightly more involved case of electron/proton plasmas. Finally, the validity of the Rankine–Hugoniot (RH) relations for a collisionless shock is discussed in Section 4.

2. PAIR PLASMAS

The encounter of two relativistic, symmetric, un-magnetized, and cold pair plasmas shells is the simplest possible setting in various respects. The absence of mass difference between species renders particle-in-cell (PIC) simulations easier. In the cold un-magnetized regime, we have exact analytical expressions for the growth rates of the instabilities involved (Bludman *et al.*, 1960; Faĭnberg *et al.*, 1970; Bret *et al.*, 2010*a*). Finally, the dominant instability (Weibel) is always the same in the relativistic regime (Bret *et al.*, 2005; Bret & Deutsch, 2006; Bret *et al.*, 2008).

The setup considered is pictured in Figure 1. In the present case of two pair plasmas, both electrons and positrons turn unstable when the shells start to overlap. All kinds of unstable modes grow out of the interaction, but the fastest growing ones are the Weibel modes. There are found for a wave vector normal to the flow, which is why they form filaments (Huntington *et al.*, 2015). In the relativistic regime, their growth rate is (Bret *et al.*, 2010*b*),

$$\delta = \sqrt{\frac{2}{\gamma_0}}\omega_p^{-1}, \tag{1}$$

where $\omega_p^2 = 4\pi n_0 e^2/m_e$ is the electronic, or positronic, plasma frequency of the shells.

We thus have the Weibel instability growing, until it saturates at the saturation time τ_s which can be determined in terms of the growth rate (1), the field at saturation, and the spontaneous fluctuations of the shells before they collide (Bret *et al.*, 2013). When Weibel reaches saturation, the density in the overlapping region is still the sum of the two shells' density. The reason for this is that just before saturation, the linear regime was still valid, imposing only slight density perturbations. If, then, the density in the central region was still $\sim 2n_0$ at $\tau_s - \epsilon$, it cannot be otherwise at τ_s. At this stage, the density "jump" is therefore only 2, far from the one expected from the RH relation (see Section 4).

How then is a shock formed from this point? By trapping the incoming flow in the overlapping region. Indeed, this region of space is now occupied by magnetic filaments, which peak field value and characteristic length are known

(Davidson *et al.*, 1972). From these data, one can compute the distance L the incoming flow can cover before being randomized (Lyubarsky & Eichler, 2006; Bret, 2015). As it turns, L is always smaller than the size of the overlapping region at saturation (Bret *et al.*, 2013, 2014). As a consequence, this region no longer expands after saturation, and the bulk velocity inside is 0. Here, the trapping time is therefore identical to the saturation time, $\tau_p = \tau_s$.

A shock has just been formed in velocity space. From this moment, the density in the central region is going to increase until it meets the expected density jump, namely 3 in 2 dimensions, and 4 in 3 dimensions (Blandford & McKee, 1976; Stockem *et al.*, 2012). If it took one saturation time τ_s to bring the central density from 1 to 2, and assuming the same region no longer expands after saturation, then we shall need to wait another saturation time to raise the density jump to 3 [in two-dimensional (2D)], and yet another saturation time to raise it to 4 (in 3D). The shock formation time is therefore given by,

$$\tau_f = d\tau_p, \tag{2}$$

where d is the dimension of the problem. Note that this line of reasoning holds regardless of the shells' composition. It only relies on the expected density jump, and on the flow being trapped at trapping time. Hence, Eq. (1) is also valid for the case we now turn to, namely, electron/proton plasmas shells.

3. ELECTRON/PROTON PLASMAS

This case can still be pictured by Figure 1, where each shell is now made of protons and electrons. Due to their smaller inertia, electrons turn Weibel unstable first, and grow magnetic filaments of size λ_e and peak field B_e. Meanwhile, protons keep ploughing through the resulting bath of hot electrons, and in turn get unstable. Here also, a number of unstable modes grow, but investigations of the full unstable spectrum showed that the fastest growing modes are, again, the Weibel's ones (Shaisultanov *et al.*, 2012).

As it is triggered, the protons' Weibel instability does not start from a perfectly clean medium, but from a plasma where magnetic filaments of size λ_e can already be found. Simply put, an unstable mode is already seeded, and it is further grown by the protons' Weibel instability. When it reaches saturation, the peak field in the filaments in now

$$B_i = \sqrt{\frac{m_p}{m_e}} B_e, \tag{3}$$

where m_p is the proton mass. But their size is still the same, since the linear regime leaves unchanged the $\mathbf{k}$ it grows.

Can magnetic filaments of size λ_e and peak field B_i efficiently trap the incoming ions? The answer is no. The main problem is that λ_e is of the order of the electronic Larmor radius in the field B_e. In order to block the protons, λ_e should also be the Larmor radius of the protons in the field B_i (Bret, 2015). Considering Eq. (3), one finds λ_e is $\sqrt{m_e/m_p}$ too short for this. Therefore, at this stage, protons keep streaming through the overlapping region.

Fortunately, another process, widely studied in relation with the non-linear evolution of the Weibel instability, allows reaching the trapping time. It has been known for long that in this regime, the magnetic filaments originated by the growth of the instability, progressively merged, and increase in size. A model for this growth, backed up by 2D and 3D PIC simulations, found that the size of the filaments grows linearly with time (Medvedev *et al.*, 2005). The same model allows us therefore to compute the time it takes for the filaments to reach the correct size. Summing up this merging time to the saturation time of the protons' Weibel instability (that of the electrons can be neglected), we find for the trapping time (Stockem Novo *et al.*, 2015),

$$\tau_p = 4.43\sqrt{\lambda_0} \ln\frac{m_p}{m_e} \omega_{pp}^{-1}, \tag{4}$$

where $\omega_{pp}^2 = 4\pi n_0 e^2/m_p$ is the protonic plasma frequency of the shells. Following the reasoning explained at the end of the previous section, the shock formation time is now,

$$\tau_f = 4.43 d\sqrt{\lambda_0} \ln\frac{m_p}{m_e} \omega_{pp}^{-1}. \tag{5}$$

Note that if $m_p = m_e$, this equation does not reduce to the shock formation time for pair plasmas since the electronic Weibel phase has been neglected. A series of 2D PIC simulations has been performed with the code OSIRIS to test our result (Fonseca *et al.*, 2002). The result can be found in Figure 2, evidencing a very good agreement theory-PIC.

At this junction, it is interesting to compare the shock formation time in pair and electron–proton plasmas. From Eqs (2) and (5), we get (Bret *et al.*, 2013),

$$\frac{\tau_{f,ep}}{\tau_{f,pairs}} = \frac{6.2}{\Pi} \ln\frac{m_p}{m_e}\sqrt{\frac{m_p}{m_e}}, \tag{6}$$

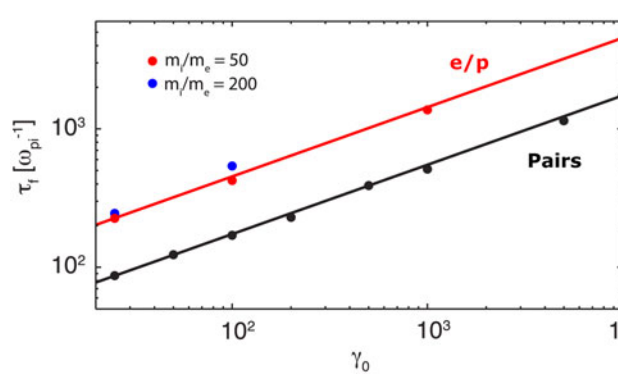

Fig. 2. Comparison of the shock formation time given by the theory with the results of PIC simulations.

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where $\Pi \sim 10$–20, is the number of exponentiations of the electronic Weibel instability. The function so defined is plotted in Figure 3 in terms of Π and m_p/m_e. For a mass ratio of 1836, a shock in electron–proton plasmas requires 100–200 more time to form than in pair plasmas.

4. RANKINE–HUGONIOT

As evidenced previously, our determination of the shock formation hinges on the validity of the RH jump condition. Within the relativistic regime where we operate, the expected RH density jump is 3 in 2D and 4 in 3D (Stockem *et al.*, 2012). To which extent should we expect its fulfilment for a collisionless shock?

These density jumps rely on the conservation of matter, momentum and energy across the shock front. If matter carries all the energy and momentum, then these jumps are correct, provided all the matter upstream goes downstream. In a collisionless chock, this latter assumption may be challenged, for at least two reasons:

1. With a mean-free path much larger than the shock front, some particles may bounce back from the upstream, or even come back upstream from the downstream (the shock front only goes at $c/2$ in 2D and $c/3$ in 3D).
2. Collisionless shocks are known to accelerate particles and grow a high-energy, non-Maxwellian tail (Krymskii, 1977; Spitkovsky, 2008). These high-energy particles may go back and forth near the shock front, before escaping upstream or downstream.

In either case, part of the initial upstream energy and momentum escapes the RH budget, blurring the jump condition. Regarding the accelerated particles, the high-energy tail grows with time like $\sqrt{t}$, gathering more and more energy (Sironi *et al.*, 2013). As a result, downstream temperatures of only 80% of the RH one have been observed in simulations (Caprioli & Spitkovsky, 2014).

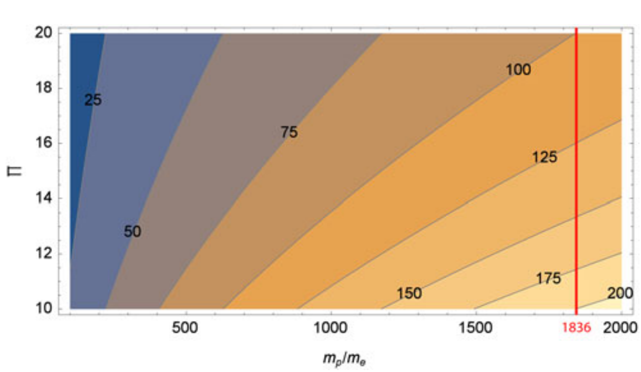

Fig. 3. Plot of Eq. (6) in terms of m_p/m_e and Π. For a mass ratio of 1836, a shock in electron–proton plasma takes 100–200 more time to form than in pair plasmas.

However, the present chock formation time relies, by definition, on the RH density jump at the very beginning of the shock history. At this stage, no particles have been accelerated yet, and we are left with the first population above. To our knowledge, there is no theory of the amount of particles departing from a pure "upstream → downstream" trajectory. For the pair case, numerical explorations of this problem found that the deviation from the RH expected jump is initially of the order of the spontaneous density fluctuations behind the front (Stockem *et al.*, 2012).

5. CONCLUSION

We have presented a theory for the formation of a collisionless Weibel shock, both in pair and electron/proton plasmas. The setup investigated consists in the generation of a shock from the encounter of two identical, cold, relativistic and unmagnetized plasma shells.

When the two shells start overlapping, the overlapping region turns unstable to the Weibel instability. A key concept related to the formation scenario is the one of trapping time τ_p, which is the time at which the instability-generated turbulence in the overlapping region is able to trap the incoming flow. Once trapping have been achieved, the shock formation time is simply $2\tau_p$ in 2D, and $3\tau_p$ in 3D, namely $d\tau_p$ where d is the dimension of the problem.

For pair plasmas, the trapping time equals the saturation time τ_s of the Weibel instability. By the time Weibel saturates, the magnetic filaments and the peak field are large enough to efficiently trap the incoming flow in the central region. Shock formation ensues at $2\tau_p$ or $3\tau_p$, depending of the geometry.

For electron/proton, the scenario leading to the shock formation is significantly longer. As the two plasma shells overlap, electrons turn Weibel unstable first. When this instability saturates, it has seeded an unstable mode that is further grown by the proton Weibel instability. This later instability then saturates in turn, but the filaments it grew are too small to trap the incoming ion flow. We must then wait for the filaments to merge, until they are big enough to trap the ion flow. This is when the trapping time τ_p is reached. As was the case for the pair plasmas, shock formation follows at $2\tau_p$ or $3\tau_p$, depending of the geometry. Equations (2) and (5) for the shock formation time in pair and electron/proton plasmas have been successfully checked by means of PIC simulations.

The ratio of the formation time in pair to the formation time in electron/proton plasmas is an interesting quantity, in particular when it comes to designing an experiment aiming at producing such shocks in the lab. Depending on the number of exponentiations of the electronic Weibel instability, and for a realistic proton to electron mass ratio, Eq. (6) shows that a shock in electron/proton plasmas requires 100–200 more time to form than in pairs. This, in turn, translates into the interaction length required between the two shells.

Although more work would be needed to determine more precisely the range of application of the RH jump conditions,

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simulations show they are tightly fulfilled at the beginning of the shock history. The reason for this is that the high-energy tail the shock grows with time, has not developed yet. Future woks may also contemplate the effect of a non-zero temperature in the shells before they collide, or that of an external magnetic field.

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