The synchrotron maser emission from relativistic magnetized shocks: Dependence on the pre-shock temperature

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ABSTRACT

Electromagnetic precursor waves generated by the synchrotron maser instability at relativistic magnetized shocks have been recently invoked to explain the coherent radio emission of Fast Radio Bursts. By means of two-dimensional particle-in-cell simulations, we explore the properties of the precursor waves in relativistic electron-positron perpendicular shocks as a function of the pre-shock magnetization $\sigma \gtrsim 1$ (i.e., the ratio of incoming Poynting flux to particle energy flux) and thermal spread $\Delta \gamma \equiv kT/mc^2 = 10^{-5} - 10^{-4}$. We measure the fraction $f_\xi$ of total incoming energy that is converted into precursor waves, as computed in the post-shock frame. At fixed magnetization, we find that $f_\xi$ is nearly independent of temperature as long as $\Delta \gamma \lesssim 10^{-1.5}$ (with only a modest decrease of a factor of three from $\Delta \gamma = 10^{-3}$ to $\Delta \gamma = 10^{-1.5}$), but it drops by nearly two orders of magnitude for $\Delta \gamma \gtrsim 10^{-4}$. At fixed temperature, the scaling with magnetization $f_\xi \sim 10^{-3} \sigma^{-1}$ is consistent with our earlier one-dimensional results. For our reference $\sigma = 1$, the power spectrum of precursor waves is relatively broad (fractional width $\sim 1 - 3$) for cold temperatures, whereas it shows pronounced line-like features with fractional width $\sim 0.2$ for $10^{-3} \lesssim \Delta \gamma \lesssim 10^{-1.5}$. For $\sigma \gtrsim 1$, the precursor waves are beamed within an angle $\sim \sigma^{-1/2}$ from the shock normal (as measured in the post-shock frame), as required so they can outrun the shock. Our results can provide physically-grounded inputs for FRB emission models based on maser emission from relativistic shocks.

Key words: magnetic fields — masers — radiation mechanisms: non-thermal — shock waves — stars: neutron

1 INTRODUCTION

Relativistic shocks are invoked as candidate sources of non-thermal particles in pulsar wind nebulae, gamma-ray bursts, and active galactic nuclei jets, and as possible accelerators of ultra-high-energy cosmic rays (e.g. Blandford & Eichler 1987). However, relativistic shocks are generally quasi-perpendicular, i.e., with pre-shock field orthogonal to the shock direction of propagation, a configuration that — if the magnetic field is sufficiently strong — inhibits efficient particle acceleration (e.g. Begelman & Kirk 1990; Sironi et al. 2013). While poor particle accelerators, relativistic magnetized perpendicular shocks can channel an appreciable fraction of the incoming flow energy into semi-coherent electromagnetic waves propagating back into the upstream medium (hereafter, “precursor waves” moving ahead of the shock). The waves are attributed to the synchrotron maser instability (Alsop & Arons 1988; Hoshino & Arons 1991). The instability is sourced by a population inversion that naturally occurs at the shock front, where a coherent ring-like distribution is constantly produced as the incoming particles gyrate in the shock-compressed field.

Recently, the discovery of Fast Radio Bursts (FRBs; for recent reviews, see Petroff et al. 2019; Cordes & Chatterjee 2019; Platts et al. 2019) has revived the interest in the synchrotron maser. FRBs are bright ($\sim 1\text{Jy}$) pulses of millisecond duration detected in the $\sim \text{GHz}$ band. Their extremely high brightness temperature, $T_B \sim 10^{37}\text{K}$, requires a coherent emission mechanism (e.g., Katz 2016). The synchrotron maser at relativistic shocks has been invoked as one of the candidate emission mechanisms within the so-called “magnetar scenario” (Lyubarsky 2014; Murase et al. 2016; Beloborodov 2017; Waxman 2017; Metzger et al. 2019; Beloborodov 2019; Margalit et al. 2020b,a), where magnetars are invoked as FRB progenitors, a hypothesis recently confirmed by the detection of FRBs from a Galactic magnetar (Scholz & CHIME/FRB Collaboration 2020; Bochenek et al. 2020). In response to motions of the magnetar crust, the above-lying magnetosphere is violently twisted and a strongly magnetized pulse is formed, which propagates away through the magnetar wind. In the shock maser scenario, the FRB is generated at ultra-relativistic shocks resulting from the collision of the magnetized pulse with the steady wind that is produced by the magnetar spin-down luminosity or by the cumulative effect of earlier flares.

The fundamental properties of the precursor waves generated by the synchrotron maser — i.e., their efficiency, power spectrum,
angular distribution and polarization — can be quantified with self-consistent particle-in-cell (PIC) simulations. PIC simulations of relativistic magnetized shocks focusing on the synchrotron maser emission have been performed both for electron-positron plasmas (Langdon et al. 1988; Gallant et al. 1992; Sironi & Spitkovsky 2009; Iwamoto et al. 2017, 2018; Plotnikov et al. 2018; Plotnikov & Sironi 2019) and electron-proton or electron-positron-proton plasmas (Hoshino et al. 1992; Amato & Arons 2006; Lyubarsky 2006; Hoshino 2008; Stockem et al. 2012; Iwamoto et al. 2019).

These works primarily focused on low magnetizations $\sigma \leq 1$, where $\sigma$ is the ratio of upstream Poynting flux to kinetic energy flux. On the other hand, FRBs are expected to originate from extreme environments where the energy content of the plasma is dominated by magnetic fields, as in magnetar winds. In Plotnikov & Sironi (2019), we performed one-dimensional (1D) PIC simulations of electron-positron shocks and investigated how the properties of the synchrotron maser depend on the flow magnetization, in the $\sigma \gtrsim 1$ regime most relevant for FRB sources. We found that the shock converts a fraction $f_s \approx 2 \times 10^{-3} \sigma^{-1}$ of the total incoming energy into the precursor waves, as measured in the post-shock (downstream) frame. At $\sigma \gtrsim 1$, we showed that the shock structure displays two solitons separated by a cavity, and the peak frequency of the spectrum corresponds to an eigenmode of the cavity. We also found that the efficiency and spectrum of the precursor waves do not depend on the bulk Lorentz factor of the pre-shock flow.

The results in Plotnikov & Sironi (2019) were obtained assuming that the pre-shock flow has small thermal spread, $\Delta y \approx kT/mc^2 \approx 10^{-4}$. In fact, with the exception of the study by Amato & Arons (2006) — who focused on nonthermal lepton acceleration in pair-proton plasmas, rather than on the properties of the precursor waves — all prior studies were conducted in the limit of negligible upstream temperatures. In this work, we discuss the dependence of the precursor waves generated by the synchrotron maser on the upstream temperature, by means of two-dimensional (2D) PIC simulations of relativistic magnetized electron-positron shocks. We focus on magnetically-dominated plasmas ($\sigma = 1$ and 3) and explore thermal spreads in the range $\Delta y \approx 10^{-5}$ to $10^{-1}$. All our simulations are evolved for sufficiently long ($\geq 4000\omega_{pe}$) so that the properties of the precursor waves, such as their Poynting flux and power spectrum, attain a steady state. At fixed magnetization, we find that the efficiency $f_s$ is nearly independent of temperature as long as $\Delta y \lesssim 10^{-1.5}$ (with only a modest decrease of a factor of three from $\Delta y = 10^{-2}$ to $\Delta y = 10^{-1.5}$), but it drops by nearly two orders of magnitude for $\Delta y \gtrsim 10^{-1}$. For our reference $\sigma = 1$, the power spectrum of precursor waves is relatively broad (fractional width $\approx 1$) for cold temperatures, whereas it shows narrow line-like features with fractional width $\approx 0.2$ for $10^{-3} \lesssim \Delta y \lesssim 10^{-1.5}$. For $\sigma \gtrsim 1$, the precursor waves are beamed within an angle $\approx \omega_{pe}^{-1/2}$ from the shock normal (as measured in the post-shock frame), as required so they can outrun the shock.

The paper is organized as follows. In Section 2 we present the numerical method and the simulation setup. We then discuss the main results of our investigation, as regard to shock structure (Section 3), precursor efficiency (Section 4) and beaming and power spectrum (Section 5). We summarize our findings in Section 6 and discuss their astrophysical implications.

2 SIMULATION SETUP

We use the three-dimensional (3D) electromagnetic PIC code TRISTAN-MP (Spitkovsky 2005) to perform simulations of relativistic magnetized shocks in pair plasmas. We perform simulations in 2D spatial domains, but all three components of particle velocities, electric currents and electromagnetic fields are retained.

The simulations are performed in the post-shock frame. The upstream flow, consisting of electrons and positrons, drifts in the $-\hat{x}$ direction with speed $-\beta_0 \hat{x}$, where $\beta_0 = (1 - 1/\gamma_0^2)^{1/2}$. For the simulations presented in this paper, we employ $\gamma_0 = 10$, but we have verified that larger values of $\gamma_0$ (up to $\gamma_0 = 80$) do not change our conclusions (see also Plotnikov & Sironi (2019), for an investigation of the dependence on $\gamma_0$ with 1D simulations). The incoming flow reflects off a wall at $x = 0$. The shock is formed by the interaction of the incoming and reflected flows and propagates along $+\hat{x}$.

The upstream temperature is cast in terms of the thermal spread $\Delta y \approx kT/mc^2$, which we vary from $\Delta y = 10^{-5}$ up to $\Delta y = 10^{-1}$. Here, $m$ is the electron mass and $c$ is the speed of light.

The pre-shock plasma carries a frozen-in magnetic field $B_0 = B_0 \hat{z}$ orthogonal to the $xy$ plane of our simulations, and the associated motional electric field $E_0 = -\beta_0 B_0 \hat{y} = -E_0 \hat{y}$. Our field configuration parallels the one employed by Iwamoto et al. (2017), and corresponds to a perpendicular shock with out-of-plane upstream field. The magnetic field strength is parameterized via the magnetization, i.e., the ratio of upstream Poynting flux to kinetic energy flux:

$$\sigma = \frac{B_0^2}{4\pi\gamma_0 N_0 mc^2} = \left( \frac{\omega_{pe}}{\omega_{p0}} \right)^2 = \left( \frac{c/\omega_p}{\omega_{p0}} \right)^2,$$

(1)

where $N_0$ is the number density of the upstream plasma (including both species), $\omega_p = eB_0/\gamma_0 mc$ is the Larmor frequency, $\omega_{p0} = (4\pi N_0 e^2/\gamma_0 m)^{1/2}$ is the plasma frequency, $c/\omega_p$ is the plasma skin depth and $\omega_{p0} = \gamma_0 mc^2/eB_0$ is the Larmor radius of particles with Lorentz factor $\gamma_0$. Here, $e$ is the electron charge. We explore two values of magnetization, $\sigma = 1$ and 3. We use the plasma skin depth $c/\omega_p$ as our unit of length and the inverse plasma frequency $\omega_{p0}^{-1}$ as our unit of time.

We employ periodic boundary conditions in the $y$ direction. The incoming plasma is injected through a "moving injector," which moves along $+\hat{x}$ at the speed of light. The simulation box is expanded in the $x$ direction as the injector approaches the right boundary of the computational domain. This permits us to save memory and computing time, while following the evolution of all the upstream regions that are causally connected with the shock (for details see, e.g. Spitkovsky 2005; Sironi & Spitkovsky 2009). Over time, the distance between the shock and the injector increases, and the incoming flow might suffer from the so-called numerical Cerenkov instability (e.g. Dieckmann et al. 2006). By employing a fourth-order spatial stencil for Maxwell’s equations (Greenwood et al. 2004), we find no evidence of numerical Cerenkov instability within the timespan covered by our simulations.

The leftmost edge of the downstream region, which is a conducting boundary for electromagnetic fields and a reflecting wall for particles (hereafter, the "wall"), is initially located at $x = 0$. The focus of this work is on upstream-propagating waves generated by the shock, rather than on the properties of the shocked plasma. In order to save memory and computing time, we choose to periodically jump the wall toward the shock, such that the average speed of the wall is $\approx 5\%$ lower than the shock speed and the wall always stays safely behind the shock (by at least a few tens of $c/\omega_p$). After every jump, we enforce conducting boundary conditions for the electromagnetic fields at the new position of the wall, and we discard particles to the left of the wall. By performing a few tests without the "jumping wall," i.e., retaining the whole downstream
region, we have verified that this strategy does not artificially affect any property of the precursor waves.\footnote{\textit{\textsuperscript{1}}}

We now describe the numerical parameters used in our work. We employ a high spatial resolution, with $c/\omega_p = 100$ cells. This ensures that the high frequency / wavenumber part of the power spectrum of precursor waves (i.e., $k c/\omega_p \gg 1$) is properly captured (Iwamoto et al. 2017). A few tests with a lower spatial resolution of $c/\omega_p = 50$ cells have shown good agreement with our production runs, so a resolution of $c/\omega_p = 50$ cells might also be sufficient. The transverse size of the box is 1440 cells, corresponding to $\sim 14 c/\omega_p$. This is sufficient to capture genuine 2D effects in the properties of the shock and of the precursor waves (e.g., filamentation of the upstream density, see Section 3). Experiments with even larger boxes give essentially the same results.

The numerical speed of light is 0.45 cells/timestep. We evolve our simulations for a few thousands of $\omega_p^{-1}$, which is sufficient to study the steady-state properties of precursor emission. Our longest simulations have been run for $\sim 6000 \omega_p^{-1}$, corresponding to more than 1.3 million timesteps.

Our 2D simulations are typically initialized with $N_0 = 16$ particles per cell (including both species) for $\sigma = 1$ and 4 particles per cell for $\sigma = 3$. For $\sigma = 1$, we have also performed simulations with $N_0 = 4$ for all the temperature values we investigated, finding excellent agreement with our reference $N_0 = 16$ cases (in fact, in Fig. 7 and Fig. 8 we use $N_0 = 4$ simulations). The simulations with the hottest upstream plasma, $\Delta \gamma = 10^{-5}$, employ 32 particles per cell (for both $\sigma = 1$ and 3), since the precursor emission is very weak, and so harder to properly capture (see Section 4). For this temperature, we have checked that simulations with even larger $N_0 = 128$ give the same results. In order to further reduce numerical noise in the simulations, we filter the electric current deposited to the grid by the particles, effectively mimicking the role of a larger number of particles per cell (Spitkovsky 2005, Belyaev 2015). We typically apply $N_{\text{sm}} = 20$ passes of a binomial low-pass digital filter at every timestep. In Appendix A, we show the effects of $N_0$ and $N_{\text{sm}}$ on the precursor wave spectrum.

In addition to 2D simulation, which constitute the bulk of this paper, we have also performed a suite of 1D simulations with $\sigma = 1$ and the same range of temperatures as in 2D. In 1D simulations, we typically employ $N_0 = 40$ particles per cell, a spatial resolution of $c/\omega_p = 112$ cells and a numerical speed of light of 0.5 cells/timestep. For the hottest temperature $\Delta \gamma = 10^{-1}$, the number of particles per cell is increased to $N_0 = 400$.

3 \textbf{SHOCK STRUCTURE}

In Figs. 1, 3 and 5 we present the 2D shock structure from our simulations with $\sigma = 1$ and three values of upstream thermal spread: respectively, $\Delta \gamma = 10^{-5}, 10^{-2}$ and $10^{-1}$. From top to bottom, in each figure we present the transverse magnetic field $B_z/B_0$, the transverse electric field $E_y/E_0$, the particle number density in units of the upstream density $N_0$, and the particle number density averaged along the $y$ direction.

\begin{figure}
\centering
\includegraphics[width=0.5\textwidth]{fig1.png}
\caption{Shock structure from the 2D PIC simulation with $\sigma = 1$ and $\Delta \gamma = 10^{-5}$ at $\omega_p t = 2000$, when the precursor emission has reached a steady state. We focus on the vicinity of the shock. We present (a) the transverse magnetic field $B_z/B_0$; (b) the transverse electric field $E_y/E_0$; (c) the particle number density, in units of the upstream density $N_0$; (d) the particle number density averaged along the $y$ direction.}
\end{figure}

\begin{figure}
\centering
\includegraphics[width=0.5\textwidth]{fig2.png}
\caption{Momentum space $\gamma \beta_y - \gamma \beta_x$ of positrons, from the 2D simulation with $\Delta \gamma = 10^{-5}$. The particles are selected at time $\omega_p t = 2000$ to be located near the soliton-like structure at the shock front, in the range $-1.5 c/\omega_p < x - x_{sh} < -1.5 c/\omega_p$. The histogram is normalized such that $N(\gamma \beta_x, \gamma \beta_y) = 1$ in the pixel with the highest value, and the color scale is stretched with 0.12 power to emphasize weak phase space structures.}
\end{figure}
Figure 3. Same as in Fig. 1, but for a 2D simulation with $\sigma = 1$ and upstream plasma temperature of $\Delta y = 10^{-2}$ at time $\omega_p t = 2000$.

Figure 4. Same as in Fig. 2, but for the 2D simulation with $\Delta y = 10^{-2}$ at time $\omega_p t = 2000$.

Figure 5. Same as in Fig. 1, but for a 2D simulation with $\sigma = 1$ and upstream plasma temperature of $\Delta y = 10^{-1}$ at time $\omega_p t = 1500$.

Figure 6. Same as in Fig. 2, but for the 2D simulation with $\Delta y = 10^{-1}$ at time $\omega_p t = 1500$. 
The shock front is at \( x = x_{\text{sh}} \). The upstream flow is on the positive side \((x - x_{\text{sh}} > 0)\) and the downstream plasma is on the negative side \((x - x_{\text{sh}} < 0)\). The existence of a well-developed shock is confirmed by the jump in number density and in \( B \), at the front location. The shock front itself exhibits a soliton-like structure, as revealed by the density spikes in panels (d) at \( x - x_{\text{sh}} \sim 0 \) (see, e.g., Alsop & Arons 1988). The density spike in the soliton is higher for colder plasmas (compare panels (d) among the three figures), as derived analytically by Chiuheh & Lai (1991). In the soliton, the incoming particles gyrate around the compressed magnetic field and form a semi-coherent ring in momentum space. As shown in Figs. 2, 4 and 6, where we plot, for different values of \( \Delta y \), the \( y_\beta \), \( -y_\beta \), momentum space of particles populating the density spike, the thickness of the ring depends on the pre-shock temperature. A cold well-defined ring appears for low temperatures (\( \Delta y = 10^{-3} \) in Fig. 2), whereas the center of the ring is nearly filled with particles for hot flows (\( \Delta y = 10^{-1} \) in Fig. 6). The radius of the ring is \(~ \gamma_0 \beta_0 \sim 10\) corresponding to the bulk four-velocity of incoming particles.

The synchrotron maser instability, and the resulting precursor waves, is believed to be sourced by the population inversion in the ring (Alsop & Arons 1988; Hoshino & Arons 2001). Such a population inversion tends to disappear for hot flows, as shown in Fig. 6. We then expect that the synchrotron maser emission will become inefficient in hot plasmas (see also Amato & Arons 2006). As a result of the synchrotron maser instability, a train of semi-coherent large-amplitude electromagnetic precursor waves is emitted toward the upstream, as shown in the \( B_z/B_0 \) and \( E_z/B_0 \) panels, for cold (\( \Delta y = 10^{-5} \) in Fig. 1) and moderate (\( \Delta y = 10^{-1} \) in Fig. 3) temperatures. As expected, no evidence of precursor waves is seen in hot plasmas (\( \Delta y = 10^{-1} \) in Fig. 5).

When the precursor emission is efficient, electromagnetic waves are seen not only in the upstream region \((x > x_{\text{sh}})\), but also right behind the leading soliton, in the density cavity at \(-2 c/\omega_p \leq x - x_{\text{sh}} \leq 0\). As discussed in Plotnikov & Sironi (2019), this cavity is a peculiarity of \( \sigma \geq 1 \) shocks. It plays an important role in setting the properties of precursor waves, since the peak frequency of the wave spectrum is observed to correspond to an eigenmode of the cavity, i.e., the precursor waves might be resonantly amplified by the density cavity. The hot case in Fig. 5 does not display such a density cavity, and in fact its precursor emission is strongly suppressed (see Section 4). For \( \sigma \geq 1 \) the precursor waves appear to be generated by an oscillating current localized near the downstream side of the cavity (at \( x \sim x_{\text{sh}} - 2 c/\omega_p \)). This is generally not accounted for within the standard description of the synchrotron maser instability (Alsop & Arons 1988; Hoshino 2001). A characterization of the cavity, and its role in setting the oscillating current that ultimately drives the precursor waves, is left for future work.

The wave vector \( k \) of the precursor waves is nearly aligned with the shock direction of propagation. The fluctuating magnetic field is along \( \xi \) (i.e., along the same direction as the upstream field \( B_0 \)), and the fluctuating electric field is perpendicular to both \( k \) and \( B_0 \). The wave is then linearly polarized and identified with the extraordinary mode (\( X \)-mode). We remark that 2D simulations with out-of-plane fields do not allow for the excitation of the ordinary mode (\( O \)-mode).

Iwamoto et al. (2017) performed 2D simulations of weakly magnetized shocks (\( \sigma \leq 1 \)) with in-plane upstream fields and showed that \( O \) modes are stronger than \( X \) modes at low magnetizations (\( \sigma \leq 10^{-2} \)), but they become weaker as \( \sigma \) increases.

\footnote{We remark that the continuous flow of plasma through the shock ensures that the population inversion is steadily maintained.}

\section{4 Precursor Efficiency}

In this Section, we quantify how the wave efficiency depends on the flow temperature and magnetization. We measure the wave intensity in a region between 5 \( c/\omega_p \) and 30 \( c/\omega_p \) ahead of the shock.
as the spatial average
\[ \langle \delta B_i^2 \rangle = \langle (B_i - B_0)^2 \rangle, \tag{2} \]
and we define the normalized wave energy as
\[ \xi_B = \frac{\langle \delta B_i^2 \rangle}{B_0^2}. \tag{3} \]

We have verified that in all our simulations \( \langle \delta B_i^2 \rangle/B_0^2 \approx (\delta E_\nu, \delta B_\nu)/(E_\nu B_0) \), i.e., the parameter \( \xi_B \) also quantifies the ratio of wave Poynting flux to incoming Poynting flux. Here, \( \delta E_\nu = E_\nu + \beta_0 B_0 = E_\nu + E_i \).

In Fig. 7, we show for different temperatures the time evolution of the normalized wave energy, for 2D simulations with \( \sigma = 1 \) (see legend). Following a transient, all the curves reach a steady state at \( \omega t \gtrsim 1000 \). The precursor efficiency is nearly independent of temperature as long as \( \Delta \gamma \lesssim 10^{-1.5} \) (with only a modest decrease of a factor of three from \( \Delta \gamma = 10^{-3} \) to \( \Delta \gamma = 10^{-1.5} \)), but between \( \Delta \gamma = 10^{-1.5} \) and \( \Delta \gamma = 10^{-1} \) drops by nearly two orders of magnitude (red curve). We have confirmed this result with 3D simulations (Sironi et al, in prep.).

The steady-state values of the normalized wave energy are shown as a function of temperature in the top panel of Fig. 8. There, we present results from 1D and 2D simulations with \( \sigma = 1 \) (blue and black points, respectively) and from 2D simulations with \( \sigma = 3 \) (red points). All the values are extracted from a time range when the precursor has achieved a steady state, more specifically at \( \omega t \gtrsim 2000 \) for \( \sigma = 1 \) and \( \omega t \gtrsim 3000 \) for \( \sigma = 3 \). As shown in the top panel of Fig. 8, 1D results for \( \sigma = 1 \) generally follow the same trend as our reference 2D simulations (compare blue and black lines), with two notable differences. At very high temperatures (\( \Delta \gamma \ll 10^{-1.5} \)), the drop in wave energy is much more dramatic in 2D than in 1D, by roughly one order of magnitude. At cold temperatures (\( \Delta \gamma \lesssim 10^{-5.5} \)), the precursor waves are stronger in 1D than in 2D. We attribute this difference in efficiency at low temperatures to the longitudinal heating induced in 2D by the filamentation mode, which is absent in 1D (see also Appendix B, for the effect of longitudinal dispersion on the precursor strength). Note that the 1D simulations presented by Plotnikov & Sironi (2019) employed \( \Delta \gamma \ll 10^{-4} \), for which 1D and 2D results only differ by \( \sim 50\% \).

As regard to the dependence on magnetization, the top panel of Fig. 8 shows that the normalized wave energy is roughly the same between \( \sigma = 1 \) and \( \sigma = 3 \): for cold and moderate temperatures (\( \Delta \gamma \ll 10^{-1.5} \)), \( \xi_B \) is a few percent, while it drops by nearly two orders of magnitude for \( \Delta \gamma \gtrsim 10^{-1} \). In the range of temperatures where the precursor is efficient, the minimum of the normalized wave energy is attained for \( \Delta \gamma \sim 10^{-1.5} \) in \( \sigma = 1 \) and for \( \Delta \gamma \sim 10^{-3.5} \) in \( \sigma = 3 \). We shall use these values to roughly distinguish between cold cases and warm cases, which, as shown in Section 5 for \( \sigma = 1 \), display different spectral properties.

From the normalized wave energy in the top panel of Fig. 8, we can extract the so-called strength parameter
\[ a = \frac{e \delta E_\nu}{mc\omega}. \tag{4} \]
At \( \sigma \gtrsim 1 \), where the typical wave frequency in cold plasmas is \( \omega \sim 3\sqrt{\gamma_0} \omega_h \) (Plotnikov & Sironi 2019), we find that \( a \sim 0.3 (\xi_B/10^{-2})^{1/2}(\gamma_0/10) \). The strength parameter could also be measured directly from the transverse motion (along \( y \)) of the upstream particles in the field of the wave, since the oscillations in \( y_\nu \) are directly related to the wave strength parameter (e.g., Lyubarsky 2006; Iwamoto et al. 2017). Notice that we expect the particle transverse oscillations to become relativistic at \( \gamma_0 \gtrsim 30 \). However, we have

\footnote{Our results do not appreciably change if the region is extended further upstream.}
performed 2D simulations with $\sigma = 1$ and $\Delta \gamma = 10^{-4}$ for pre-shock Lorentz factors up to $\gamma_0 = 80$ and we do not find that this changes the precursor dynamics or efficiency, in agreement with our earlier 1D results (see figure 8 in Margalit et al. 2020b).

The second panel in Fig. 8 shows the shock velocity in units of the speed of light, as measured in the downstream frame. Its dependence on dimensionality and temperature (in the regime $\Delta \gamma \lesssim 10^{-1}$ we have explored) is minimal. The marginal reduction at large $\Delta \gamma$ is expected based on MHD jump conditions. The fact that for $\sigma = 1$, 1D shocks in very cold plasmas are slightly slower than 2D shocks mirrors the different efficiency in precursor emission: in 1D the precursor takes away a larger fraction of the plasma energy, thereby slowing down the shock. The dependence on magnetization follows the expectation of MHD jump conditions. Assuming the adiabatic index of a relativistic gas with two degrees of freedom ($\gamma_{ad} = 3/2$; in fact, our particles only isotropize in the $xy$ plane orthogonal to the magnetic field), the dimensionless 4-velocity of ultra-relativistic magnetized (i.e., $\gamma_0 >> 1$ and $\sigma >> 1$) shocks is expected to be $\gamma_{ad} \beta_{sh} = (5\sigma/4 + 7/20)^{1/2}$ (e.g., Petri & Lyubarsky 2007; Plotnikov et al. 2018). At $\sigma >> 1$, this scales as $\gamma_{ad} \beta_{sh} \propto \sigma^{1/2}$. As we discuss in Section 5, this has important implications for the beaming of the precursor emission.

The data in the first and second panels are used to compute the wave efficiency $f_i$ shown in the third panel. This is defined as the fraction of incoming total energy (electromagnetic and kinetic) that is converted into precursor wave energy. In the downstream frame of the simulations, the energy that has flown into the shock per unit area up to time $t$ is

$$E_{in} = \gamma_0(1 + \sigma)N_0 m c^2 (\beta_0 + \beta_{sh})ct$$  \hspace{1cm} (5)

where the flux factor $(\beta_0 + \beta_{sh})$ accounts for the fact that the shock is moving towards the upstream. The energy converted into precursor waves per unit shock area is

$$E_{out} = \langle \delta B_x^2 \rangle_{4\pi} (1 - \beta_{sh})ct$$  \hspace{1cm} (6)

where we have assumed that the whole region between the shock and the leading edge of the precursor (moving at $c$) is occupied by precursor waves with uniform energy density. The efficiency is then

$$f_i = \frac{E_{out}}{E_{in}} = \xi_B \left( \frac{1}{\sigma} \frac{1 - \beta_{sh}}{\beta_0 + \beta_{sh}} \right)$$  \hspace{1cm} (7)

Notice that in the $\sigma >> 1$ limit this scales as $f_i \propto \xi_B \sigma^{-1}$. In the shock maser scenario for FRBs, this quantifies the fraction of blast wave energy that is converted into precursor wave energy (i.e., the candidate FRB). The ratio of precursor power to blast wave power can be obtained by accounting for the duration of the precursor emission, which in the limit $\sigma >> 1$ is a factor of $\sim \sigma^{-1}$ shorter than the blast wave ejection duration. It follows that in the limit of high magnetizations the ratio of emitted precursor power to blast wave power is $\sim \xi_B$.

The bottom panel of Fig. 8 shows that the precursor efficiency in 2D simulations scales as $f_i \sim 10^{-3}\sigma^{-1}$ as long as $\Delta \gamma \lesssim 10^{-15}$, whereas it drops abruptly to $f_i \lesssim 10^{-5}$ for $\Delta \gamma \gtrsim 10^{-4}$.

5 PRECURSOR BEAMING AND POWER SPECTRUM

To characterize the spectral properties of precursor waves, we have constructed the 2D wavenumber spectrum $P_{2D}(k_x, k_y)$ of $|\delta B_x(k_x, k_y)|^2$, by taking the Fourier transform $\delta B_x(k_x, k_y)$ of the fluctuating magnetic field $\delta B_x(x, y)$. Our spectra are computed in the post-shock frame, by extracting $\delta B_x(x, y)$ in the same region ahead of the shock ($5c/\omega_p < x - x_{sh} < 30c/\omega_p$) where we compute the precursor efficiency. We will show both the 2D power spectrum $P_{2D}(k_x, k_y)$ and the $k_x$-integrated 1D power spectrum $P(k_x) = \int P_{2D}(k_x, k_y)dk_y$. The power spectra are normalized such that $\int P_{2D}(k_x, k_y)dk_xdk_y = \int P(k_x)dk_x = \xi_B$.

In Plotnikov & Sironi (2019), we showed that for 1D simulations, wavenumber spectra ($k_x$-spectra) and frequency spectra ($\omega$-spectra) nearly overlap, when accounting for the dispersion relation of X modes.
In Figs. 9-11, we present the 2D power spectrum for three representative simulations, having different values of magnetization and pre-shock temperature. In Fig. 9 we show results for $\sigma = 1$ and $\Delta \gamma = 10^{-5}$ (the cold case described in Fig. 1 and Fig. 2), in Fig. 10 for $\sigma = 1$ and $\Delta \gamma = 10^{-7}$ (the warm case in Fig. 3 and Fig. 4), and in Fig. 11 for $\sigma = 3$ and $\Delta \gamma = 10^{-5}$ (a cold case with higher magnetization).

In each of the plots, the power at $k_x \sim 0$ and $k_y \gtrsim \omega_p/c$ is attributed to wave filamentation associated to the density filaments observed in panel (c) of Figs. 1 and 3. Most of the spectral power, however, resides at higher $k_x$, within the region delimited by the green dashed line (defined below). By comparing the figures, one sees that for cold plasmas (Fig. 9 and Fig. 11) the power is distributed over a wide range of longitudinal wavenumbers ($k_x \sim 5 - 10 \omega_p/c$), whereas for warm plasmas (Fig. 10) the spectrum is sharply peaked at $k_x \sim 5 \omega_p/c$. This will be further discussed below, where we show the $k_x$-integrated spectrum $P(k_x)$.

By comparing Fig. 9 and Fig. 11, which differ in magnetization, we find that the precursor emission is beamed within a narrower angle $\theta = \arctan(k_y/k_x)$ for larger $\sigma$. As we now discuss, this follows from the requirement that the waves be able to escape ahead of the shock, which moves faster for higher magnetizations (see middle panel of Fig. 8).

The dispersion relation of the extraordinary mode (X mode) in cold plasmas in the frame where the background plasma is at rest reads (see, e.g., Hoshino & Arons 1991)

$$k^2 c_s^2 = 1 - \frac{\omega_p^2}{\omega^2 - \sigma^2 \omega_p^2},$$

where double primed quantities are measured in the upstream rest frame. In the limit $\gamma_0^* \gg \sigma$, the dispersion relation in the downstream frame becomes

$$k^2 c_s^2 = \omega^2 - \omega_p^2,$$

which is identical to the dispersion relation of a simple electromagnetic wave propagating in an unmagnetized plasma.

The motion of the shock front imposes a cutoff below which the wave cannot escape into the upstream medium. This cutoff is obtained by equating the projection of the group velocity $v_g = \partial \omega/\partial k$ onto the shock normal with the shock speed, i.e.,

$$v_g \cos \theta \gtrsim \beta_{sh} c,$$

which leads to

$$k_x \gtrsim \gamma_0 \beta_{sh} \sqrt{k_Y^2 + \omega_p^2 / c^2}.$$  \hspace{1cm} (10)

This inequality identifies the values of $(k_x, k_y)$ for which the wave can successfully outrun the shock (Iwamoto et al. 2017).

This has two main consequences. First, the range of $k_x$ allowed for wave propagation has an absolute lower limit, $k_{x\text{cut}} = \gamma_0 \beta_{sh} \omega_p/c$. In the limit of high magnetizations, this scales as $k_{x\text{cut}} \simeq \sigma^{1/2} \omega_p/c$, in agreement with a comparison of Fig. 9 and Fig. 11. Second, for each given $k_y$, the allowed range of $k_x$ is constrained by Eq. (10), which is indicated by the dashed green lines in Figs. 9-11. At $k_x, k_y \gg \omega_p/c$ this corresponds to precursor emission being confined within an angle $\theta_{\text{crit}} = \arctan(1/\gamma_0 \beta_{sh})$ from the shock normal. For $\sigma \gg 1$, this scales as $\theta_{\text{crit}} \simeq \sigma^{-1/2}$. So, precursor waves from more strongly magnetized shocks will be directed closer to the shock normal.

This is demonstrated in Fig. 12, by computing the angular distribution of precursor power $dP_\theta/d\theta$, obtained by integrating the 2D wavenumber spectrum $P_\alpha(k_x, k_y)$ by integrating along lines of constant $\theta$. The two curves correspond to the two cases in Figs. 9 and 11 (see legend), and are normalized to their respective peak values. The vertical dashed lines (same color coding as the solid lines) correspond to $\theta_{\text{crit}} = \arctan(1/\gamma_0 \beta_{sh})$. The plot shows that the precursor emission at $\theta > \theta_{\text{crit}}$ is indeed negligible (the power at $\theta = \pi/2$ is contributed by the non-propagating filamentation mode, as described above). As expected, precursor waves in flows with higher magnetizations...
and varying temperatures is presented in Fig. 13, where we show the \( k_2 \)-integrated wavenumber spectrum \( P(k_2) \) (more precisely, we show \( k_2 P(k_2) \)) to emphasize where most of the power resides.

We focus our discussion on the range \( k_s < 20 \omega_p/c \) which is robust against variations of numerical parameters, see Appendix A.

Once the precursor efficiency settles to a steady state (see Fig. 7), the spectral shape is also nearly time independent. Regardless of the pre-shock temperature, the spectral shape presents some common features: (i) the range of longitudinal wavenumbers has a sharp cutoff at \( k_{\text{x,cut}} \approx 2 \omega_p/c \), which descends from the constraint of Eq. (10) for \( k_s = 0 \); (ii) the spectral shape at high wavenumbers \( (k_s \geq 5 \omega_p/c) \) resembles a power law \( P(k_s) \propto k_s^{-2} \).

Despite these similarities, sharp differences exist between cases with different pre-shock temperatures. For cold temperatures, the spectrum peaks at \( k_s \sim 5 - 10 \omega_p/c \) and is relatively broad, with fractional width \( \sim 1 - 3 \). This is common to all cases with \( \Delta \gamma \leq 10^{-3} \) (we show \( \Delta \gamma = 10^{-5} \) in black and \( \Delta \gamma = 10^{-4} \) in blue). In this temperature range, the spectral power at high wavenumbers \( (k_s \geq 10 \omega_p/c) \) gets increasingly suppressed for larger thermal spreads, as discussed by Amato & Arons (2006). For warm temperatures \( (10^{-3} \leq \Delta \gamma \leq 10^{-1.5}) \), the spectrum shows pronounced line-like features with fractional width \( \sim 0.2 \). The line-like features are located at the low-wavenumber end of the spectrum, at \( k_s \sim 3 - 5 \omega_p/c \). We show spectra for \( \Delta \gamma = 10^{-3} \) (green), \( \Delta \gamma = 10^{-2.5} \) (yellow) and \( \Delta \gamma = 10^{-2} \) (red). The spectrum for \( \Delta \gamma = 10^{-1.5} \) (not shown) is very similar to the \( \Delta \gamma = 10^{-2} \) case, whereas we remind that the precursor efficiency is strongly suppressed for even hotter temperatures \( (\Delta \gamma \gtrsim 10^{-1}) \).

6 SUMMARY AND DISCUSSION

In this work we have investigated by means of 2D PIC simulations the physics of the precursor waves emitted by perpendicular relativistic electron-positron shocks with out-of-plane upstream fields. We have focused on the high magnetization regime \( \gamma \geq 1 \) appropriate for magnetar winds, motivated by the shock-powered synchrotron maser scenario proposed for FRBs (Lyubarsky 2014; Murase et al. 2016; Beloborodov 2017; Waxman et al. 2017; Metzger et al. 2019; Beloborodov 2019; Margalit et al. 2020b,a). We have explored the efficiency and spectrum of the precursor waves as a function of the pre-shock thermal spread \( \Delta \gamma = kT/\rho c^2 \) in the range \( \Delta \gamma = 10^{-3} - 10^{-1} \). All our simulations have been run for a sufficiently long time \( (\sim 4000 \omega_p^{-1}) \) that the precursor emission has achieved a steady state. Our main results are:

(i) By measuring the fraction \( f_p \) of total (i.e., electromagnetic and kinetic) incoming energy that is converted into precursor waves, as computed in the post-shock frame, we can quantify the efficiency of precursor emission. At fixed temperature, the scaling with magnetization \( f_p \propto \gamma^{-1} \) at \( \gamma \gtrsim 1 \) is consistent with our earlier 1D results (Plotnikov & Sironi 2019).

(ii) At fixed magnetization, the precursor efficiency is nearly independent of temperature as long as \( \Delta \gamma \leq 10^{-3} \) (with only a modest decrease of a factor of three from \( \Delta \gamma = 10^{-5} \) to \( \Delta \gamma = 10^{-1.5} \)), but between \( \Delta \gamma = 10^{-1.5} \) and \( \Delta \gamma = 10^{-1} \) it drops by nearly two orders of magnitude. We have confirmed this result with dedicated 3D simulations (Sironi et al., in prep.) So, shocks propagating in hot plasmas with \( \Delta \gamma \gtrsim 10^{-1} \) are unlikely to power FRBs.

(iii) For \( \gamma \gtrsim 1 \), the precursor waves are beamed within a cone of half-opening angle \( \theta_{\text{out}} = \arctan(1/\gamma a \beta_0) \approx \gamma^{-1/2} \) around the shock normal (as measured in the post-shock frame). This stems from the fact that only the waves whose group velocity projected along the shock normal is larger than the shock speed can outrun the shock. More precisely, the width at half maximum of the angular distribution of precursor power is \( \sim 0.7 \theta_{\text{out}} \).

(iv) For \( \gamma = 1 \), we have compared the power spectrum \( P(k_s) \) of precursor waves (integrated over the transverse wavenumber \( k_t \)) among different \( \Delta \gamma \). For cold temperatures, the spectrum peaks at \( k_s \sim 5 - 10 \omega_p/c \) and is relatively broad, with fractional width \( \sim 1 - 3 \). In contrast, for warm temperatures \( (10^{-3} \leq \Delta \gamma < 10^{-1}) \) it shows pronounced line-like features with fractional width \( \sim 0.2 \). The line-like features are located at the low-wavenumber end of the spectrum, at \( k_s \sim 3 - 5 \omega_p/c \). For both cold and warm flows, the high-wavenumber part at \( k_s > 10 \omega_p/c \) can be roughly modeled as a power law \( P(k_s) \propto k_s^{-2} \).

Our simulations employ 2D computational domains initialized with out-of-plane magnetic fields. This configuration only allows for the excitation of the X mode (and not of the O mode), so in our case the precursor waves are 100% linearly polarized, with fluctuating magnetic field along the same direction as the upstream mean field. A 3D investigation of the dependence of the precursor emission (in terms of both O and X modes) on the pre-shock temperature will be presented elsewhere.

The shocks investigated in this work propagate in an electron-positron plasma. The physics of electron-proton shocks will quali-
Here, “longitudinal” is along the shock direction of propagation (x direction), while “transverse” is along the wave electric field (y direction).

5 Here, “longitudinal” is along the shock direction of propagation (x direction), while “transverse” is along the wave electric field (y direction).
Figure A1. 1D wavenumber power spectrum \( P(k_x) = \int P_{2D}(k_x, k_y) \, dk_y \) for 2D simulations with \( \sigma = 1 \) and different choices of particles per cell \( N_0 \) and number of passes of the smoothing filter of electric currents \( N_{\text{sm}} \) (see legend). The spectrum is extracted from the same region ahead of the shock where \( \xi_{\text{sh}} \) is computed, i.e. \( 5 \, c/\omega_p < x < x_{\text{sh}} < 30 \, c/\omega_p \). At each time, the spectrum is normalized such that its integral equals \( \xi_{\text{sh}} \). Here, each curve is obtained by averaging the spectra in the time interval \( 2000 < \omega_p \tau < 2500 \).

The vertical dashed line indicates the boundary \( k_{\text{max}} = 20 \, \omega_p/c \) beyond which our spectra differ because of numerical effects.

APPENDIX B: EFFECT OF LONGITUDINAL AND TRANSVERSE DISPERSION

In this Appendix, we present dedicated 1D simulations to clarify whether it is the longitudinal or the transverse dispersion that is most detrimental for the efiiciency of the precursor emission. We employ \( N_0 = 400 \) particles per cell, a spatial resolution of \( c/\omega_p = 112 \) cells and a numerical speed of light of 0.5 cells/timestep. The magnetic field initialization is the same as in the 2D simulations presented in the main body of the paper, and we focus on \( \sigma = 1 \).

Fig. B1 shows that the efficiency drops by nearly two orders of magnitude from \( \Delta y = 10^{-6} \) (black) to \( \Delta y = 10^{-8.5} \) (red). Indeed, as Fig. 8 shows, the drop occurs between \( \Delta y = 10^{-5.5} \) and \( \Delta y = 10^{-4} \). We perform two additional simulations, in which we start with a hot upstream flow \( (\Delta y = 10^{-8.5}) \); as for the red line), but for \( \omega_p \tau \gtrsim 700 \) we artificially change its properties right ahead of the shock, in the region \( 4 \, c/\omega_p < x < x_{\text{sh}} < 22 \, c/\omega_p \). For the yellow line, we suppress the transverse momentum dispersion (i.e., along \( y \)), while for the green line we suppress the longitudinal momentum dispersion (i.e., along \( x \)). If the region where we enforce the suppression were to be far ahead of the shock, gyration around the upstream magnetic field would interchange the transverse and longitudinal motions on a timescale \( (\sigma/2) \gamma_0 \sigma^{-1} \omega_p^{-1} \), or equivalently on a distance \( (\sigma/2) \gamma_0 \sigma^{-1} \omega_p^{-1} \). For \( \gamma_0 = 10 \) and \( \sigma = 1 \), the choice of suppressing one momentum component between 4 and 22 skin depths is then sufficient to guarantee that the flow entering the shock preserves its imposed temperature anisotropy.

Fig. B1 shows that, if we suppress the transverse dispersion (yellow line), the shock retains the low efficiency of the hot case (red curve). In contrast, if we suppress the longitudinal dispersion (green line), the efficiency increases and eventually settles to the level corresponding to the cold case (black line). We have also verified that at late times the power spectra corresponding to the black and green lines are similar (and the same applies to the spectral comparison of yellow and red lines).

This demonstrates that, if the upstream plasma were to be continuously heated in an anisotropic way, it is the longitudinal dispersion that determines the efficiency of the precursor emission. Such anisotropic heating may result from the propagation of the precursor waves themselves, since they preferentially induce particle motions in the transverse direction.

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Figure B1. Time evolution of the normalized precursor wave energy, $\xi_B = \langle \delta B_z^2 \rangle / B_0^2$ for 1D simulations with $\sigma = 1$. The black and red lines respectively correspond to a cold case ($\Delta \gamma = 10^{-6}$) and a hot case ($\Delta \gamma = 10^{-5.5}$). The yellow and green lines respectively correspond to cases where we start with a hot thermal spread ($\Delta \gamma = 10^{-5.5}$), but for $\omega_p t > 700$ we suppress by hand either the longitudinal momentum dispersion (i.e., along $x$) or the transverse one (i.e., along $y$) just ahead of the shock. Here, “longitudinal” and “transverse” refer to the shock direction of propagation.

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