Precursor Wave Emission Enhanced by Weibel Instability in Relativistic Shocks

Masanori Iwamoto1, Takanobu Amano1, Masahiro Hoshino1, and Yosuke Matsumoto2

1 Department of Earth and Planetary Science, University of Tokyo, 7-3-1 Hongo, Bunkyo-ku, Tokyo 113-0033, Japan; iwamoto@eps.s.u-tokyo.ac.jp
2 Department of Physics, Chiba University, 1-33 Yayoi, Inage-ku, Chiba, Chiba 263-8522, Japan

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Abstract

We investigated the precursor wave emission efficiency in magnetized purely perpendicular relativistic shocks in pair plasmas. We extended our previous study to include the dependence of upstream magnetic field orientations. We performed two-dimensional particle-in-cell simulations and focused on two magnetic field orientations: the magnetic field in the simulation plane (i.e., in-plane configuration) and that perpendicular to the simulation plane (i.e., out-of-plane configuration). Our simulations in the in-plane configuration demonstrated that not only extraordinary but also ordinary mode waves are excited. We quantified the emission efficiency as a function of the magnetization parameter \( \sigma_e \), and found that the large-amplitude precursor waves are emitted for a wide range of \( \sigma_e \). We found that especially at low \( \sigma_e \), the magnetic field generated by Weibel instability amplifies the ordinary mode wave power. The amplitude is large enough to perturb the upstream plasma, and transverse density filaments are generated as in the case of the out-of-plane configuration investigated in the previous study. We confirmed that our previous conclusion holds regardless of upstream magnetic field orientations with respect to the two-dimensional simulation plane. We discuss the precursor wave emission in three dimensions and the feasibility of wakefield acceleration in relativistic shocks based on our results.

Key words: acceleration of particles – cosmic rays – plasmas – shock waves

1. Introduction

Observations of active galactic nuclei (AGNs) and gamma-ray bursts (GRBs) usually show broad nonthermal spectra (e.g., Kaneko et al. 2006; Abdo et al. 2010), which are believed to originate from synchrotron radiation and inverse Compton scattering of relativistic electrons. Since the relativistic outflow from the central compact object is the common feature in AGNs and GRBs (e.g., Gehrels et al. 2009; Lister et al. 2016), relativistic shocks can be formed upon interaction between the jets and the interstellar medium. The relativistic shocks are assumed to play an important role for generating such nonthermal electrons.

Previous one-dimensional (1D) particle-in-cell (PIC) simulations showed that synchrotron maser instability (SMI) is the significant dissipation mechanism for relativistic magnetized shocks (e.g., Langdon et al. 1988; Gallant et al. 1992; Hoshino et al. 1992; Amato & Arons 2006). The SMI is driven by particles reflected off the shock-compressed magnetic field in the shock-transition region and emits electromagnetic waves of extraordinary mode (\( X \)-mode) both upstream and downstream (Hoshino & Arons 1991). Since the electromagnetic precursor waves have a non-negligible fraction of the upstream kinetic energy, the upstream flow is significantly perturbed by the precursor wave (Lyubarsky 2006). Hoshino (2008) demonstrated that the wave power is strong enough to induce wakefield in the upstream and that nonthermal electrons are generated by wakefield acceleration (WFA; Tajima & Dawson 1979; Chen et al. 2002) in 1D relativistic shocks propagating in magnetized ion–electron plasmas.

In multidimensional systems, it is well known that Weibel instability (WI; Fried 1959; Weibel 1959) develops in the transition region of weakly magnetized shocks. The WI is widely studied in laser plasma as well as astrophysics (e.g., D’Angelo et al. 2015; Huntington et al. 2015, 2017; Park et al. 2015). Previous PIC simulation studies in multiple dimensions demonstrated that the WI grows into substantial amplitude in the shock-transition region at low magnetization \( \sigma_e = \omega_{ce}^2/\omega_{pe}^2 < 10^{-2} \) (e.g., Spitkovsky 2005; Sironi et al. 2013). Here, \( \omega_{ce} \) is the relativistic electron cyclotron frequency and \( \omega_{pe} \) is the proper electron plasma frequency. The effective temperature anisotropy induced by reflected particles in the shock-transition region provides the free energy source for the development of the WI (e.g., Kato 2007; Chang et al. 2008). The linear theory, including relativistic effects showed that the maximum growth rate of the WI is on the order of \( \omega_{pe} \) (see, e.g., Yang et al. 1993; Achterberg et al. 2007; Schueber-Rolffs & Tautz 2008), whereas that of the SMI is on the order of \( \omega_{ce} \) (Hoshino & Arons 1991). Since both instabilities are excited from the same free energy source in the same region and \( \omega_{pe} \) is much greater than \( \omega_{ce} \) for \( \sigma_e < 1 \), it was believed that the WI dominates over the SMI and the precursor wave emission could be shut off in multidimensional shocks.

Recently, by using two-dimensional (2D) PIC simulations, we have shown that the SMI can coexist with the WI and that the precursor wave emission continues to persist even in the Weibel-dominated regime (Iwamoto et al. 2017). We also showed that the wave power is sufficient enough to induce wakefield for a wide range of magnetization parameter \( \sigma_e \). Based on the results, we suggested that external shocks in the relativistic jets from GRBs may be important sites for the production of ultra-high-energy cosmic rays via WFA.

However, in the previous work, we focused only on perpendicular shock with the upstream ambient magnetic field perpendicular to the simulation plane (i.e., the out-of-plane configuration). One may also choose the upstream ambient magnetic field to be in the simulation plane (i.e., in-plane configuration), which may, in general, change the shock dissipation physics because the degree of freedom in this case becomes three rather than two in the out-of-plane configuration (e.g., Amano & Hoshino 2009). In fact, Sironi et al. (2013) reported that the particle acceleration efficiency in 2D
perpendicular shocks depends on the orientation of the pre-shock magnetic field. Therefore, in this study, we consider the in-plane configuration and investigate the physics of magnetized perpendicular shocks, especially the electromagnetic wave emission by the SMI. We quantify the precursor wave emission efficiency and discuss the effects of the magnetic field configuration on the WFA, combining this study with our previous results.

This paper is organized as follows. First, Section 2 describes our simulation setup. In Section 3, we show the global structure of relativistic magnetized shocks for relatively high and low magnetization, respectively. In Section 4, the properties of precursor waves are analyzed. In Section 5, we discuss the wave excitation mechanism and the feasibility of the WFA in relativistic magnetized shocks. Finally, Section 6 summarizes this study.

2. Simulation Setup

We carried out simulations of 2D perpendicular shocks in electron–positron plasmas using an electromagnetic PIC code (Matsumoto et al. 2013, 2015). The basic configuration of our simulations is almost identical to our previous simulation (Iwamoto et al. 2017) and schematically illustrated in Figure 1. We changed only the direction of the upstream ambient magnetic field \( B_1 \) from the out-of-plane direction \( (\zeta \text{ direction in our coordinate system}) \) to the in-plane direction \( (y \text{ direction}) \).

Our simulation domain is in the \( x-y \) plane with periodic boundary conditions in the \( y \) direction and the number of grids in each direction is \( N_x \times N_y = 20,000 \times 1,680 \). A cold pair stream is continuously injected along the \( -x \) direction with a bulk Lorentz factor \( \gamma_1 = 40 \) from the right-hand boundary and elastically reflected at the left-hand boundary. The shock wave is excited by the interaction between the returning particles and the incoming plasma flow and propagates toward \( +x \) direction. The number of particles per cell in the upstream is \( N_e \Delta x^2 = 64 \) for both electrons and positrons, where \( \Delta x \) is the grid size. The grid size is fixed to \( \Delta x/(c/\omega_{pe}) = 1/40 \) throughout this study, where \( c \) is the speed of light and the \( \omega_{pe} \) is the proper electron plasma frequency.

The proper electron plasma frequency is defined as follows:

\[
\omega_{pe} = \sqrt{\frac{4\pi N_e e^2}{\gamma_1 m_e}}. \tag{1}
\]

The number of particles per cell and the grid size are motivated by the numerical convergence study of 1D simulations (see Iwamoto et al. 2017, Appendix A). The time step is set to \( \omega_{pe}\Delta t = 1/40 \) in order to minimize the effect of the numerical Cherenkov instability (Ikeya & Matsumoto 2015). For more details, please refer to our previous paper (Iwamoto et al. 2017).

As in our previous study, we investigated the dependence of the precursor wave emission on the magnetization parameter \( \sigma_e \):

\[
\sigma_e = \frac{B_1^2}{4\pi \gamma_1 N_e m_e c^2} \approx \frac{\omega_{ce}^2}{\omega_{pe}^2}. \tag{2}
\]

where \( \omega_{ce} \) is the relativistic electron cyclotron frequency:

\[
\omega_{ce} = \frac{eB_1}{\gamma_1 m_e c}. \tag{3}
\]

More specifically, we discuss the results obtained from the following eight runs: \( \sigma_e = 1, 3 \times 10^{-1}, 3 \times 10^{-2}, 1 \times 10^{-2}, 3 \times 10^{-3}, 1 \times 10^{-3}, \) and \( 3 \times 10^{-4} \).

3. Global Shock Structure

3.1. High-\( \sigma_e \) Case

First, we discuss the overview of the global shock structure for relatively high \( \sigma_e \). Figure 2 is the global shock structure at \( \omega_{pe}l = 500 \) for \( \sigma_e = 3 \times 10^{-1} \). The electron number density \( N_e \), the electron number density averaged along the \( y \) axis \( \langle N_e \rangle \), the \( x \) component of the magnetic field \( B_x \), 1D cut along \( y = 21c/\omega_{pe} \) for \( B_x \), the in-plane magnetic field \( B_{xy} \), 1D cut along \( y = 21c/\omega_{pe} \) for \( B_{xy} \), the out-of-plane magnetic field \( B_{z} \), 1D cut along \( y/(c/\omega_{pe}) = 21 \) for \( B_z \), and the electron phase-space density \( x-u_{xe}, x-u_{ye}, \) and \( x-u_{ce} \) integrated over the \( y \) direction are shown from top to bottom. All quantities are normalized by the corresponding upstream values. Note that our 2D simulations track all three components of the particle velocity and electromagnetic field. A well-developed shock structure is formed at this time, and the shock front is clearly seen at \( x/(c/\omega_{pe}) \sim 235 \).

At the shock front, fluctuations in \( B_z \) are generated. We think the magnetic field fluctuations may be attributed to instabilities excited in the shock-transition region. One of the possible instabilities for this case is the Alfvén-ion-cyclotron instability, which is an electromagnetic instability on the Alfvén mode branch driven by a temperature anisotropy (e.g., Winske & Quest 1988). We perform linear analysis for a relativistic pair plasma with a cold ring distribution and indeed find a similar instability. This instability may be the cause of fluctuations in \( B_z \) and the magnetic field energy is eventually amplified up to 10%–20% of the upstream kinetic energy. Although fluctuations in \( B_z \) at the shock front may also be generated by the instability, the fluctuations start decreasing in time after \( \omega_{pe}l \sim 140 \) and are not clearly seen at this time. The wave magnetic fields \( \delta B_{e} \) are visible in the upstream region. The electromagnetic waves are continuously emitted from the shock front and persist with large amplitude. Remember that the upstream ambient magnetic field is in the \( y \) direction. The wave magnetic field is polarized in the \( y \) direction and parallel to the ambient magnetic field, which is the signature of the X-mode wave (see Section 4.1). This result is consistent with both the linear theory (Hoshino & Arons 1991) and the previous 2D simulation (Iwamoto et al. 2017). The oblique propagation of these X-mode waves may be responsible for the
x component of the fluctuating magnetic field $\delta B_x$ in the upstream region. Since $\delta B_x$ is very small compared to $\delta B_y$, we mainly consider $\delta B_y$ in our analysis. We think that the waves in the region $x/(c/\omega_{pe}) \gtrsim 460$ are contaminated by the initial and boundary conditions. Therefore, we excluded this region from our analysis presented below.

The wave magnetic fields $\delta B_z$ are identified in the upstream region. They also appear to be electromagnetic precursor waves emitted from the shock front. The wave magnetic field is polarized in the $z$ direction, and thus the wave mode is the ordinary mode (O-mode; see Section 4.1). This is unexpected because the linear theory of the SMI showed that the growth rate of the O-mode is finite at oblique propagation but much smaller than that of the X-mode (see, e.g., Wu & Lee 1979; Lee et al. 1980; Melrose et al. 1982, 1984). The amplitude of the O-mode wave is smaller than that of the X-mode wave but non-negligible. The tip of the O-mode wave is behind that of the X-mode wave. This delay should result from the difference of the generation time since both have group velocities almost equal to the speed of light. The X-mode waves are generated by
the SMI soon after the shock formation in the initial phase of the simulation. In contrast, the generation of the O-mode waves seems to become effective after \( \omega_{pe}f \sim 80 \), which is estimated from the time evolution of the wave magnetic field \( \delta B_y \). We discuss how the O-mode waves are excited in Section 5.1.

As in the case of our previous simulation, transverse density filaments are formed in the upstream region. This again indicates that the precursor waves remain large amplitude and coherent in 2D systems.

3.2. Low-\( \sigma_e \) Case

Here we discuss the overall shock structure for relatively low \( \sigma_e \). Figure 3 is the global shock structure at \( \omega_{pe}f = 500 \) for \( \sigma_e = 3 \times 10^{-3} \). The format is the same as that of Figure 2. A well-developed shock front is distinctly visible at \( x/(c/\omega_{pe}) \sim 160 \).

The filamentary magnetic field, expected for the structure of the Weibel-generated magnetic field, is seen at the shock front in the \( x \) and \( z \) direction. In the in-plane configuration, previous works indeed showed that the WI excites \( B_z \) as well as \( B_x \). (e.g., Matsukiyo & Scholer 2006). Our linear analysis arrives at the same conclusion (see Appendix A). Thus we think that the fluctuations in \( B_x \) and \( B_z \) near the shock front is attributed to the WI. The maximum magnetic field energies for both components reach about 10%–20% of the upstream kinetic energy, which is consistent with the previous studies (Kato 2007; Chang et al. 2008; Sironi & Spitkovsky 2011). However, the length of the Weibel filaments are shorter than that in the out-of-plane configuration (see Iwamoto et al. 2017, Figure 3). As we already explained, the WI is driven to be unstable by the effective field diffusion in the in-plane configuration.

The suppression of the cross-field diffusion in the in-plane configuration may also contribute to the relatively short reflected particle beam. Jokipii et al. (1993) and Jones et al. (1998) mathematically proved that charged particles cannot move further than one Larmor radius from a given magnetic field if there are one or more ignorable coordinates. A notable exception is a 2D system with the out-of-plane magnetic field, which thus allows particles to diffuse across the magnetic field. In contrast, the diffusion of particles back into the upstream is prohibited in the in-plane configuration. Therefore, the length of the reflected particle beam in the in-plane case may become shorter than that in the out-of-plane case.

The precursor waves are observed both in \( B_z \) and \( B_x \), and the delay of the O-mode precursor wave is identified in this case as well. The generation time of the O-mode wave may be estimated to be \( \omega_{pe}f \sim 40 \). The amplitude of the O-mode wave is comparable to that of the X-mode wave unlike the high-\( \sigma_e \) case. The \( \sigma_e \) dependence is discussed in Section 4.4 in more detail. Notice that clear density filaments are observed in the precursor region in this case as well.

4. Precursor Wave

4.1. Wave Mode

As we mentioned in Section 3, the X-mode and O-mode electromagnetic waves are observed in the in-plane configuration. Both of the waves propagate perpendicular to the ambient magnetic field and are linearly polarized in pair plasmas. The wave magnetic field of the X-mode is parallel to the ambient magnetic field, whereas that of the O-mode is perpendicular.

Figure 4 is the enlarged view of the region in \( 300 \leq x/(c/\omega_{pe}) \leq 320 \) for \( \sigma_e = 3 \times 10^{-1} \) (left) and \( 200 \leq x/(c/\omega_{pe}) \leq 220 \) for \( \sigma_e = 3 \times 10^{-3} \) (right), and shows the \( y \) and \( z \) components of the wave electromagnetic fields at \( \omega_{pe}f = 500 \). The electromagnetic fields are normalized by the upstream ambient magnetic field \( B_1 \). The top panels show the \( y \) component of the wave magnetic field \( \delta B_y \) and the \( z \) component of the wave electric field \( \delta E_z \). The red and blue lines indicate the magnetic field and electric field, respectively. Recall that the upstream ambient magnetic field \( B_1 \) is oriented along the \( y \) axis. The anticorrelation between \( \delta B_z \) and \( \delta E_z \) in phase are clearly seen in both cases, and the amplitude of the magnetic field is almost identical to that of the electric field. It is easy to confirm that the waves carry the positive Poynting flux, indicating that the waves propagate toward the \(+x\) direction. All of these results show that the X-mode and O-mode electromagnetic waves travel upstream.

4.2. Time Evolution

Now we discuss time evolution of the precursor wave power. Figure 5 shows the time evolution of the wave energy from \( \omega_{pe}f = 300 \) up to \( \omega_{pe}f = 500 \) for \( \sigma_e = 3 \times 10^{-1} \) (left) and \( \sigma_e = 3 \times 10^{-3} \) (right). The time evolution is determined by the same method as that in our previous study (Iwamoto et al. 2017). The wave energy is given in units of the upstream bulk kinetic energy, and \( y \) and \( z \) components are shown by the solid and dashed lines, respectively. As shown in Figure 4, the amplitude of the electric field is comparable to that of the magnetic field. Thus the same plots for the electric field is almost identical and we show only those for the magnetic field.

For \( \sigma_e = 3 \times 10^{-1} \), although \( \delta B_y \) gradually declines in time, it still remains finite and gets saturated at around \( \omega_{pe}f = 460 \). In contrast, \( \delta B_z \) shows a continuous decrease. Although the O-mode wave emission might be shut off after long-term evolution, the X-mode wave emission has already reached a quasi-steady state by the end of our simulation and the wave amplitude is comparable to that in the out-of-plane configuration (see Section 4.4). Therefore, the coherent electromagnetic precursor wave emission continues in the in-plane as well as out-of-plane configuration.

For \( \sigma_e = 3 \times 10^{-3} \), both \( \delta B_x \) and \( \delta B_y \) are already saturated in this time range. Considering that \( \delta B_y \) is the component expected from the linear theory of the SMI, it is somewhat surprising that \( \delta B_z \) is always greater than \( \delta B_y \). We discuss the \( \sigma_e \) dependence in Section 4.4 in detail.
4.3. Wavenumber Spectra

Figure 6 shows the precursor wave power spectra for each component in wavenumber space normalized by the upstream ambient magnetic field energy density. The left column shows the spectra of $\delta B_y$ (top) and $\delta B_z$ (bottom) for $\sigma_e = 3 \times 10^{-1}$, whereas the right column shows the spectra of $\delta B_y$ (top) and $\delta B_z$ (bottom) for $\sigma_e = 3 \times 10^{-3}$. The spectra are obtained in the same manner as in our previous study (Iwamoto et al. 2017). Note that the Nyquist wavenumber for our simulation is $k_x c / \omega_{pe} \approx 120$ and both X-mode ($\delta B_y$) and O-mode ($\delta B_z$) precursor waves are well resolved.

The white solid line indicates the theoretical cutoff wavenumber:

$$k_x = \beta_{sh} \gamma_{sh} \sqrt{k_x^2 \gamma_{sh} + \frac{2 \omega_{pe}^2}{c^2}},$$  \hspace{1cm} (4)

where $\beta_{sh}$ is the shock velocity normalized by the speed of light and $\gamma_{sh}$ is the Lorentz factor of the shock velocity (see Iwamoto et al. 2017, Appendix B). This theoretical cutoff wavenumber comes from the wavenumber below which the group velocity of the precursor wave is smaller than the shock velocity. Therefore, only those waves with $k_x$ greater than the threshold...
can escape from the shock in the upstream direction. The dispersion relation of the X-mode in a cold magnetized pair plasma is used to derive Equation (4). For $\gamma_1$ and $\sigma_e$ used in our simulation, the dispersion relation in the simulation frame can be written as

$$\omega^2 \simeq 2\omega_{pe}^2 + k^2c^2.$$  

(5)

This dispersion relation is identical to that of the O-mode in a cold pair plasma, and we use Equation (4) for the O-mode wave as well. The shock propagation velocity is determined from the time evolution of the $y$-averaged electron number density $\langle N_e \rangle$, which is then used for calculation of the theoretical cutoff wavenumber. The result shows that the precursor waves are indeed propagating away from the shock, suggesting that they are generated at the shock front.

4.4. $\sigma_e$ Dependence

Now we discuss the $\sigma_e$ dependence of the precursor wave amplitude. The wave amplitude was calculated by integrating the power spectra (Figure 6) over the whole wavenumber space. Figure 7 shows the precursor wave energy as a function of $\sigma_e$ normalized by the upstream ambient magnetic field energy (left) and the upstream bulk kinetic energy (right). The latter may be understood as the energy conversion rate from the upstream bulk kinetic energy to the precursor wave energy. Red, blue, and magenta indicate X-mode wave energy $\delta B_x^2$, O-mode wave energy $\delta B_y^2$, and total wave energy $\delta B_x^2 + \delta B_y^2$, respectively. The simulation results in the out-of-plane configuration by Iwamoto et al. (2017) are also shown in green for comparison. Note that only the X-mode precursor waves ($\delta B_x$) are excited in the out-of-plane configuration.

For $\sigma_e \gtrsim 10^{-2}$, the X-mode wave energy $\delta B_x^2$ in the in-plane configuration shows the same tendency as that in the out-of-plane configuration. This may be understood as follows. The ambient magnetic field is larger than the magnetic field fluctuations generated by the instability in the shock-transition region and almost unperturbed for high $\sigma_e$. Thus the X-mode wave excitation via the SMI is nearly identical between the in-plane and out-of-plane configurations.

For $\sigma_e \lesssim 10^{-2}$, the X-mode wave energy $\delta B_x^2$ in the in-plane configuration is greater than that in the out-of-plane configuration. This may be explained in terms of the coherence of the particle gyromotion in the shock-transition region. The WI generates strong magnetic field fluctuations for low $\sigma_e$ and the shock-transition region is dominated by the Weibel-generated magnetic field in both of the configurations. While charged particles, on average, gyrate in the $x-z$ plane for the in-plane configuration, they always gyrate in the $x-y$ plane for the out-of-plane configuration. Since the $z$ direction is ignored in our 2D simulations, the particle gyromotion in the in-plane case is less perturbed by the Weibel-generated turbulence than that in the out-of-plane case. In the in-plane case, therefore, the electromagnetic wave emission may be sufficiently amplified by the SMI and the wave amplitude may grow larger than that in the out-of-plane case.

The O-mode wave energy $\delta B_y^2$ is smaller than the X-mode wave energy $\delta B_x^2$ for $\sigma_e \gtrsim 10^{-2}$, whereas it exceeds the X-mode for $\sigma_e \lesssim 10^{-2}$. This tendency cannot be explained by the above argument. We discuss a possible excitation mechanism of O-mode waves and its relation to the $\sigma_e$ dependence in Section 5.1 in detail.

In conclusion, the simulation results have demonstrated that regardless of the orientation of the upstream ambient magnetic field, the precursor waves remain finite amplitude and coherent in 2D. This is true even for relatively low $\sigma_e$ cases where the WI grows into substantial amplitude in the shock-transition region. The results confirm the idea that the coherent electromagnetic precursor wave emission is the real nature of the relativistic magnetized shocks.

5. Discussion

5.1. Excitation Mechanism of O-mode Waves

Our simulation results show that the O-mode as well as X-mode electromagnetic waves are excited in relativistic shocks. As we already mentioned, the linear theory of the SMI (Wu & Lee 1979; Lee et al. 1980; Melrose et al. 1982, 1984) predicts that the X-mode wave emission overwhelms the O-mode wave emission. However, in the in-plane configuration, the O-mode precursor wave is clearly identified in the upstream region. This may be explained qualitatively as follows. The WI grows into substantial amplitude in the shock-transition region. The results confirm the idea that the coherent electromagnetic precursor wave emission is the real nature of the relativistic magnetized shocks.
magnetic field parallel to the ambient magnetic field, only $\delta B_y$ is excited in this stage. When the fluctuations in $B_z$ generated by the plasma instabilities in the shock-transition region have grown to be a non-negligible fraction of the ambient magnetic field, the net ambient magnetic field in the shock-transition region is undulated in the $y$-$z$ plane. If the SMI is induced by particles gyrating around the net ambient magnetic field, the X-mode wave in the shock-transition region will have $\delta B_t$, as well as $\delta B_y$. Such an X-mode wave experiences changes in the direction of the ambient magnetic field during its propagation upstream. If the polarization of the wave electromagnetic field remains unchanged during the propagation, the X-mode wave may be mode-converted into an O-mode wave when it reaches upstream. By performing simple PIC simulations, we have confirmed that this hypothesis is indeed correct. That is, an X-mode wave component keeps its polarization and is converted into an O-mode as it propagates through a layer of magnetic field rotation. Therefore, we believe that O-mode waves observed in the precursor region are the result of mode conversion from the X-mode generated by the SMI in the turbulent shock-transition region.

The delay of the O-mode wave in our simulation provides indirect evidence for this model. The excitation mechanism shows that the O-mode wave is excited after the generation of the strong magnetic field fluctuations by the instabilities in the shock-transition region. In fact, the generation time of the O-mode waves ($\omega_{pe} t \sim 80$ for $\sigma_e = 3 \times 10^{-1}$ and $\omega_{pe} t \sim 40$ for $\sigma_e = 3 \times 10^{-3}$) is roughly identical to the saturation time of the plasma instabilities (see Appendix A).

The $\sigma_e$ dependence of the O-mode wave amplitude in Figure 7 may be explained by this excitation mechanism. For $\sigma_e \gtrsim 10^{-2}$, the ambient magnetic field is much larger than the magnetic field fluctuations and almost aligned in the $y$ direction. Thus $\delta B_y$ should be the main component of the X-mode wave in the shock-transition region, and $\delta B_z$, which is observed as the O-mode wave in the upstream region may be much smaller for high $\sigma_e$.

For $\sigma_e \lesssim 10^{-2}$, the Weibel-generated magnetic field dominates over the ambient magnetic field and the effective $\sigma_e$ becomes much larger in the shock-transition region. The higher effective $\sigma_e$ allows a wave generated via a lower-order cyclotron harmonic resonance $n$ to satisfy the condition $\omega = n\omega_{ce} \gtrsim \sqrt{2(1 + \beta_{sh}/\gamma_{sh})}\omega_{pe}$ such that it can propagate upstream (see Iwamoto et al. 2017). We think that the lower cyclotron harmonics contribution may be the reason for the enhanced power of $\delta B_z$. In other words, the Weibel-generated magnetic field plays the role for the enhanced O-mode wave power. This model indicates that the O-mode wave continues to exist with a finite amplitude for considerably lower $\sigma_e$.

5.2. Implication for 3D

Based on the 2D simulation results obtained with both the in-plane and out-of-plane configurations, we now discuss implications for three-dimensional (3D) systems. As discussed in Section 5.1, the O-mode precursor wave emission is attributed to the large-amplitude magnetic field fluctuations generated in the shock-transition region, particularly by the WI at low $\sigma_e$. Since we can naturally expect the presence of such fluctuations in 3D, the O-mode precursor waves will also be excited. It is, however, not easy to estimate the relative emission efficiency between the O-mode and X-mode. Since the particle gyromotion in 3D should be less coherent than the 2D with the in-plane configuration, the O-mode wave power will become smaller. Concerning the gyromotion in the shock-transition region, the out-of-plane configuration may better represent 3D.

We have found that the particle acceleration efficiency also depends on the magnetic field configuration. Nonthermal particles are not generated in the in-plane configuration in 2D (see Appendix B), whereas a clear nonthermal tail is observed in the energy spectra for low $\sigma_e$ in the out-of-plane configuration (see Iwamoto et al. 2017, Figure 10). Considering the suppression of the cross-field diffusion in the in-plane configuration, again the out-of-plane configuration may be closer to 3D concerning the particle acceleration efficiency.

In any case, the important fact is that the intense coherent precursor wave can be excited for a wide range of $\sigma_e$ in both of the configurations. This strongly indicates that the intense coherent precursor wave emission is intrinsic to relativistic magnetized shocks and even in 3D.

5.3. Implication for WFA in Relativistic Shocks

Now we discuss the feasibility of the WFA in relativistic shocks. The WFA requires an intense electromagnetic wave in
the sense that the wave strength parameter $a = e \delta E / m_e c \omega$ is greater than unity, where $\delta E$ is the amplitude of the wave electric field and $\omega$ is the wave frequency (Kuramitsu et al. 2008). We estimated the strength parameter of the precursor wave with two different methods; one based on the oscillation amplitude of the transverse particle velocity, the other based on the wave amplitude. The details can be found in our previous paper (Iwamoto et al. 2017) except that the total wave power $\sqrt{\delta B_z^2 + \delta B_e^2}$ is used here. The results are shown in Figure 8, which demonstrates that the amplitudes of the precursor waves are indeed quite large.

Assuming a linear scaling of the strength parameter with respect to the Lorentz factor $\gamma_1$ (see, e.g., Hoshino 2008), we may estimate the region in the $\sigma_e - \gamma_1$ parameter space where the WFA is effective. In Figure 9, the solid and dashed lines indicate the estimates obtained by using the simulation results for the out-of-plane and in-plane configuration, respectively. This clearly indicates that higher Lorentz factors and moderate magnetizations are favorable for the WFA model. Again, we draw the same conclusion as Iwamoto et al. (2017) that highly relativistic external shocks of GRBs are candidate sites for the acceleration of ultra-high-energy cosmic rays.

The above discussion primarily focused on the wave amplitude. Although the large-amplitude precursor waves will induce wakefield in ion–electron plasmas, it is not clear yet whether the wakefield can sufficiently accelerate particles. The previous study demonstrated the WFA using a Gaussian laser pulse (Kuramitsu et al. 2008). However, the actual precursor waves are a superposition of waves continuously emitted from different positions in the shock front, and the spectra are rather broadband in wavenumber as shown in Figure 6. A ponderomotive force exerted by such waves should become weaker (e.g., Krue 1988; Hoshino 2008). Therefore, the generation of wakefield and particle acceleration may be less efficient. Also, different properties of the WI in pair and ion–electron plasmas, energy exchange between ions and electrons (e.g., Kumar et al. 2015) may influence the efficiency of the particle acceleration. These issues should be examined by performing shock simulations in ion–electron plasmas in the future.

6. Summary

In this work, we performed 2D simulations of relativistic perpendicular shocks in pair plasmas with the in-plane ambient magnetic field, and investigated the physics of the intense coherent precursor wave emission. In the in-plane configuration, O-mode as well as X-mode electromagnetic precursor waves are excited. We think that the O-mode waves are initially excited as X-mode by the SMI in the shock-transition region. Since the instabilities in the shock-transition region generate fluctuations in $B_z$ and disturb the ambient magnetic field, the SMI should excite X-mode waves that have $\delta B_z$, as well as $\delta B_e$. The generated waves having $\delta B_e$ may be mode-converted into O-mode during the propagation to the upstream region. The delay of the O-mode wave identified in our simulation is consistent with this model. We quantified the precursor wave amplitude as a function of the magnetization parameter $\sigma_e$ and compared the simulation results with that in the out-of-plane configuration by Iwamoto et al. (2017). The wave amplitude is sufficiently large to disturb the upstream plasma even in the Weibel-dominated regime and the transverse density filaments are generated as in the case of the out-of-plane configuration. We thus conclude that the precursor wave emission is the real nature of the realistic magnetized shocks.

In the range of $\sigma_e$ used in our simulations, the precursor wave keeps coherent and its amplitude is large enough to induce the wakefield. Therefore, the WFA may operate in relativistic ion–electron shocks.

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Appendix A

Electromagnetic Instabilities in the Shock-transition Region

Here we present a linear analysis of the electromagnetic instabilities excited in the shock-transition region. The dispersion relation for electromagnetic waves propagating parallel to an ambient magnetic field is given by (see, e.g., Yoon & Davidson 1987)

\[
D(k, \omega) = 1 - \frac{c^2k^2}{\omega^2} + \sum_s \frac{\Omega_{ps}^2}{\omega^2} \int_{-\infty}^{\infty} \int_0^\infty (\gamma_0 - c\mu_{1j})\frac{\partial F_{0j}}{\partial u_{0j}} + c\mu_{0j}\frac{\partial F_{00}}{\partial u_{01}} \pi \omega^2 \, du_{0j} \, du_{01},
\]

where \( \gamma = \sqrt{1 + u_{01}^2 + u_{02}^2} \) is the Lorentz factor, \( \Omega_{ps} \) is the nonrelativistic plasma frequency, \( \Omega_{cs} \) is the nonrelativistic cyclotron frequency, and \( F_{00} \) is the unperturbed distribution function normalized as follows:

\[
\int_{-\infty}^{\infty} \int_0^\infty F_{0s}(u_{0j}, u_{02}) \pi \omega^2 \, du_{0j} \, du_{01} = 1.
\]

The subscript \( s \) indicates particle species (i.e., electron and positron). In Equation (6), the positive (negative) sign corresponds to the right-hand (left-hand) polarization.

We assume a cold ring distribution for both electrons and positrons,

\[
F_{0s} = \frac{1}{2\pi\Omega_{0s}} \delta(u_{01} - u_{02}) \delta(u_{0j}).
\]

By substituting Equation (8) for (6), we obtain

\[
D(k, \omega) = 1 - \frac{c^2k^2}{\omega^2} - \sum_s \frac{\omega_{ps}^2}{\omega(\omega \pm \omega_{cs})} + \frac{1}{2} \left(1 - \frac{c^2k^2}{\omega^2}\right) \sum_s \left(1 - \frac{1}{\gamma_{0s}^2}\right) \frac{\omega_{ps}^2}{(\omega \pm \omega_{cs})^2},
\]

where \( \gamma_{0s} = \sqrt{1 + u_{0s}^2} \) is the initial Lorentz factor, \( \omega_{ps} = \Omega_{ps}/\sqrt{\gamma_{0s}} \) is the relativistic plasma frequency, and \( \omega_{cs} = \Omega_{cs}/\sqrt{\gamma_{0s}} \) is the relativistic cyclotron frequency. When \( \gamma_{0s} = 1 \), Equation (9) is identical to the dispersion relation in a cold magnetized plasma. Introducing \( \omega_{pe} = \omega_{ps} = \omega_{pc} \) and \( \omega_{ce} = -\omega_{ce} = \omega_{ce} > 0 \), Equation (9) reduces to

\[
D(k, \omega) = 1 - \frac{c^2k^2}{\omega^2} - \frac{2\omega_{pe}^2}{\omega^2 - \omega_{ce}^2} + \beta_0^2 \left(1 - \frac{c^2k^2}{\omega^2}\right) \frac{\omega_{pe}^2(\omega^2 + \omega_{ce}^2)}{(\omega^2 - \omega_{ce}^2)^2},
\]

Figure 7. Energy of the precursor wave emission given in units of the ambient upstream magnetic field energy (left) and the upstream kinetic energy (right) as a function of \( \sigma_e \). The red, blue, and magenta indicate \( \delta B_{ps}^2, \delta B_{ce}^2 \), and \( \delta B_{ps}^2 + \delta B_{ce}^2 \), respectively. The simulation results by Iwamoto et al. (2017) are shown in green.

Figure 8. Strength parameter of the precursor wave as a function of \( \sigma_e \). The solid and dashed lines indicate the estimation method based on the particle quiver velocity and the wave amplitude, respectively.

Figure 9. Parameter space plot in \( \sigma_e \) and \( \gamma_{1s} \). The solid and dashed lines indicate the out-of-plane and in-plane cases, respectively. Strength parameter \( \alpha \) is greater than unity in the region above each line.
where $\beta_0 = u_0 c / \gamma_0 e$. $D(k, \omega) = 0$ is expressed as follows:

\[
\left( \frac{\omega}{\omega_{pe}} \right)^6 - \left( \frac{c^2 k^2 \gamma_0}{\omega_{pe}^2} + 2 \sigma_e - \beta_0^2 \right) \left( \frac{\omega}{\omega_{pe}} \right)^4 \\
+ \left[ \left( \frac{2\sigma_e - \beta_0^2}{\omega_{pe}^2} \right) c^2 k^2 + \sigma_e (\sigma_e + \beta_0^2) + 2 \right] \left( \frac{\omega}{\omega_{pe}} \right)^2 \\
- \sigma_e (\sigma_e + \beta_0^2) \frac{c^2 k^2}{\omega_{pe}^2} = 0, \tag{11}
\]

where $\sigma_e = \omega_{ce}^2 / \omega_{pe}^2$. In this linear analysis, we assume that the wavenumber $k$ is a real number and consider only the region where $k \geq 0$ and $\omega \geq 0$ because of the symmetry of $k$ and $\omega$ for a pair plasma.

First, we study the case of $\sigma_e > (2 - \beta_0^2)^2 / 8 \beta_0^2$, $\omega$ becomes a complex number at $k \leq k_1$ and $k \geq k_2$. Here, the threshold wavenumbers $k_1$ and $k_2$ are determined from

\[
\left( \frac{c k_1}{\omega_{pe}} \right)^2 = f \left( \frac{\omega_1^2}{\omega_{pe}^2} \right), \tag{12}
\]

\[
\left( \frac{c k_2}{\omega_{pe}} \right)^2 = f \left( \frac{\omega_2^2}{\omega_{pe}^2} \right), \tag{13}
\]

\[
k_1 < k_2, \tag{14}
\]

\[
f(x) = x - 2 - \frac{2(\sigma_e - \beta_0^2) x - 2 \sigma_e (\sigma_e + \beta_0^2)}{x^2 - (2\sigma_e - \beta_0^2) x + \sigma_e (\sigma_e + \beta_0^2)}. \tag{15}
\]

$\omega_1$ and $\omega_2$ satisfy

\[
f' \left( \frac{\omega_1^2}{\omega_{pe}^2} \right) = f' \left( \frac{\omega_2^2}{\omega_{pe}^2} \right) = 0, \tag{16}
\]

\[
\omega_1 > \omega_2. \tag{17}
\]

Figure 10 shows the dispersion relation with $\sigma_e = 0.3$ and $\gamma_0 = 40$ numerically obtained from Equation (11). The real and imaginary part of the frequency is shown by the solid and dashed lines, respectively. The unstable branch for $k \leq k_1$ and $k \geq k_2$ is connected to the electromagnetic and Alfvén mode branch, respectively. We think the unstable mode for $k \geq k_2$ corresponds to the Alfvén-ion-cyclotron instability in ion-electron plasmas. The growth rate of the mode $k \leq k_1$ has its maximum at $k = 0$,

\[
\frac{\text{Re}(\omega_{\text{max}})}{\omega_{pe}} = \sqrt{\frac{1}{2} \sigma_e - \frac{1}{4} \beta_0^2 + \frac{1}{2} + \frac{1}{2} \sqrt{\sigma_e (\sigma_e + \beta_0^2) + 2}}, \tag{18}
\]

\[
\frac{\text{Im}(\omega_{\text{max}})}{\omega_{pe}} = \sqrt{\frac{1}{2} \sigma_e - \frac{1}{4} \beta_0^2 - \frac{1}{2} + \frac{1}{2} \sqrt{\sigma_e (\sigma_e + \beta_0^2) + 2}}. \tag{19}
\]

The maximum growth rate of the mode $k \geq k_2$ occurs for $ck/\omega_{pe} \gg 1$,

\[
\frac{\text{Re}(\omega_{\text{max}})}{\omega_{pe}} = \sqrt{\frac{1}{2} \sigma_e - \frac{1}{4} \beta_0^2 + \frac{1}{2} + \frac{1}{2} \sqrt{\sigma_e (\sigma_e + \beta_0^2) + 2}}, \tag{20}
\]

\[
\frac{\text{Im}(\omega_{\text{max}})}{\omega_{pe}} = \sqrt{\frac{1}{2} \sigma_e + \frac{1}{4} \beta_0^2 + \frac{1}{2} \sqrt{\sigma_e (\sigma_e + \beta_0^2) + 2}}. \tag{21}
\]

If $\sigma_e \gg 1$, then the maximum growth rates for both modes are written as follows:

\[
\frac{\text{Re}(\omega_{\text{max}})}{\omega_{pe}} \sim \omega_{ce}, \tag{22}
\]

\[
\frac{\text{Im}(\omega_{\text{max}})}{\omega_{pe}} \sim \frac{\beta_0}{\sqrt{2}} \omega_{pe}. \tag{23}
\]

Second, we study the case of $\beta_0^2 / 8 \leq \sigma_e \leq (2 - \beta_0^2)^2 / 8 \beta_0^2$. For highly relativistic plasma $\beta_0 \sim 1$, the condition is satisfied within a very narrow range of $\sigma_e$. In this case, only the instability on the Alfvén mode branch exists at $k \geq k_2$. The dispersion relation is shown in Figure 11. The maximum growth rate is identical to Equations (20) and (21).
Finally, we study the case of $\sigma_\epsilon < \beta_0^2/8$. $\omega$ is a complex number and a pure imaginary number for $k_2 \leq k < k_3$ and $k > k_3$, respectively. Here, the threshold wavenumber $k_3$ is determined by

$$\left(\frac{ck_3}{\omega_{pe}}\right)^2 = f\left(\frac{\omega^2}{\omega_{pe}^2}\right).$$

The threshold frequency $\omega_3$ satisfies

$$f'(\frac{\omega^2}{\omega_{pe}^2}) = 0,$$

$$\omega_- < \text{Im}(\omega_3) < \omega_+,$$

$$\text{Re}(\omega_3) = 0,$$

where

$$\frac{\omega_+}{\omega_{pe}} = \sqrt{-\sigma_e + \frac{1}{2}\beta_0^2 \pm \frac{1}{2}\beta_0 \sqrt{\beta_0^2 - 8\sigma_e}}.$$  

Figure 12 shows the dispersion relation with $\sigma_e = 3 \times 10^{-3}$, $\gamma_0 = 40$. For $\sigma_e \leq \beta_0^2/8$, there are two purely growing modes when the wavenumber is greater than the threshold wavenumber $k_3$. We think the upper unstable branch corresponds to the WI. The growth rates of these modes asymptotically approach $\omega_\pm$ as the wavenumber increases. The maximum growth rate is expressed as follows:

$$\text{Im}(\omega_{\text{max}}) \sim \omega_+.$$  

If $\sigma_e \ll \beta_0^2$, then the maximum growth rate is written as follows:

$$\text{Im}(\omega_{\text{max}}) \sim \beta_0 \omega_{pe}.$$  

We now compare our simulation results with the maximum linear growth rate. Our linear analysis indicates that the electromagnetic fields perpendicular to the ambient magnetic field are induced by the electromagnetic instabilities. In fact, previous simulations showed that the instabilities in the shock-transition region excite $B_y$ as well as $B_z$ in the in-plane configuration (e.g., Winske & Quest 1988; Matsukiyto & Scholer 2006). The maximum values of the $x$ and $z$ components of the magnetic field energy averaged over the $y$ axis are determined for each snapshot, and are shown in Figure 13 for $\sigma_e = 3 \times 10^{-1}$ (left) and $\sigma_e = 3 \times 10^{-3}$ (right). The red and blue solid lines indicate the $x$ and $z$ components of the magnetic field energy, respectively. The magnetic field energy $\epsilon_B = B^2/8\pi\mu_0 c^2$ are expressed in units of the upstream kinetic energy. The maximum linear growth rate determined by Equations (21) and (29) is also shown in Figure 13 with the black dashed lines.

For $\sigma_e = 3 \times 10^{-1}$, although the maximum linear growth rate is consistent with the simulation result, it is difficult to differentiate which instabilities generate the magnetic field fluctuations because the maximum growth rates (Equations (21) and (29)) are almost the same. Furthermore, our analysis assumes a cold ring distribution and ignores possible kinetic effects that should become important at relatively short wavelengths. In any case, we think that the instabilities excited in the shock-transition region generate the magnetic field fluctuations in our simulation.

For $\sigma_e = 3 \times 10^{-3}$, the maximum linear growth rate gives good agreement with the simulation result. In addition, the maximum energy of the fluctuating magnetic field saturates about 10%–20% of the upstream bulk kinetic energy. This result is consistent with the previous studies (Kato 2007; Chang et al. 2008; Sironi & Spitkovsky 2011). Therefore, we conclude that the fluctuations in $B_y$ and $B_z$ in the shock-transition region result from the WI.

**Appendix B**

Particle Energy Spectra

Figure 14 shows the downstream energy spectra of electrons for $\sigma_e = 3 \times 10^{-1}$ and $3 \times 10^{-3}$, which are normalized as follows:

$$\int f_\gamma(\gamma)d\gamma = 1.$$  

The energy spectra of positrons are identical to those of electrons. We followed the time evolution from $\omega_{pe}t = 100$ up to $\omega_{pe}t = 500$. For both $\sigma_e$, the measured distribution reaches a steady state by the end of our simulation. The energy spectra can be well-fitted with 3D relativistic Maxwellian,

$$f(\gamma)d\gamma \propto \gamma \sqrt{\gamma^2 - 1} \exp\left(-\frac{\gamma mc^2}{kT}\right).$$

Note that the degree of freedom is three in the in-plane configuration. The fitting result indicates that the downstream
particles are completely thermalized. A clear suprathermal tail is not observed in the range of $\sigma_e$ used in our simulations.

In the out-of-plane configuration, however, a suprathermal tail is visible for $\sigma_e = 3 \times 10^{-3}$ (see Iwamoto et al. 2017, Figure 10). This agrees with the simulation result by Sironi et al. (2013). They suggested that the particle acceleration can be explained in terms of a Fermi process due to the strong turbulence generated by the WI and that the suppression of the cross-field diffusion (Jokipii et al. 1993; Jones et al. 1998) may result in low-level injection of particles into the Fermi process. Our result also confirms that the orientation of the ambient magnetic field affects the efficiency of particle acceleration.

**ORCID iDs**

Masanori Iwamoto  
https://orcid.org/0000-0003-2255-5229

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Figure 14. Downstream energy spectra of electrons: $\sigma_e = 3 \times 10^{-1}$ (left) and $\sigma_e = 3 \times 10^{-3}$ (right). The black dashed lines indicate a 3D relativistic Maxwellian fitting result.