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RESEARCH ARTICLE

Low frequency Whistler waves excited in fast magnetic reconnection processes

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Whistler waves generated in fast magnetic reconnection processes of collisionless high beta plasmas are reviewed in experiments and satellite observations, as well as in theory and simulation, and further studied in the two-fluid theory. It is found that low frequency whistler waves can be excited in the ion inertial range of the reconnection region. The wave is found right-handed polarized with a quadrupolar out-of-plane magnetic perturbation, in accord with satellite observations in the geomagnetosphere.

Keywords Whistler waves, mode conversion, Hall MHD, magnetic reconnection

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1 Introduction

Magnetic reconnection is of essential importance in many laboratory and space plasma physics processes associated with configuration relaxations and fast energy releases of magnetic fields. The process of reconnection occurs on a thin layer called reconnection layer or diffusion region where kinetic and/or dissipative effects become substantially important. In the last decade, whistler waves generated in magnetic reconnection layers have been taken as a major signature of collisionless magnetic reconnection. Observationally whistler waves have been detected on current structures and magnetic reconnection regions in laboratory and magnetosphere plasmas [1–12], where the plasma beta is on the order of unity. It was found that the wave propagated obliquely to the ambient magnetic field, while the polarity of the wave was right-handed in the plane perpendicular to the ambient magnetic field, a typical electron mode feature [9–12]. On the other hand however, the earlier theory predicted that the fast rate of reconnection in collisionless plasmas was a direct consequence of the quadratic nature of the dispersion character of whistler waves, $\omega(\approx k^2 d_e^2 \Omega_{ce}) \sim k^2$ where $\Omega_{ce} = eB/(m_e c)$ is the electron cyclotron frequency and $d_e = c/\omega_{pe}$ is the electron skin depth with the speed of light c and the plasma frequency of $\omega_{pe} = (4\pi n_0 e^2/m_e)^{1/2}$, which controlled the plasma dynamics at the small scale δ with $k \sim 1/\delta$ [13–15]. Nevertheless, later theoretical analysis showed that the whistler wave should be obliquely propagated with a much lower frequency of $\omega \approx k_{/\!/} k_\perp d_e^2 \Omega_{ce} = k_{/\!/} k_\perp d_i^2 \Omega_{ci} = k_{/\!/} V_A k_\perp d_i$, due to $k_{/\!/} \ll k_\perp \approx k \sim 1/\delta$, with $V_A = B/\sqrt{4\pi m_i n_0}$ being the local Alfvén velocity and $d_i = c/\omega_{pi}$ being the ion inertial length with the ion plasma frequency $\omega_{pi} = (4\pi n_0 e^2/m_i)^{1/2}$ [16, 17]. This dispersion relation of the obliquely propagated low frequency whistler mode was recently confirmed by k-filter analysis for Cluster observation data [12]. In a recent review on fast magnetic reconnection however, the whistler wave was thought of as a standing wave, with the reconnected magnetic field lines in the outflow region as its half spatial period, and the field line relaxation as the wave oscillation amplitude [18]. Clearly, this picture cannot explain the excitation mechanism of the whistler wave and its polarization property, as well as the mode structure observed. Therefore these issues are still open questions for further understanding in theory and in comparison with observations.

We thus further study the whistler wave generation in fast magnetic reconnection processes and its propagation and polarity properties in this paper. The wave equation is derived from a Harris-sheet-like equilibrium in Section 2. The singularity of magnetic reconnection is found in magnetohydrodynamics (MHD) approximation. We then resolve the singularity by WKB method and obtain the whistler wave solution on the reconnection layer. The properties of the whistler wave are further discussed in Section 3. The paper is then concluded with a summary in Section 4.

2 Singularities in MHD regime and dispersion relation of the mode

In high beta plasmas, such as in the interplanetary plasma medium and the geomagnetosphere, the initial static state equilibrium for magnetic reconnection can be approximately modeled in a two dimensional (2D) configuration

$$\boldsymbol{B}_0 = \hat{y}B_0(x) = \hat{z}\times\nabla\psi_0(x), \quad |x|\leqslant a$$

with $\psi_0(x) = B_0x^2/2a$, $B_0(x) = B_0x/a$, generated by an approximate constant current $\boldsymbol{J}_0 \approx \hat{z}cB_0/(4\pi a)$. It can be seen as the central segment of a Harris sheet. In such an equilibrium, if the plasma is isothermal, the density then has a distribution of $n_0(x) \approx n_0[1-x^2/(2a^2)]$, with $\beta_0 \equiv \beta(a) = 4\pi n_0(T_{e0}+T_{i0})/B_0^2 = 1$, where the plasma beta $\beta = \beta(x) \equiv 8\pi n_0(x)(T_{e0}+T_{i0})/B_0^2(x)$. Thus, for the equilibrium we have $\beta \geqslant 1$ everywhere in $|x| \leqslant a$. For perturbations in a form of $f_1(x,y,t) = \tilde{f}(x)\mathrm{e}^{-\mathrm{i}(\omega_0 t - k_{/\!/}y)}$, the velocity and magnetic field perturbations are

$$\boldsymbol{v}_1(x,y,t) = (\hat{z}\times\nabla\tilde{\varphi}(x) + \hat{z}\tilde{v}_z(x))\mathrm{e}^{-\mathrm{i}(\omega_0 t - k_{/\!/}y)} \tag{1}$$

$$\boldsymbol{B}_1(x,y,t) = (\hat{z}\times\nabla\tilde{\psi}(x) + \hat{z}\tilde{B}_z(x))\mathrm{e}^{-\mathrm{i}(\omega_0 t - k_{/\!/}y)} \tag{2}$$

The wave equation in the ideal MHD approximation can then be easily derived as

$$\frac{\mathrm{d}}{\mathrm{d}x}\left(\varepsilon_A(x)\frac{\mathrm{d}\tilde{\xi}_x}{\mathrm{d}x}\right) - k_{/\!/}^2\varepsilon_A(x)\tilde{\xi}_x = 0 \tag{3}$$

with the x-component displacement perturbation $\tilde{\xi}_x = \mathrm{i}\tilde{v}_x/\omega = k_{/\!/}\tilde{\varphi}/\omega$, and

$$\varepsilon_A(x) \equiv \frac{\omega_0^2}{k_{/\!/}^2} - V_A^2(x) = \frac{\omega_0^2}{k_{/\!/}^2} - \frac{B_0^2(x)}{4\pi\rho(x)} = \frac{\omega_0^2}{k_{/\!/}^2} - \frac{B_0^2x^2}{4\pi\rho_0a^2(1-x^2/a^2)} \tag{4}$$

where $\rho_0 \equiv n_0m_i$. In the small x region then, $x^2/a^2 \ll 1$, we approximately have

$$\varepsilon_A(x) \approx \frac{\omega_0^2}{k_{/\!/}^2} - \frac{V_{A0}^2x^2}{a^2} \tag{4$'$}$$

Clearly, at $\varepsilon_A(x_0) = 0$, i.e., $x_0/a = \omega_0/(k_{/\!/}V_{A0})$ with $V_{A0} = B_0/\sqrt{4\pi m_i n_0}$, we have a logarithm singularity for $\tilde{\xi}_x$ (as well as $\tilde{\varphi}$ and $\tilde{\psi}$), and a $1/(x-x_0)$ singularity for $\tilde{\xi}_y \sim \tilde{\varphi}'$ [19]. This singularity leads to the Alfvén resonance. Thus, we call it the Alfvén resonance singularity. Additionally, however, there is another singularity at $x = 0$, if $\omega_0 = 0$ and thus

$$\varepsilon_A(x) = -V_A^2(x) = -V_{A0}^2x^2/a^2 \tag{5}$$

Clearly, it is a stronger singularity than that of Alfvén resonances, with a $1/x$ singularity for $\tilde{\xi}_x$ (as well as $\tilde{\varphi}$ and $\tilde{\psi}$) and a $1/x^2$ singularity for $\tilde{\xi}_y \sim \tilde{\varphi}'$. We then call it the reconnection singularity due to the fact of its inducing magnetic reconnection on the singular surface $x = 0$.

To resolve the singularity, kinetic and/or dissipation effects should be introduced [19]. Here we apply the two-fluid theory to the reconnection layer (near $x = 0$) to get the leading order equations in the dimensionless form of

$$\frac{\partial}{\partial t}\frac{\partial^2}{\partial x^2}\tilde{\varphi}^{(0)} = \pm_x\mathrm{i}k_{/\!/}V_A\frac{\partial^2}{\partial x^2}\tilde{\psi}^{(0)} \tag{6}$$

$$\frac{\partial}{\partial t}\tilde{\psi}^{(0)} - \mathrm{i}(\pm_x k_{/\!/}V_A)\tilde{\varphi}^{(0)} = \pm_x\mathrm{i}k_{/\!/}V_Ad_i\tilde{B}_z^{(0)} \tag{7}$$

$$\frac{\partial}{\partial t}\tilde{v}_z^{(0)} = \pm_x\mathrm{i}k_{/\!/}V_A\tilde{B}_z^{(0)} \tag{8}$$

$$\frac{\partial}{\partial t}\tilde{B}_z^{(0)} = -\mathrm{i}(\pm_x k_{/\!/}V_A)d_i\frac{\partial^2}{\partial x^2}\tilde{\psi}^{(0)} \pm_x \mathrm{i}k_{/\!/}V_A(x)\tilde{v}_z^{(0)} \tag{9}$$

where "$\pm_x$" corresponds to the sign of x, and the dimensional translations are

$$\frac{\tilde{\psi}}{B_0a}\to\tilde{\psi},\quad \frac{\tilde{\varphi}}{V_{A0}a}\to\tilde{\varphi},\quad \frac{\tilde{B}_z}{B_0}\to\tilde{B}_z,\quad \frac{\tilde{v}_z}{V_{A0}}\to\tilde{v}_z$$

$$a\nabla\to\nabla,\quad \frac{tV_{A0}}{a}\to t,\quad \frac{d_i}{a}\to d_i$$

To solve Eqs. (6)–(9), we use WKB method by applying the "eikonal" approximation

$$\{\tilde{\varphi}^{(0)}(x,t),\ \tilde{\psi}^{(0)}(x,t)\} = \{\hat{\varphi},\ \hat{\psi}\}\mathrm{e}^{\mathrm{i}S(x,t)} \tag{10}$$

with

$$k \approx k_\perp = k_x = \frac{\partial S}{\partial x},\quad \left|\frac{\partial k_x}{\partial x}\right| \ll k_x^2 \tag{11}$$

$$\omega = -\frac{\partial S}{\partial t},\quad \left|\frac{\partial \omega}{\partial t}\right| \ll \omega^2 \tag{12}$$

Substituting (10)–(12) to (6)–(9) we then have the dispersion relation

$$k_\perp^2 = \frac{(\omega^2 - k_{/\!/}^2V_A^2)^2}{\omega^2d_i^2k_{/\!/}^2V_A^2} \tag{13}$$

In the regime of $k_\perp^2d_i^2 \gg 1$, approximately

$$\omega \approx \mp_0 k_{/\!/}V_Ak_\perp d_i = \mp_0 k_{/\!/}k_\perp d_i^2\Omega_{ci} = \mp_0 k_{/\!/}k_\perp d_e^2\Omega_{ce} \tag{14}$$

Here the sign of "$\mp_0$" corresponds to the antisymmetry about the origin, depending on the propagation prop-

Frontiers of Physics

erty of the wave to make the frequency positive in the propagation cases. Due to the feature of the reconnection process, we have: $k_\perp < 0$, $k_{/\!/} > 0$ in Quadrant I ($x > 0$, $y > 0$); $k_\perp < 0$, $k_{/\!/} < 0$ in Quadrant II ($x > 0$, $y < 0$); $k_\perp > 0$, $k_{/\!/} < 0$ in Quadrant III ($x < 0$, $y < 0$); and $k_\perp > 0$, $k_{/\!/} > 0$ in Quadrant IV ($x < 0$, $y > 0$). Then the sign of (14) is the same as the quadrupolar Hall field with "$\mp_0$" corresponding to "–" in Quadrants I & III, and "+" in Quadrants II & IV. Eq. (14) is then the dispersion relation of low frequency whistler waves, the same as the low frequency whistler mode predicted and observed in the collisionless reconnection layer [12, 16].

3 Properties of the Whistler modes

From the approximation in deriving (14), it is obvious that wave is generated in a region inside of the reconnection layer, $d_e < 1/|k_\perp| < d_i$, called ion inertial layer. We then analyze the features of the modes in this region.

3.1 The group velocity and reconnection rate

From Eq. (14), no losing of generality we can write out the dispersion relation $\omega \approx k_{/\!/} k_\perp d_e^2 \Omega_{ce}$ in Quadrant IV. It is different from the conventional whistler waves with a dispersion relation of $\omega = k^2 d_e^2 \Omega_{ce}$. The phase velocity of the latter is $V_{ph} = k d_e^2 \Omega_{ce}$, while its group velocity is $V_{gr} = 2k d_e^2 \Omega_{ce}$. On the other hand, the whistler wave in the reconnection layer is obliquely propagating with its phase velocity almost perpendicular to the equilibrium magnetic field due to $k \approx k_\perp = k_x$. Nevertheless, similar to the kinetic Alfvén wave (KAW), perpendicular group velocity of the oblique whistler is small, as

$$V_{gr\perp} = \frac{\partial \omega}{\partial k_\perp} \approx k_{/\!/} d_e^2 \Omega_{ce} \tag{15}$$

while its parallel group velocity

$$V_{gr/\!/} = \frac{\partial \omega}{\partial k_{/\!/}} = k_\perp d_e^2 \Omega_{ce} \gg V_{gr\perp} \tag{16}$$

Thus, the wave energy propagates mainly along the magnetic field.

By the wave energy conservation, we can approximately get the relation

$$V_{gr\perp} L = k_{/\!/} d_e^2 \Omega_{ce} L \sim V_{gr/\!/} \delta = k_\perp d_e^2 \Omega_{ce} \delta \tag{17}$$

leading to the estimate of the reconnection rate, $\delta/L \sim |k_{/\!/}/k_\perp|$.

3.2 Electric vectors and polarity of the mode

From Eqs. (2) and (10)–(12), we can get $\tilde{B}_x^{(0)} = -\mathrm{i}k_{/\!/}\tilde{\psi}^{(0)}$ and $\tilde{B}_y^{(0)} = \mathrm{i}k_\perp \tilde{\psi}^{(0)}$ easily. Then by Eqs. (8), (9) and (14), we further obtain $\tilde{B}_z^{(0)} = \pm_y k_\perp \tilde{\psi}^{(0)}$, where

$$\pm_y = - \pm_x \cdot \mp_0 = \pm_x \cdot \pm_0$$

corresponding to the sign of y. Then by the Faraday law

$$\mathrm{i}k_{/\!/}\tilde{E}_z^{(0)} = \mathrm{i}\frac{\omega}{c}\tilde{B}_x^{(0)} = \frac{\omega}{c}k_{/\!/}\tilde{\psi}^{(0)} \tag{18}$$

$$-\mathrm{i}k_\perp \tilde{E}_z^{(0)} = \mathrm{i}\frac{\omega}{c}\tilde{B}_y^{(0)} = -\frac{\omega}{c}k_\perp \tilde{\psi}^{(0)} \tag{19}$$

$$\begin{aligned}\mathrm{i}k_\perp \tilde{E}_y^{(0)} - \mathrm{i}k_{/\!/}\tilde{E}_x^{(0)} &\approx -\mathrm{i}k_{/\!/}\tilde{E}_x^{(0)} = \mathrm{i}\frac{\omega}{c}\tilde{B}_z^{(0)} \\ &= \pm_y \mathrm{i}\frac{\omega}{c}k_\perp \tilde{\psi}^{(0)}\end{aligned} \tag{20}$$

we can then have

$$\tilde{E}_z^{(0)} = -\mathrm{i}\frac{\omega}{c}\tilde{\psi}^{(0)} \tag{21}$$

and in Hall MHD regime, $\tilde{E}_y^{(0)} \ll \tilde{E}_x^{(0)}$,

$$\begin{aligned}\tilde{E}_x^{(0)} &\approx \mp_y \frac{k_\perp \omega}{k_{/\!/} c}\tilde{\psi}^{(0)} = \mp_y \mp_0 \left|\frac{k_\perp \omega}{k_{/\!/} c}\right| \tilde{\psi}^{(0)} \\ &= \pm_x \left|\frac{k_\perp \omega}{k_{/\!/} c}\right| \tilde{\psi}^{(0)}\end{aligned} \tag{22}$$

Then the polarity of the wave can be determined by

$$\frac{\mathrm{i}\tilde{E}_z^{(0)}}{\tilde{E}_x^{(0)}} = \pm_x \left|\frac{k_{/\!/}}{k_\perp}\right| \tag{23}$$

It is clearly a right handed polarized wave with respect to the equilibrium magnetic field. On the upper half of the domain ($x > 0$), $\mathrm{i}\tilde{E}_z^{(0)}/\tilde{E}_x^{(0)} > 0$, the wave is right handed elliptically polarized. On the lower half, however, the equilibrium magnetic field is pointed to the $-y$ direction. Then $(z, x, -y)$ is a left hand system and $\mathrm{i}\tilde{E}_z^{(0)}/\tilde{E}_x^{(0)} < 0$ is again right handed elliptically polarized. Then the wave can be decomposed into an electrostatic component of $\tilde{E}'_x = \tilde{E}_x^{(0)} - \pm_x|\omega/c|\tilde{\psi}^{(0)}$ and a right handed circularly polarized "extraordinary" wave of $\pm_x|\omega/c|\tilde{\psi}^{(0)}\hat{x} + \tilde{E}_z^{(0)}\hat{z}$. The latter is the obliquely propagated and right handed polarized low frequency whistler wave.

3.3 The standing wave solution and Hall field structure

The pure growth solution of the mode is clearly *a standing wave* with $\omega^2 = -\gamma^2$. And the dispersion relation has a form of

$$k_\perp^2 = -\frac{(\gamma^2 + k_{/\!/}^2 V_A^2)^2}{\gamma^2 d_i^2 k_{/\!/}^2 V_A^2} < 0, \quad k_\perp = \pm_x \mathrm{i}|k_\perp| \tag{24}$$

Therefore such a pure growth mode of the standing wave solution has a spatial decay factor of $e^{\mp_x |k_\perp| x}$, and $1/|k_\perp| \sim d_i$ can be seen as the *skin depth* of the mode.

Thus, the mode structure should be confined in the skin depth region. Then the relation (14) is reduced to

$$\gamma \approx k_{/\!/} V_A |k_\perp| d_i = k_{/\!/} |k_\perp| d_i^2 \Omega_{ci} = k_{/\!/} |k_\perp| d_e^2 \Omega_{ce} \tag{25}$$

where $k_{/\!/} > 0$ always for the standing wave solution. According to Eqs. (8), (9) and (24), (25), the standing wave mode then has an out-of-plane magnetic component of

$$\tilde{B}_z^{(0)} \approx -i \frac{\pm_x k_{/\!/} V_A d_i |k_\perp|^2}{\gamma} \tilde{\psi}^{(0)} = \mp_x i |k_\perp| \tilde{\psi}^{(0)} \tag{26}$$

Since ψ_1 is symmetric on $y = 0$, and the factor of $i = \mathrm{e}^{\mathrm{i}\pi/2}$ makes a phase shift of $\pi/2$, then in the concerned region it is equivalent to the sign of $\pm_y$. Therefore $\mp_x i \rightarrow \mp_x \pm_y = \mp_0$. Thus, $\tilde{B}_z^{(0)}$ is antisymmetric about the origin with the "–" sign in Quadrants I & III, and the "+" in Quadrants II & IV. It is accorded to the quadrupolar out-of-plane field observed in satellite data and experiments.

4 Conclusion and summary

We in this paper analyze the properties, such as dispersion, generation, propagation, and mode structure etc., of the obliquely propagated low frequency whistler modes in fast reconnection processes for high beta plasmas. The basic picture of the low frequency whistler modes generation can be summarized as follows.

An outside excited incident wave of $\boldsymbol{\xi}_1(x, y, t) = \tilde{\boldsymbol{\xi}}(x) \exp(-\mathrm{i}(\omega_0 t - k_{/\!/} y))$ can be transmitted in a Harris sheet like plasma, either on upper and lower boundaries or by antenna coupling. It then has two Alfvén resonance singularities on the layer of $x = \pm x_0$, where $x_0 = a\omega_0/(k_{/\!/} V_{A0})$. If the incident perturbation is a standing wave with $\omega_0 = 0$, it has a reconnection singularity at $x = 0$ in the ideal MHD regime. The singularity is $1/x$ for $\tilde{\xi}_x$ (as well as $\tilde{\varphi}$ and $\tilde{\psi}$) and $1/x^2$ for $\tilde{\xi}_y \sim \tilde{\varphi}'$, much stronger than that of Alfvén resonances.

To solve the singularity, we apply the WKB method and derive the dispersion relation of $\omega \approx \mp_0 k_{/\!/} V_A k_\perp d_i = \mp_0 k_{/\!/} k_\perp d_i^2 \Omega_{ci} = \mp_0 k_{/\!/} k_\perp d_e^2 \Omega_{ce}$, with the sign of "$\mp_0$" having the antisymmetric about the origin with a "–" sign in Quadrants I & III, and a "+" in Quadrants II & IV. From the dispersion relation, we find that the propagation features of the wave are different from that of the conventional whistler waves. While the phase velocity of the whistler wave in the reconnection layer is now oblique, almost perpendicular to the magnetic field, the group velocity of the wave is however almost along the field line. Furthermore, the wave is found right handed polarized, a feature detected in spacecraft observations.

A standing wave solution for the pure growth mode is also found. The analysis for the mode structure finds that the out-of-plane magnetic component of the mode is quadrupolar, as observed for fast magnetic reconnection in space plasmas and found in Hall MHD reconnection in simulations.

The above features of the low frequency whistler modes in the reconnection layer are in accord with experiments, observations, and simulations. However, the magnitude of the modes, both for the right handed polarized wave and the out-of-plane components, should be calculated in the nonlinear regime. Furthermore, the cause of the whistler wave polarity and excitation of modes in the electron skin depth regime are also open questions to be answered.

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