Chaos-Induced Resistivity in the Collisionless Reconnection Region (Postprint)

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Abstract

Collisionless magnetic reconnection, as a mechanism for efficiently converting magnetic energy into plasma kinetic and thermal energy, has been widely applied to explain various explosive plasma activities such as solar flares and geomagnetic storms. However, the microscopic physical mechanism of anomalous resistivity in collisionless reconnection regions remains an unresolved fundamental issue. Among the many formation mechanisms of anomalous resistivity, chaos-induced resistivity—generated based on the chaoticity of particle orbits near magnetic null points—although not the most popular formation mechanism, possesses the clearest microscopic physical picture. This paper reviews the early studies and basic theoretical models of chaos-induced resistivity in collisionless reconnection regions, introduces recent advances in research on chaos-induced resistivity, and discusses future research directions for chaos-induced resistivity.

Full Text

Chaos-induced Resistivity in Collisionless Magnetic Reconnection Region

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Abstract

Collisionless magnetic reconnection, which efficiently converts magnetic energy into plasma kinetic and thermal energy, has been widely applied to explain various eruptive phenomena in plasmas such as solar flares and geomagnetic storms. However, the microscopic physical mechanism of anomalous resistivity in collisionless reconnection regions remains an unresolved fundamental problem. Among the many mechanisms proposed for anomalous resistivity formation, chaos-induced resistivity—arising from the chaotic nature of particle orbits near magnetic null points—is not the most commonly discussed, yet it offers the clearest microscopic physical picture. This paper reviews early studies and the basic theoretical model of chaos-induced resistivity in collisionless reconnection regions, introduces recent progress in this area, and outlines future research directions.

Keywords: magnetic reconnection, anomalous resistivity, chaos-induced resistivity, Sun: coronal mass ejections

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The concept of magnetic reconnection was first introduced by Giovanelli in 1946 to explain the mechanism of solar flares. Today, magnetic reconnection is widely recognized as an effective mechanism for the rapid release of magnetic energy in phenomena including solar flares, coronal mass ejections, magnetic storms, magnetospheric substorms, and auroras. According to the fundamental concept of magnetic reconnection, the primary characteristics of reconnection are the change in magnetic topology and the acceleration and heating of plasma, which necessarily requires breaking the frozen-in condition. In classical plasma theory, plasma resistivity arises from Coulomb collisions between charged particles, known as classical Spitzer resistivity or classical collisional resistivity. This resistivity can break the frozen-in condition and enable magnetic reconnection. However, in space plasmas, the characteristic timescale of plasma parameter variations is much shorter than the mean collision time between charged particles, making the resistivity from Coulomb collisions extremely small and insufficient to break the frozen-in condition. Consequently, magnetic reconnection would hardly occur. This raises the question: what breaks the frozen-in condition in collisionless magnetic reconnection? To address this, researchers proposed the concept of anomalous resistivity.

The earliest model of magnetic reconnection was the Sweet-Parker model, but studies revealed that it predicts a very low reconnection rate and long reconnection times. For example, in coronal jet ejection events, the theoretically predicted reconnection time is on the order of months, whereas typical observed coronal jet timescales are minutes. Therefore, the Sweet-Parker model is often referred to as slow reconnection. Subsequent work on magnetic reconnection has focused largely on increasing the reconnection rate, with some studies replacing the classical collisional resistivity in the Sweet-Parker model with anomalous resistivity to achieve higher rates. In 1964, Petschek proposed the Petschek reconnection model, in which magnetic energy can be dissipated not only through Ohmic heating but also through slow shocks that convert magnetic energy into plasma kinetic and thermal energy. This model is known as fast reconnection. However, numerical simulations found the Petschek model to be unstable, and later studies showed that stable fast reconnection requires non-uniform resistivity. In summary, both classical magnetic reconnection models require anomalous resistivity beyond classical collisional resistivity, and this anomalous resistivity should be non-uniformly distributed in the reconnection region.

The essence of anomalous resistivity in collisionless magnetic reconnection is the randomization of current-carrying charged particle motion in the reconnection current sheet, which transforms macroscopically ordered energy into microscopically disordered energy. The most common approach to studying anomalous resistivity mechanisms involves wave-particle interactions with turbulent plasma waves, such as ion-acoustic turbulence, kinetic Alfvén waves, lower-hybrid drift waves, and whistler waves. Since real magnetic energy dissipation occurs at plasma kinetic scales, the motion of individual charged particles plays a crucial role. Based on the observation that strong magnetic field gradients near magnetic null points cause nonlinear effects that make charged particle trajectories chaotic, Yoshida et al. (1998) proposed a new mechanism for collisionless anomalous resistivity, suggesting that chaotic particle motion leads to randomization of current-carrying particles' directed motion. The core of this mechanism is that chaotic motion of charged particles in non-uniform electromagnetic fields creates an effective "collision," thereby generating anomalous resistivity. This resistivity produced by particle orbit chaos is called chaos-induced resistivity.

This paper primarily introduces early research and recent progress on using chaos-induced resistivity to explain anomalous resistivity in collisionless reconnection regions. Section 2 reviews early studies of chaos-induced resistivity, Section 3 presents the basic theoretical model, Section 4 introduces recent advances in finite-width current sheets and island-chain current sheets, and Section 5 provides a summary and proposes future research directions.

2 Early Research on Chaos-induced Resistivity

In 1962, Schmidt discovered that near magnetic null points where the magnetic field is spatially non-uniform, the presence of strong magnetic field gradient nonlinearity causes charged particle trajectories to become disordered and chaotic. The left panel of [Figure 1: see original paper] shows the time evolution of kinetic entropy S for chaotic motion (solid line) and periodic motion (dashed line), while the right panel displays the variation of chaos-induced resistivity (solid line) and classical collisional resistivity (dashed line) with mean free path.

Yoshida et al. (1998) pointed out that strong magnetic field variations near null points violate magnetic moment conservation, inducing chaotic motion of charged particles that introduces new degrees of freedom and ultimately leads to entropy increase, as shown in the left subplot of [Figure 1: see original paper]. They calculated chaos-induced resistivity from charged particle chaotic motion and classical collisional resistivity from Coulomb collisions for different mean free paths, finding that the effective resistivity is 1-2 orders of magnitude higher than classical collisional resistivity, as shown in the right subplot of [Figure 1: see original paper]. This was the first proposal to explain anomalous resistivity in collisionless magnetic reconnection through chaos-induced resistivity arising from particle orbit chaos.

Numata et al. recognized that energy dissipation persists only in open systems, so they studied chaos-induced resistivity in collisionless reconnection regions under open system conditions, using test-particle simulations and ensemble statistics to quantitatively calculate chaos-induced resistivity. In their numerical simulations, Numata et al. used Lyapunov exponents and local maximum Lyapunov exponents to measure the degree of chaos for single particles and multiple particles, respectively. The left subplot of [Figure 2: see original paper] shows the time evolution of the Lyapunov exponent for a single particle ($\lambda(t) \equiv (1/\Delta t)\ln(|\delta x(t+\Delta t)|/|\delta x(t)|)$, where δx is the separation between adjacent charged particle trajectories), while the right subplot displays the time evolution of the local maximum Lyapunov exponent ($\tilde{\lambda}(t) \equiv \langle \lambda(t) \rangle$) for multiple particles in chaotic regions of different radii (denoted as R in the figure). By selecting the chaotic region radius corresponding to a stable local maximum Lyapunov exponent, they statistically determined the average velocity of charged particles in the accelerating electric field direction $\bar{v}_z(t')$ and the number of charged particles in the chaotic region $n(t')$, fitting these with a linear function ($\bar{v}_z(t') = \alpha t'$, where α is the average acceleration) and a single exponential function ($n(t') = n_0 \exp(-\beta t')$, where n_0 is the initial number of charged particles and β is the relative escape rate). This yielded the average acceleration of charged particles along the electric field direction and their relative escape rate from the chaotic region. Theoretically, when the system reaches equilibrium, the saturation velocity in the accelerating electric field direction can be determined and combined with the saturation velocity obtained from the dissipation equation to give the effective collision frequency, thereby calculating the chaos-induced resistivity for specific plasma conditions. Using typical coronal plasma parameters, they computed chaos-induced resistivity for X-type magnetic configurations and found it to be four orders of magnitude higher than classical collisional resistivity.

The magnetic topology in Earth' s magnetotail often consists of X-line current sheets and bifurcated current sheets. In 2014, Andriyas et al. studied chaos-induced resistivity for protons and oxygen ions in these two magnetic configurations. They used a double exponential function ($n(t') = \sum_j n_j \exp(-\beta_j t')$; $j = 1, 2$) to fit the relative escape rate of charged particles and analyzed Poincaré sections and maximum Lyapunov exponents. The results showed maximum Lyapunov exponents of 0.34 and 0.25 for bifurcated and X-line current sheets, respectively, indicating that charged particle motion is more chaotic in bifurcated current sheets. Using typical magnetotail plasma parameters, they calculated chaos-induced resistivity in bifurcated current sheets and found it to be 9–10 orders of magnitude larger than classical collisional resistivity.

Magnetic reconnection can be classified as complete or partial. Complete reconnection involves two strictly anti-parallel magnetic fields that completely cancel at the reconnection point, forming a true magnetic null. However, most magnetic reconnection events in space plasmas are partial, where the magnetic fields at the reconnection point are not strictly anti-parallel and do not completely cancel, leaving a residual component called the guide field. In 2017, Shang et al. studied the effect of guide fields on chaos-induced resistivity in X-type magnetic configurations. The first row of [Figure 3: see original paper] shows three charged particle orbits (R_0 is the characteristic length scale of magnetic field variation): oscillatory particle orbit, magnetic drift particle orbit, and chaotic particle orbit. The second row displays the time evolution of the magnetic moment (thick line, μ_M) and the distance between the particle and magnetic null (thin line) for each orbit (τ_A is the characteristic time). The third row shows the time evolution of the Lyapunov exponent for each orbit. The results show that the Lyapunov exponent for oscillatory orbits decays to below 0.05, for magnetic drift orbits stabilizes near 0.05, and for chaotic orbits stabilizes above 0.1, indicating that chaotic particle orbits contribute most significantly to chaos-induced resistivity.

Furthermore, Shang et al. investigated the variation of chaos-induced resistivity with guide field strength in multi-particle systems. [Figure 4: see original paper] shows the average acceleration, relative escape rate, and effective collision frequency as functions of guide field for different accelerating electric fields. The results reveal that when the guide field is half the background magnetic field ($B_z = 0.5B_0$), the effective collision frequency in the collisionless reconnection region reaches its maximum, corresponding to the highest chaos-induced resistivity.

3 Basic Theoretical Model of Chaos-induced Resistivity

To quantitatively calculate chaos-induced resistivity from chaotic charged particle motion, multiple charged particles must be considered through test-particle simulations and ensemble statistics. In the simulations, the magnetic field is given by $\mathbf{B} = B_0(b_x, b_y, \epsilon)$, where $b_x = B_x/B_0$, $b_y = B_y/B_0$, and ϵ is a constant representing the relative strength of a constant guide field. The x-y plane is the reconnection plane, z is the guide field direction, and B_x, B_y, B_z denote the magnetic field strengths in the reconnection plane and guide field direction, respectively. A uniform, constant accelerating electric field is applied along the guide field direction, $\mathbf{E} = (0, 0, E_z)$. Particle collisions and collective interactions are neglected, so the equation of motion for a single charged particle is

$$\frac{d\mathbf{r}}{dt} = \mathbf{v}, \quad \frac{d\mathbf{v}}{dt} = \frac{q}{m}(\mathbf{E} + \mathbf{v} \times \mathbf{B}),$$

where $\mathbf{r}$ and $\mathbf{v}$ are the particle position and velocity, and q and m are the charge and mass, respectively. For numerical convenience, the following normalized parameters are used to normalize the equations of motion, magnetic field, and electric field:

$$\mathbf{B}' = \mathbf{B}/B_0, \quad \mathbf{r}' = \mathbf{r}/R_0, \quad t' = t/\tau_A, \quad \tau_A = R_0/v_A, \quad \mathbf{v}' = \mathbf{v}/v_A, \quad E_z' = E_z/v_A B_0 = M_A,$$

where B_0 is the characteristic reconnection magnetic field strength, v_A is the Alfvén speed, and M_A is the relative strength of the accelerating electric field. The simulations adopt $R_0 = v_A/\omega_c$, where ω_c is the cyclotron frequency of charged particles in the electromagnetic field.

Using these normalized parameters, the normalized equation of motion becomes

$$\frac{d\mathbf{r}'}{dt'} = \mathbf{v}', \quad \frac{d\mathbf{v}'}{dt'} = M_A + \mathbf{v}' \times \mathbf{B}',$$

with normalized magnetic and electric fields $\mathbf{B}' = (b_x, b_y, \epsilon)$ and $\mathbf{E}' = (0, 0, M_A)$.

In test-particle simulations, the average velocity of charged particles in the z-direction within the chaotic region and the remaining number of particles in the chaotic region can be obtained. Linear and single exponential functions are used to fit the time evolution of the average particle velocity in the z-direction and the number of particles in the chaotic region, respectively. [Figure 5: see original paper] shows the time evolution of the average velocity in the electric field direction and the number of charged particles in the chaotic region (the hat notation ˆ indicates normalization, following the convention in reference [40] and equivalent to the prime notation).

$$\bar{v}_z'(t') = \alpha t', \quad n(t') = n_0 \exp(-\beta t'),$$

yielding the acceleration α of charged particles along the electric field direction and the escape rate β from the chaotic region, where $\bar{v}_z'$ and $n(t')$ are the average velocity in the z-direction and the particle number density in the chaotic region, respectively.

Theoretically, steady-state reconnection is achieved by continuously injecting new zero-velocity particles into the chaotic region. After injection, the total average velocity of charged particles in the chaotic region becomes

$$\bar{V}_z'(t) = \frac{\int_0^t \bar{v}_z'(t') n_0 \beta \exp(-\beta t') dt' + n_0 \exp(-\beta t') \bar{v}_z'(t')}{\int_0^t n_0 \beta \exp(-\beta t') dt' + n_0 \exp(-\beta t')}.$$

The saturation velocity can also be calculated from the dissipation equation

$$\frac{d\bar{V}_z'(t)}{dt'} = M_A - \nu_{\text{eff}}' \bar{V}_z'(t),$$

where ν'_{eff} is the dimensionless effective collision frequency. When $d/dt \to 0$, the saturation velocity is

$$\bar{V}'_z(t')|_{d/dt\to 0} = \frac{M_A}{\nu'_{\mathrm{eff}}}.$$

Simplifying the expression for $\bar{V}'_z(t')$ gives

$$\bar{V}'_z(t') = \frac{\alpha}{\beta}[1 - \exp(-\beta t')].$$

When $t' \to \infty$, the saturation velocity at steady state is

$$\bar{V}'_z(t')|_{t'\to\infty} = \frac{\alpha}{\beta}.$$

Equating the two saturation velocities from equations (8) and (10) yields the effective collision frequency

$$\nu'_{\mathrm{eff}} = \frac{M_A \beta}{\alpha}.$$

Therefore, the chaos-induced resistivity in a specific plasma environment can be expressed as

$$\eta_{\mathrm{eff}} = \frac{m\nu_{\mathrm{eff}}}{n_0 e^2} = \frac{m\omega_c \nu'_{\mathrm{eff}}}{n_0 e^2}.$$

4.1 Chaos-induced Resistivity in Finite-Width Current Sheets

When studying chaos-induced resistivity in finite-width current sheets, the normalized magnetic field configuration $\mathbf{B}'$ is expressed as

$$\mathbf{B}' = \begin{cases} (\tanh(y'), \tanh(x' - l/2), \epsilon), & x' > l/2, \\ (\tanh(y'), \tanh(x' + l/2), \epsilon), & x' \le -l/2, \\ (\tanh(y'), 0, \epsilon), & |x'| < l/2, \end{cases}$$

where the x-y plane is the reconnection plane, z is the guide field direction, and l is the distance between two Y-type neutral points. When $l = 0$, the configuration is an X-type magnetic field; when $l \gg 1$, it becomes a one-dimensional Hall current sheet; and when $0 < l < 1$, a finite-width current sheet of width lR_0 exists, known as a double Y-type magnetic configuration.

In the numerical simulations, 2×10^5 charged particles are used. For $l = 0$ (X-type configuration), particles are initially uniformly distributed in $-1.0 <$

$(x', y') < 1.0$ with $z' = 0$, and have an isotropic Maxwellian velocity distribution $f(v) = (2\pi v_T^2)^{-3/2} \exp(-v^2/v_T^2)$, where $v_T = 0.3v_A$. [Figure 6: see original paper] shows the average acceleration α, relative escape rate β, and effective collision frequency ν'_{eff} as functions of the chaotic region radius R_c/R_0 for a fixed guide field $B_z = 0.5B_0$ in the X-type configuration. Based on the chaotic region size, three zones are identified: the non-magnetized zone ($R_c < R_0$), the transition zone ($R_c \approx R_0$), and the magnetized zone ($R_c > R_0$). In the non-magnetized zone, the average acceleration α increases with R_c/R_0 while the relative escape rate β remains nearly constant, causing ν'_{eff} to decrease with increasing R_c/R_0. In the magnetized zone, both α and β decrease with R_c/R_0, leading to an increase in ν'_{eff} with R_c/R_0. Moreover, chaos-induced resistivity in both the non-magnetized ($R_c < R_0$) and magnetized ($R_c > R_0$) zones is higher than in the transition zone ($R_c \approx R_0$).

For $l = 8$ (double Y-type configuration), particles are uniformly distributed in $-5.0 < x' < 5.0$, $-1.0 < y' < 1.0$, $z' = 0$ with an isotropic Maxwellian velocity distribution. [Figure 7: see original paper] illustrates the double Y-type magnetic configuration and three different types of chaotic regions (shaded areas), which from top to bottom are designated as Y-type, S-type, and O-type chaotic regions.

[Figure 8: see original paper] shows α, β, and ν'_{eff} as functions of R_c/R_0 for the three chaotic regions in the double Y-type configuration, with the top, middle, and bottom rows corresponding to Y-type, S-type, and O-type regions, respectively. The results show that α and β are primarily determined by the accelerating electric field M_A, consistent with Shang et al. (2017). Both α and β decrease with R_c/R_0 in the magnetized zone, while being less affected by R_c/R_0 in the non-magnetized zone, mainly because the reconnection field influences particle gyration, hindering acceleration and escape. Furthermore, ν'_{eff} varies non-uniformly with R_c/R_0 across the three chaotic regions: it increases with R_c/R_0 in Y-type regions but decreases in O-type regions, demonstrating significant spatial variation of chaos-induced resistivity across different chaotic zones in reconnection current sheets.

Using typical pre-flare coronal plasma parameters ($B_0 \approx 100$ Gs, plasma density $n_0 \approx 10^{16}$ m^{-3}, temperature $T \approx 500$ eV), presents the weighted average effective collision frequencies for non-magnetized and magnetized cases in X-type, Y-type, S-type, and O-type chaotic regions. The chaos-induced resistivity η_{eff} is calculated using equation (12). Classical collisional resistivity is given by $\eta_{\text{cal}} = 5.2 \times 10^{-5} Z \ln\Lambda / T^{3/2}$ $[\Omega \cdot \text{m}]$, where Z is the ion charge number and the Coulomb logarithm is $\ln\Lambda \approx 10$. The results show that chaos-induced resistivity in pre-flare coronal plasma is 5-7 orders of magnitude larger than classical collisional resistivity, confirming that chaos-induced resistivity can provide effective dissipation in collisionless reconnection regions. Additionally, chaos-induced resistivity in X-type chaotic regions is 1-2 orders of magnitude higher than in O-type regions, indicating non-uniform spatial distribution of chaos-induced resistivity.

4.2 Chaos-induced Resistivity in Island-Chain Current Sheets

Numerous theoretical analyses and simulations have shown that finite-width reconnection current sheets are often unstable to resistive tearing modes and can evolve into island-chain structures with alternating X-type and O-type neutral points. To study chaos-induced resistivity near X-type and O-type neutral points in island-chain current sheets, the normalized magnetic field configuration is given by

$$\mathbf{B}' = \left(\frac{\sinh(y'/2)}{\cosh(y'/2) + \epsilon \cos(x'/2)}, \frac{\epsilon \sin(x'/2)}{\cosh(y'/2) + \epsilon \cos(x'/2)}, \epsilon \right),$$

where lengths are normalized by the particle inertial length $R_A = v_A/\omega_c$, and the magnetic island parameter ϵ determines the strength and configuration of magnetic islands in the chain structure. Here, R_A represents the scale of magnetic field inhomogeneity. For $\epsilon = 0.4$, the magnetic configuration is shown in [Figure 9: see original paper], with the shaded areas in the top and bottom panels representing X-type and O-type regions, respectively.

In the simulations, 2×10^5 particles are used, initially uniformly distributed in $-1.0 < x', y' < 1.0$, $z' = 0$ with an isotropic Maxwellian velocity distribution of thermal speed $v_T = 0.3v_A$. The chaotic region radius near X-type and O-type neutral points is set to $R_c = 2.0R_A$. Figure 10: see original paper and (b) show α, β, and ν'_{eff} as functions of guide field for $\epsilon = 0.4$ and $\epsilon = 0.8$, respectively, with the top and bottom rows in each panel corresponding to X-type and O-type regions. The average acceleration α increases with M_A in all cases, indicating that particle acceleration is primarily determined by the electric field. Strong guide fields cause particles to remain longer in chaotic regions, making acceleration more effective. However, in X-type regions with weak guide fields ($B_z < B_0$), α exhibits complex dependence on B_z/B_0. The relative escape rate β in O-type regions decreases sharply with increasing guide field, showing that strong guide fields effectively inhibit particle escape from O-type regions. In X-type regions, β shows little variation for $B_z > 0.5B_0$, especially for $M_A = 0.001$ and 0.0005. For $B_z < 0.5B_0$, β in X-type regions also shows complex behavior similar to α, leading to complex variation of ν'_{eff} with guide field. In O-type regions, ν'_{eff} remains essentially constant with increasing guide field for $\epsilon = 0.4$ but decreases for $\epsilon = 0.8$. Comparing Figure 10: see original paper and (b) reveals that except for weak guide fields ($B_z < 0.5B_0$), ν'_{eff} (i.e., chaos-induced resistivity η_{eff}) is generally larger in X-type regions than in O-type regions. For weak guide fields and low magnetic island parameter ($\epsilon = 0.4$), ν'_{eff} in X-type regions approaches or falls below that in O-type regions.

For a fixed guide field $B_z = 1.0B_0$ and accelerating electric field $M_A = 0.001$, [Figure 11: see original paper] shows α, β, and ν'_{eff} as functions of ϵ. The average acceleration α in O-type regions increases significantly with ϵ, while showing little variation in X-type regions, indicating that acceleration along the guide field can be hindered by strong magnetic islands. Both β and ν'_{eff} in O-type regions increase slightly with ϵ for weak islands ($\epsilon < 0.5$) but decrease rapidly for strong islands ($\epsilon > 0.5$). Furthermore, ν'_{eff} (i.e., η_{eff}) in X-type regions is about one order of magnitude higher than in O-type regions for weak islands ($\epsilon < 0.5$), and about two orders of magnitude higher for strong islands ($\epsilon > 0.5$).

These results demonstrate that chaos-induced resistivity plays an important role in collisionless reconnection current sheets and may be much larger than classical collisional resistivity in typical space plasma environments. For example, using typical coronal parameters ($B_0 \approx 500$ Gs, plasma density $n_0 \approx 10^9$ cm^{-3}, temperature $T \approx 200$ eV), lists the effective collision frequencies ν'_{eff} in X-type and O-type regions for different ϵ values, along with the ratios of chaos-induced resistivity to classical collisional resistivity ($\eta_{X\text{eff}}/\eta_{X\text{cal}}$ and $\eta_{O\text{eff}}/\eta_{O\text{cal}}$). The ratios are approximately 10^7 in X-type regions and 10^4–10^6 in O-type regions, satisfying the anomalous resistivity requirements for solar flares. The last column of gives the ratio $\eta_{X\text{eff}}/\eta_{O\text{eff}}$, showing that chaos-induced resistivity in X-type regions is 1–2 orders of magnitude higher than in O-type regions, indicating non-uniform distribution of chaos-induced resistivity concentrated in X-type regions of island-chain current sheets.

5 Summary and Outlook

Collisionless magnetic reconnection is widely studied as an effective mechanism for magnetic energy release and is closely related to many explosive phenomena observed in solar, space, and laboratory plasmas. However, an unresolved problem in collisionless magnetic reconnection is the generation mechanism of anomalous resistivity. The concept of explaining anomalous resistivity through chaos-induced resistivity from particle orbit chaos was proposed by Yoshida et al. (1998). Numata et al. (2002, 2003) further developed quantitative calculations for chaos-induced resistivity in open systems. Andriyas et al. (2014) found through simulations that chaos-induced resistivity may significantly contribute to anomalous resistivity in magnetotail reconnection current sheets. Shang et al. (2017) extended the work of Numata et al. to cases with non-zero guide fields. In the most recent work, Wang et al. investigated chaos-induced resistivity in finite-width current sheets and island-chain current sheets, finding that chaos-induced resistivity is much larger than classical collisional resistivity and is non-uniformly distributed in reconnection regions, which matches the requirements for anomalous resistivity. However, current studies have only examined chaos-induced resistivity in fixed electromagnetic fields, whereas real space plasma environments involve much more complex fields. Important unanswered questions remain: How do perturbation fields in magnetic reconnection affect individual particle trajectories? Will chaos-induced resistivity increase? How much can chaos-induced resistivity contribute to the anomalous resistivity required for magnetic reconnection in full particle simulations? These issues warrant further attention and research.

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Note: Figure translations are in progress. See original paper for figures.