Shifting and Splitting of Resonance Lines due to Dynamical Friction in Plasmas

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A quasilinear plasma transport theory that incorporates Fokker-Planck dynamical friction (drag) and pitch angle scattering is self-consistently derived from first principles for an isolated, marginally unstable mode resonating with an energetic minority species. It is found that drag fundamentally changes the structure of the wave-particle resonance, breaking its symmetry and leading to the shifting and splitting of resonance lines. In contrast, scattering broadens the resonance in a symmetric fashion. Comparison with fully nonlinear simulations shows that the proposed quasilinear system preserves the exact instability saturation amplitude and the corresponding particle redistribution of the fully nonlinear theory. Even in situations in which drag leads to a relatively small resonance shift, it still underpins major changes in the redistribution of resonant particles. This novel influence of drag is equally important in plasmas and gravitational systems. In fusion plasmas, the effects are especially pronounced for fast-ion-driven instabilities in tokamaks with low aspect ratio or negative triangularity, as evidenced by past observations. The same theory directly maps to the resonant dynamics of the rotating galactic bar and massive bodies in its orbit, providing new techniques for analyzing galactic dynamics.

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Introduction.—Resonances are channels for nonadiabatic energy exchange between particles and waves in kinetic systems. In linear theory, resonant interactions occur when a specific synchronization condition is satisfied exactly. In reality, however, this condition is smeared out due to effects such as finite mode amplitude, turbulence, and collisions. In plasmas, the effect of resonance broadening has been historically investigated in strong turbulence [1–5] and quasilinear [6–11] approaches, with important implications to the dynamics of plasma echoes [12] and Alfvén waves [13]. Previously unexplored aspects, but similarly important for influencing the resonant particle relaxation, are mechanisms responsible for the shifting and splitting of discrete resonances.

The competition between convective (dynamical friction, also known as drag) and diffusive (scattering) collisions plays a key role in determining the dynamical behavior of wave-particle resonant systems [14]. For instance, it is crucial to explaining the observation of Alfvén eigenmode frequency chirping in tokamaks [15]. Although global nonresonant effects of drag are well known [16], the influence of drag on the structure of narrow resonant layers has not previously been determined.

In this Letter, we explore kinetic instabilities close to their threshold to analytically show that coherent Fokker-Planck drag breaks the symmetry of the resonances with respect to their original location. This occurs not only by shifting and splitting them, but also by altering the relative strength of regions in the vicinity of a resonance. The skewed dependence of the resonance due to drag is, in turn, responsible for significant modifications to the particle distribution function. Moreover, it is demonstrated that the fully nonlinear collisional kinetic system naturally reduces to the form of a quasilinear (QL) theory in the limit of marginal stability and when stochastic processes dominate the relaxation.

An important application of these findings in fusion plasmas is the forecasting of deleterious fast ion transport by Alfvén eigenmodes (AE). Fully nonlinear simulations are numerically costly and, therefore, it is of practical interest to develop reduced models, such as quasilinear theory, capable of reproducing essential features of more complete descriptions [17,18]. The drag-induced modifications of the instability saturation level, resonance function, and fast ion redistribution are anticipated to be more significant in tokamaks with low aspect ratio or negative triangularity, as evidenced by past observations of a greater propensity for dynamical Alfvén eigenmode behavior in these configurations [19–23], consistent with theoretical predictions [24–26].

A deep connection exists between kinetic processes in plasmas and those present in self-gravitating systems, as both are well described by mean field theories governed by long-range, inverse square laws. In fact, the respective distributions obey the same evolution equation in phase space, giving rise to analogous phenomena such as Landau damping and resonant relaxation in both plasmas [27–29] and self-gravitating systems [30–33]. Consequently, understanding the structure of collisional wave-particle resonances is equally relevant to kinetic plasma instabilities and galactic dynamics, for instance, in determining the torque applied to the rotating galactic bar by orbiting heavy bodies [34].

Self-consistent transport theory with dynamical friction and scattering.—Krook or scattering collisions are known to lead to an antisymmetric modification of the particle distribution around a resonance [9]. As will be demonstrated, drag introduces a distinct, asymmetric response. Hence, we investigate the 1D nonlinear dynamics of an electrostatic wave resonating with a hot minority species, i.e., the canonical bump-on-tail problem, in the presence of both scattering and drag. The generalization to more complicated geometries and to waves of any polarization can be achieved with the methods in Refs. [9,35]. The wave field is represented as $E(x,t) = \mathrm{Re}[\hat{E}(t)e^{i(kx-\omega t)}]$, with complex amplitude $\hat{E}(t)$. The hot minority species is described via a distribution function $f(x,v,t)$ with an initial discrete resonance at $kv_{\mathrm{res}} = \omega$. For resonances narrow with respect to the velocity scale of the distribution, it is sufficient to evaluate the Fokker-Planck coefficients at the resonance center, and the kinetic equation becomes [14,35]

$$\frac{\partial f}{\partial t} + v\frac{\partial f}{\partial x} + \frac{1}{k}\mathrm{Re}\left[\omega_b^2 e^{i(kx-\omega t)}\right]\frac{\partial f}{\partial v} = C[f - F_0], \quad (1)$$

where

$$C[g] = \frac{\nu^3}{k^2}\frac{\partial^2 g}{\partial v^2} + \frac{\alpha^2}{k}\frac{\partial g}{\partial v}. \quad (2)$$

Here, $F_0(v)$ is the equilibrium distribution function, which is assumed to have an approximately constant positive slope in the vicinity of the resonance. The rate of mode drive at $t = 0$ in the absence of damping is proportional to this slope: $\gamma_{L,0} = (2\pi^2 q^2\omega/mk^2)\partial F_0/\partial v$, while the mode damping rate due to interaction with the background thermal plasma is given by a constant γ_d. The nonlinear bounce frequency for deeply trapped resonant particles, $\omega_b(t) \equiv \sqrt{qk\hat{E}(t)/m}$, is a convenient measure of the mode amplitude.

The coefficients ν and α are the effective scattering and drag collision frequencies, respectively, which are enhanced with respect to the 90° pitch angle scattering rate and the inverse slowing down time [7,12,25,36] (see Ref. [37] for heuristic arguments). Because of the narrowness of the resonances, the resonant particle dynamics can be strongly controlled by collisions even if their nonresonant dynamics hardly feel their effect.

Because of spatial periodicity, the distribution can be assumed of the form $f(x,v,t) = f_0(v,t) + \sum_{l=1}^{\infty}[f_l(v,t)e^{il(kx-\omega t)} + \mathrm{c.c.}]$. In the collisional regime, an asymptotic expansion exists in orders of the small parameter $|\omega_b^2|/\nu^2 \ll 1$. Rewriting Eq. (1) in orders of such a small parameter naturally implies the ordering $|f_0'^{(0)}| \gg |f_1'^{(1)}| \gg |f_0'^{(2)}|, |f_2'^{(2)}|$ [39]. The prime denotes the derivative with respect to v, while the superscript denotes the order in the wave amplitude (proportional to $|\omega_b^2|/\nu^2$). Note that $f_0^{(0)}$ corresponds to the equilibrium distribution F_0. This ordering of the perturbation theory is satisfied for the entire time evolution of the system so long as (i) steady asymptotic solutions exist [guaranteed when $\alpha/\nu < 0.96$ and $\nu > \nu_{\mathrm{crit}} \approx 2.05(\gamma_{L,0} - \gamma_d)$ [14]] and (ii) the system is sufficiently close to marginal stability, i.e., $\gamma_{L,0} - \gamma_d \ll \gamma_{L,0}$. When conditions (i) and (ii) are not satisfied, the mode may grow to an amplitude that violates the $|\omega_b^2| \ll \nu^2$ assumption. From Eq. (1), the f_l satisfy

$$\frac{\partial f_l}{\partial t} + il(kv - \omega)f_l + \frac{(\omega_b^2 f_{l-1}' + \omega_b^{2*} f_{l+1}')}{2k} = C[f_l]. \quad (3)$$

In general, the solution of Eq. (3) is an integral over the time history of the system that involves delays in the argument of $\omega_b(t)$. However, when stochastic collisions dominate the system's dynamics, the system's memory is poorly retained, and the dynamics instead become essentially local in time [38,40]. This occurs when $\nu \gg \gamma_{L,0} - \gamma_d$. The constraints on these parameters become more restrictive as the ratio of drag to scattering is increased [37]. Then, to first order in $|\omega_b^2|/\nu^2$, one can disregard the time derivative in Eq. (3), yielding

$$f_1^{(1)} = -\frac{F_0'\omega_b^2(t)}{2k\nu}\int_0^\infty ds\, e^{-i\left(\frac{kv-\omega}{\nu}\right)s - i\frac{\alpha^2 s^2}{2\nu^2} - \frac{s^3}{3}}. \quad (4)$$

Using the reality rule $f_{-1}^{(1)} = f_1^{(1)*}$, Eq. (3) can be written to second order in $|\omega_b^2|/\nu^2$:

$$\frac{\partial f_0^{(2)}}{\partial t} + \frac{1}{2k}(\omega_b^2[f_1'^{(1)}]^* + \omega_b^{2*}f_1'^{(1)}) = C[f_0^{(2)}]. \quad (5)$$

Substitution of Eq. (4) into Eq. (5) shows that only the spatially averaged distribution $f(v,t) \equiv (k/2\pi)\times\int_0^{2\pi/k} f(x,v,t)dx = F_0(v) + f_0^{(2)}(v,t)$ explicitly appears in the wave-particle power exchange to second order in $|\omega_b^2|/\nu^2$. Since $\partial F_0/\partial t = 0$ and $|F_0'| \gg |f_0'^{(2)}|$, one then obtains from Eq. (5) that the resonant particle response is regulated by a QL diffusion-advection transport equation:

$$\frac{\partial f(v,t)}{\partial t} - \frac{\partial}{\partial v}\left[\frac{\pi}{2k^3}|\omega_b^2|^2\mathcal{R}(v)\frac{\partial f}{\partial v}\right] = C[f - F_0], \quad (6)$$

with the resonance function given by

$$\mathcal{R}(v) = \frac{k}{\pi\nu}\int_0^\infty ds \cos\left(\frac{(kv-\omega)}{\nu}s + \frac{\alpha^2}{\nu^2}\frac{s^2}{2}\right)e^{-s^3/3}. \quad (7)$$

$\mathcal{R}(v)$ gives the velocity-dependent strength of the resonant wave-particle interaction. In the absence of collisions, $\mathcal{R}(v) = \delta(v - \omega/k)$ would describe an exact, unbroadened resonance. For all values of α/ν, the property

$\int_{-\infty}^{\infty} \mathcal{R}(v)dv = 1$ is exactly satisfied. The complete set of equations describing the QL system is then given by Eqs. (6) and (7), coupled with the amplitude evolution equation

$$\frac{d|\omega_b^2(t)|^2}{dt} = 2[\gamma_L(t) - \gamma_d]|\omega_b^2(t)|^2, \tag{8}$$

where

$$\gamma_L(t) = \frac{2\pi^2 q^2 \omega}{mk^2} \int_{-\infty}^{\infty} dv \mathcal{R}(v) \frac{\partial f(v,t)}{\partial v}. \tag{9}$$

A first-principles derivation of Eqs. (8) and (9) is shown in Supplemental Material [37].

Remarkably, a QL transport theory has emerged spontaneously from the nonlinear one under the assumption that stochasticity dominates over the relaxation timescale ($\nu \gg \gamma_{L,0} - \gamma_d$), even in the presence of coherent drag which acts to preserve phase space correlations. Furthermore, the QL system was obtained for an isolated, discrete resonance without invoking any overlapping condition.

Shift of the resonance function due to symmetry breaking.—Plotting the resonance function $\mathcal{R}(v)$ for different values of α/ν, as done in Fig. 1(a), demonstrates that drag shifts the location of the strongest wave-particle interaction in phase space. Moreover, drag acts to increase the strength of the interaction downstream of the resonance and diminish it upstream. Although the quantitative changes in the resonance function due to drag are seemingly small, this fundamental change in the character of the wave-particle resonance has substantial consequences for how the energetic particles are redistributed. This shift is not analogous to a simple Doppler shift but rather the result of asymmetry present in the collisional dynamics due to drag, which has a preferred direction. Without drag, $\mathcal{R}(v)$ is perfectly symmetric about $kv - \omega = 0$.

The shift of the resonance function implies a previously unrecognized collisional modification to the resonance condition, which can be obtained by calculating the location of the peak of $\mathcal{R}(v)$, leading to $\int_0^\infty \sin(s[(kv-\omega)/\nu]_{\rm peak} + \alpha^2 s^2/2\nu^2) s \exp(-s^3/3) ds = 0$. Noting that only small s contributes due to the strongly decaying cubic exponential, the drag-modified resonance condition becomes $(kv-\omega)_{\rm peak} \approx -3^{1/3}\Gamma(4/3)\alpha^2/2\nu$, which is accurate to within 1.5% for $\alpha/\nu < 1$. Shifted resonance lines, captured here using drag in a considerably reduced theory, are also known to exist in the strong turbulence framework [49] as a result of turbulence modification of ensemble average orbits.

Particle relaxation.—The derived QL system also reproduces key features of the complete nonlinear system, namely, the perturbed distribution function ($\delta f \equiv f - F_0$) and wave saturation amplitude. When stochasticity regulates the timescale for the mode growth, Eq. (6) can be further simplified to $(\nu^3/k)\delta f' + \alpha^2 \delta f = -(\pi|\omega_b^2|^2/2k^2) \times \mathcal{R}(v)F_0'$, which can be solved for by direct integration:

$$\delta f(v,t) = \frac{|\omega_b^2(t)|^2 F_0'}{2k\nu^3}\left\{c\left(\frac{\alpha}{\nu}\right) - \int_0^\infty \frac{ds e^{-s^3/3}}{\alpha^4/\nu^4 + s^2}\left[\frac{\alpha^2}{\nu^2}\cos\left(\frac{(kv-\omega)s}{\nu} + \frac{\alpha^2}{\nu^2}\frac{s^2}{2}\right) + s\sin\left(\frac{(kv-\omega)s}{\nu} + \frac{\alpha^2}{\nu^2}\frac{s^2}{2}\right)\right]\right\}. \tag{10}$$

The integration constant $c(\alpha/\nu)$ is determined by enforcing particle conservation: $\int_{-v_{\max}}^{v_{\max}} \delta f dv = 0$. The velocity structure of δf does not evolve in time, as all time dependence is contained in the overall factor $|\omega_b^2(t)|^2$. The relaxed distribution is plotted in Fig. 1(b) for several values of $\alpha/\nu < 1$. Remarkably, a small quantitative asymmetry in the collisional dynamics due to drag can

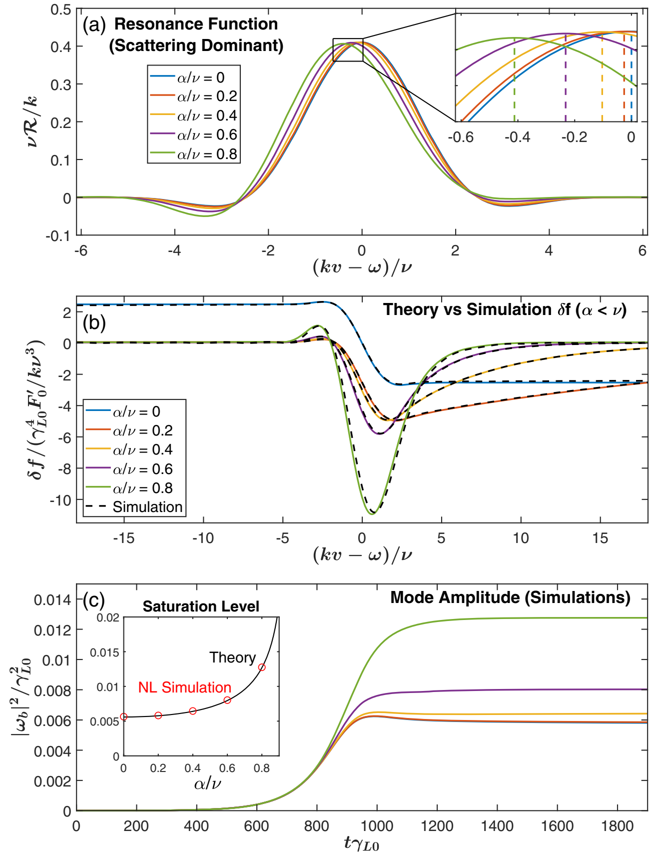

FIG. 1. (a) Resonance function [Eq. (7)], (b) saturated distribution modification with respect to the initial equilibrium distribution, $\delta f = f - F_0$ [Eq. (10)], and (c) time evolution of the mode amplitude (proportional to $|\omega_b^2|$). In (a) and (b), solid curves represent analytic expressions, while dashed curves are simulation results from the nonlinear Vlasov code BOT using $\nu/(\gamma_{L,0} - \gamma_d) = 20$ and $\gamma_d/\gamma_{L,0} = 0.99$. Colored curves in (c) are simulation results. The inset plots the saturation level from these simulations against the analytic prediction.

have a large qualitative effect on the saturated distribution, even though the corresponding changes in the resonance function are less dramatic. When no drag is present ($\alpha = 0$), δf is antisymmetric with constant plateaus outside of a narrow transition region near the peak of the resonance. As the ratio of drag to scattering is increased, the plateau upstream of the resonance ($kv > \omega$) instead decays in velocity space at a rate proportional to α^2/ν^2. Particle conservation shifts the entire distribution downward once the symmetry is broken, eliminating nearly all of the downstream particle redistribution, even for very small amounts of drag. The drag-induced modifications to δf were compared against simulations performed with the nonlinear 1D Vlasov code BOT, which solves the plasma kinetic equation [Eq. (1)] directly [24,50]. The dashed black curves in Fig. 1(b) show the simulation results for each value of α/ν, demonstrating excellent agreement.

Instability saturation level.—Drag has a destabilizing effect on the underlying instability, leading the wave to saturate at a larger amplitude with increasing α/ν. An analytically tractable example is the first-order correction to the saturation amplitude due to drag. Substituting Eq. (10) into Eq. (9), to lowest order in α^2/ν^2, one finds (the details of the derivation are given in Supplemental Material [37])

$$\gamma_L(t) \simeq \gamma_{L,0}\left\{1 - \frac{|\omega_b^2(t)|^2}{2\nu^4}\left[\Gamma\left(\frac{1}{3}\right)\left(\frac{3}{2}\right)^{1/3}\frac{1}{3} - \frac{\pi\alpha^2}{2\nu^2}\right]\right\}. \quad (11)$$

At saturation, i.e., when $\gamma_L(t) = \gamma_d$, then the mode amplitude $\hat{E}$ is given by

$$|\omega_{b,\mathrm{sat}}| = \sqrt{\frac{qk\hat{E}_{\mathrm{sat}}}{m}} \simeq \left[\frac{2(1-\gamma_d/\gamma_{L,0})}{\Gamma(\frac{1}{3})(\frac{3}{2})^{1/3}\frac{1}{3} - \frac{\pi}{2}\frac{\alpha^2}{\nu^2}}\right]^{1/4}\nu, \quad (12)$$

which is the same as calculated directly from nonlinear theory (Ref. [38]). In the absence of drag, the saturation in Eq. (12) recovers the level calculated in Ref. [35]. While not possible to express in terms of elementary functions, it can nonetheless be proven analytically that the QL saturation level also reproduces the fully nonlinear one whenever a steady state solution exists.

Figure 1(c) shows the time evolution of the mode amplitude in fully nonlinear simulations. The saturation levels are in excellent agreement with the unapproximated analytic expression [37] and demonstrate that even moderate amounts of drag can appreciably increase the saturated mode amplitude. For example, $\alpha/\nu = 0.6$ yields a 40% increase in the saturated amplitude relative to if drag were neglected. Consequently, this increase in mode amplitude enhances the resonant particle transport, as shown in Fig. 1(b). As α/ν is increased from 0.6 to 0.8, the magnitude of δf nearly doubles.

Resonance splitting due to large drag.—It is also interesting to consider the structure of the resonance when drag dominates over scattering, corresponding to instabilities which will eventually reach a strongly nonlinear regime, with the potential for substantial transport. Although for $\alpha/\nu > 0.96$ no steady state solution is allowed within the near-threshold perturbation theory [14], all of the derivations to this point nonetheless remain valid during the early growth phase, up until the mode amplitude exceeds the assumed $|\omega_b^2|/\nu^2 \ll 1$ ordering.

For $\alpha \gg \nu$, a pronounced splitting of the resonance function occurs, as shown in Fig. 2(a). Simultaneously, the resonance function broadens beyond its original width, proportional to ν, instead becoming proportional to α. It can be shown from Eq. (7) that, for $kv < \omega$ in the limit of $(kv-\omega)^2 \gg \alpha^2 \gg \nu^2 \gg \gamma^2$, the resonance function has the following asymptotic behavior:

$$\mathcal{R}(v) = \frac{1}{\alpha}\sqrt{\frac{2}{\pi}}\exp\left[\frac{1}{3}\left(\frac{\nu}{\alpha}\frac{kv-\omega}{\alpha}\right)^3\right]\cos\left[\frac{(kv-\omega)^2}{2\alpha^2} - \frac{\pi}{4}\right]. \quad (13)$$

Hence, the resonance function's extrema are given by $(kv-\omega)_{\mathrm{crit}} = -\alpha\sqrt{2\pi}\sqrt{n+1/4}$ for $n \geq 0$. In particular, $n = 0$ gives the central shift in the resonance function, showing that it becomes proportional to α when drag dominates instead of α^2 when scattering dominates. Strictly speaking, similar extrema are present even when $\alpha \ll \nu$; however, they are substantially smaller than the primary peak.

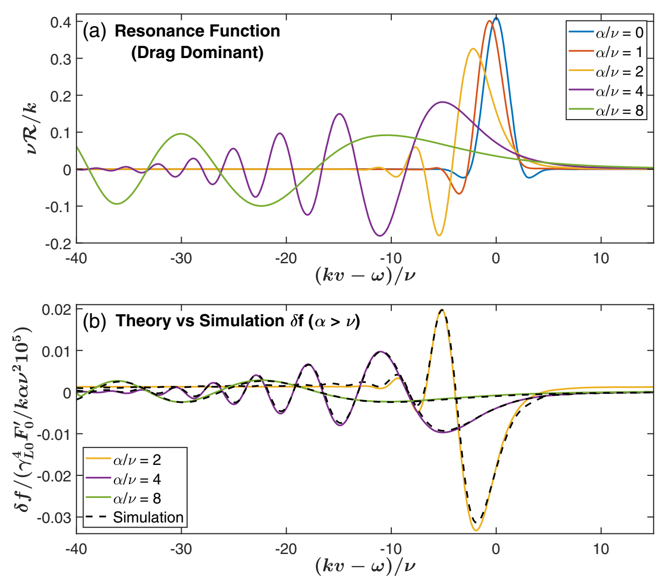

FIG. 2. (a) Resonance function [Eq. (7)] and (b) saturated distribution modification with respect to the initial equilibrium distribution, $\delta f = f - F_0$. The dashed curves in (b) were produced at $t\gamma_{L,0} = 300$ from nonlinear BOT simulations with $\nu/(\gamma_{L,0} - \gamma_d) = 20$, $\gamma_d/\gamma_{L,0} = 0.99$, and initial condition $\omega_b/\gamma_{L,0} = 10^{-8}$, while the solid curves represent Eq. (10).

The splitting of the resonance function is a novel quasilinear effect not previously identified in plasmas. This behavior implies that for sufficiently large drag there can be multiple regions of phase space where particles are efficiently interacting with the wave, even for the idealized case of a monochromatic wave with a uniform background.

During the early growth phase in the drag-dominated regime, δf has identical asymptotic behavior as the resonance function: $\delta f(v) \propto -\mathcal{R}(v)$, inheriting the same pattern of decaying oscillations. Strong agreement is found between the analytic δf calculated with Eq. (10) and nonlinear BOT simulations with $\alpha \gg \nu$, as shown in Fig. 2(b) for the cases with clear splitting. Consequently, large amounts of drag act to extend the downstream region of phase space where significant transport occurs.

Relevance to fusion plasmas.—The effects of drag derived in this Letter are most prominent when drag approaches or exceeds scattering within a narrow resonance. A heuristic scaling useful for guiding intuition is given by $\alpha/\nu \sim \mathcal{E}_{\mathrm{res}}^{1/2} n_e^{1/6}/T_e^{3/4}\omega^{1/6}$ [38]. Quantitatively, α/ν can be numerically calculated with rigor in fusion plasmas using a kinetic equilibrium reconstruction, solving for the eigenmode structure with a magnetohydrodynamics (MHD) code, following realistic orbits, and averaging over resonant surfaces. This method is outlined in Supplemental Material [37] and was previously applied to NSTX, DIII-D, and ITER [15,25,26], including the effect of enhanced diffusion due to microturbulence [41,42].

An increase in α/ν due to the reduction of turbulence has been previously identified as a reliable indicator of the onset of chirping behavior for Alfvénic modes [15], which is ubiquitous in spherical tokamaks but rarely observed in conventional tokamaks. Turbulence is typically lower in spherical tokamaks such as NSTX due to favorable curvature and enhanced rotation shear [43], allowing $\alpha \sim \nu$. Similar findings exist on other devices. Calculations by Lilley, Breizman, and Sharapov found that $\alpha/\nu = 0.6$–5 at the toroidicity-induced Alfvén eigenmode (TAE) resonance for beam-heated MAST plasmas [14]. This had the practical consequence of explaining why those discharges often feature bursting TAEs, in contrast to radio-frequency-heated discharges with larger scattering rates. In addition, Lesur *et al.* analyzed a case of chirping in JT-60U by fitting the observed frequency sweeping to a theoretical model, enabling the extraction of the collisional coefficients, finding α/ν in the range of 0.2–0.7—consistent with measured plasma parameters [51]. Furthermore, simulations of an International Tokamak Physics Activity (ITPA) tokamak benchmark case and the W7-X stellarator using a global hybrid MHD-kinetic code performed by Slaby *et al.* concluded that including realistic amounts of drag in the Fokker-Planck collision operator affected both the AE saturation level and its long-term nonlinear behavior [52]. Lastly, negative triangularity experiments on DIII-D observed a greater tendency for chirping than in matched positive triangularity discharges [23]. Van Zeeland *et al.* attributed this difference to the lower level of turbulent scattering in negative triangularity, increasing α/ν and making nonsteady behavior more likely according to theory. The relative propensity for chirping in several distinct scenarios empirically supports the relevance of drag to wave-particle dynamics in fusion plasmas, especially tokamaks with low aspect ratio or negative triangularity.

Connection to galactic dynamics.—Beyond fusion plasmas, the methods presented in this Letter are directly applicable to the resonant gravitational interaction of the rotating galactic bar and heavy bodies in its orbit such as black holes, massive astrophysical compact halo objects, and stars in a tepid disk [53]. As recently demonstrated by Hamilton *et al.* [34], the collisional bump-on-tail problem in plasmas (studied in this Letter) can be made isomorphic to this application. In this formalism, the role of the AE is played by the rigidly rotating galactic bar, the resonant fast ions are replaced by the orbiting bodies, and the electrostatic potential is replaced by a gravitational one. Diffusive scattering occurs due to stochastic potential fluctuations, while heavy bodies experience drag when passing through a background of lighter masses. The resonant interaction exerts a torque on the bar and leaves a collisionally dependent imprint on the spatial distribution of surrounding bodies in the galaxy, which in tandem with observations can be used to constrain theoretical models of dark matter. Specifically, the collisions experienced by orbiting heavy compact objects are dominated by drag when their mass ratio to the background particles is sufficiently large [54,55]. Hence, the novel influence of drag on wave-particle interactions derived in this Letter is relevant for an accurate description of the resonant galactic bar–heavy body system in the presence of collisions.

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[1] T. H. Dupree, Phys. Fluids **9**, 1773 (1966).
[2] J. Weinstock, Phys. Fluids **11**, 1977 (1968).

[3] A. Salat, Phys. Fluids **31**, 1499 (1988).
[4] O. Ishihara, H. Xia, and A. Hirose, Phys. Fluids B **4**, 349 (1992).
[5] J. A. Krommes, J. Plasma Phys. **81**, 205810601 (2015).
[6] H. Berk, B. Breizman, J. Fitzpatrick, and H. Wong, Nucl. Fusion **35**, 1661 (1995).
[7] J. D. Callen, Phys. Plasmas **21**, 052106 (2014).
[8] N. H. Bian, E. P. Kontar, and H. Ratcliffe, J. Geophys. Res. **119**, 4239 (2014).
[9] V. N. Duarte, N. N. Gorelenkov, R. B. White, and H. L. Berk, Phys. Plasmas **26**, 120701 (2019).
[10] P. J. Catto, J. Plasma Phys. **86**, 815860302 (2020).
[11] P. J. Catto and E. A. Tolman, J. Plasma Phys. **87**, 905870606 (2021).
[12] C. H. Su and C. Oberman, Phys. Rev. Lett. **20**, 427 (1968).
[13] R. B. White, V. N. Duarte, N. N. Gorelenkov, and G. Meng, Phys. Plasmas **26**, 032508 (2019).
[14] M. K. Lilley, B. N. Breizman, and S. E. Sharapov, Phys. Rev. Lett. **102**, 195003 (2009).
[15] V. N. Duarte, H. L. Berk, N. N. Gorelenkov, W. W. Heidbrink, G. J. Kramer, R. Nazikian, D. C. Pace, M. Podestà, B. J. Tobias, and M. A. Van Zeeland, Nucl. Fusion **57**, 054001 (2017).
[16] J. D. Gaffey, J. Plasma Phys. **16**, 149 (1976).
[17] N. N. Gorelenkov, V. N. Duarte, M. Podestà, and H. L. Berk, Nucl. Fusion **58**, 082016 (2018).
[18] M. Podestà, M. Gorelenkova, and R. B. White, Plasma Phys. Controlled Fusion **56**, 055003 (2014).
[19] M. Gryaznevich and S. Sharapov, Nucl. Fusion **46**, S942 (2006).
[20] M. Podestà, R. Bell, A. Bortolon, N. Crocker, D. Darrow, A. Diallo, E. Fredrickson, G.-Y. Fu, N. Gorelenkov, W. Heidbrink, G. Kramer, S. Kubota, B. LeBlanc, S. Medley, and H. Yuh, Nucl. Fusion **52**, 094001 (2012).
[21] Y. V. Petrov, N. N. Bakharev, V. K. Gusev, V. B. Minaev, V. A. Kornev, G. S. Kurskiev, M. I. Patrov, N. V. Sakharov, S. Y. Tolstyakov, P. B. Shchegolev, J. Plasma Phys. **81**, 515810601 (2015).
[22] K. G. McClements and E. D. Fredrickson, Plasma Phys. Controlled Fusion **59**, 053001 (2017).
[23] M. Van Zeeland, C. Collins, W. Heidbrink, M. Austin, X. Du, V. Duarte, A. Hyatt, G. Kramer, N. Gorelenkov, B. Grierson, D. Lin, A. Marinoni, G. McKee, C. Muscatello, C. Petty, C. Sung, K. Thome, M. Walker, and Y. Zhu, Nucl. Fusion **59**, 086028 (2019).
[24] M. K. Lilley, B. N. Breizman, and S. E. Sharapov, Phys. Plasmas **17**, 092305 (2010).
[25] V. N. Duarte, H. L. Berk, N. N. Gorelenkov, W. W. Heidbrink, G. J. Kramer, R. Nazikian, D. C. Pace, M. Podestà, and M. A. Van Zeeland, Phys. Plasmas **24**, 122508 (2017).
[26] V. N. Duarte, N. N. Gorelenkov, M. Schneller, E. D. Fredrickson, M. Podestà, and H. L. Berk, Nucl. Fusion **58**, 082013 (2018).
[27] L. D. Landau, Sov. Phys. JETP **10**, 25 (1946).
[28] T. O'Neil, Phys. Fluids **8**, 2255 (1965).
[29] R. K. Mazitov, J. Appl. Mech. Tech. Phys. **6**, 22 (1965).
[30] D. Lynden-Bell, Mon. Not. R. Astron. Soc. **124**, 279 (1962).
[31] P. A. Sweet, Mon. Not. R. Astron. Soc. **125**, 285 (1962).
[32] A. A. Mikhailovskii and A. M. Fridman, Sov. Phys. JETP **34**, 243 (1972).
[33] D. D. Ryutov, Plasma Phys. Controlled Fusion **41**, A1 (1999).
[34] C. Hamilton, E. Tolman, L. Arzamasskiy, and V. Duarte, arXiv:2208.03855.
[35] H. L. Berk, B. N. Breizman, and M. Pekker, Plasma Phys. Rep. **23**, 778 (1997).
[36] H. L. Berk, B. N. Breizman, and H. Ye, Phys. Rev. Lett. **68**, 3563 (1992).
[37] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.130.105101, which includes Refs. [44–48].
[38] J. B. Lestz and V. N. Duarte, Phys. Plasmas **28**, 062102 (2021).
[39] H. L. Berk, B. N. Breizman, and M. Pekker, Phys. Rev. Lett. **76**, 1256 (1996).
[40] V. N. Duarte and N. N. Gorelenkov, Nucl. Fusion **59**, 044003 (2019).
[41] J. Lang and G.-Y. Fu, Phys. Plasmas **18**, 055902 (2011).
[42] W. Zhang, Z. Lin, and L. Chen, Phys. Rev. Lett. **101**, 095001 (2008).
[43] S. Kaye *et al.*, Nucl. Fusion **47**, 499 (2007).
[44] Y. Zhang, W. W. Heidbrink, H. Boehmer, R. McWilliams, G. Chen, B. N. Breizman, S. Vincena, T. Carter, D. Leneman, W. Gekelman, P. Pribyl, and B. Brugman, Phys. Plasmas, **15**, 012103 (2008).
[45] W. W. Heidbrink, H. Boehmer, R. McWilliams, A. Preiwisch, Y. Zhang, L. Zhao, S. Zhou, A. Bovet, A. Fasoli, I. Furno, K. Gustafson, P. Ricci, T. Carter, D. Leneman, S. K. P. Tripathi, and S. Vincena, Plasma Phys. Controlled Fusion **54**, 124007 (2012).
[46] W. Gekelman, P. Pribyl, Z. Lucky, M. Drandell, D. Leneman, J. Maggs, S. Vincena, B. Van Compernolle, S. K. P. Tripathi, G. Morales, T. A. Carter, Y. Wang, and T. DeHaas, Rev. Sci. Instrum. **87**, 025105 (2016).
[47] M. Sarfaty, S. De Souza-Machado, and F. Skiff, Phys. Rev. Lett. **80**, 3252 (1998).
[48] J. W. R. Schroeder, F. Skiff, G. G. Howes, C. A. Kletzing, T. A. Carter, and S. Dorfman, Phys. Plasmas **24**, 032902 (2017).
[49] T. J. Birmingham and M. Bornatici, Phys. Fluids **15**, 1785 (1972).
[50] The BOT code repository, https://github.com/mklilley/BOT.
[51] M. Lesur, Y. Idomura, K. Shinohara, and X. Garbet (The JT-60 Team), Phys. Plasmas **17**, 122311 (2010).
[52] C. Slaby, A. Koenies, R. Kleiber, and H. Leyh, Nucl. Fusion **59**, 046006 (2019).
[53] J. B. Fouvry, C. Pichon, J. Magorrian, and P. H. Chavanis, Astron. Astrophys. **584**, A129 (2015).
[54] J. Binney and S. Tremaine, *Galactic Dynamics* (Princeton University Press, Princeton, NJ, 2008).
[55] L. Ciotti, *Introduction to Stellar Dynamics* (Cambridge University Press, Cambridge, England, 2021).