The geometry of the non-neutral experiments is cylindrical with an axial magnetic field. The magnetic field is typically strong enough that the larmor radius is much smaller than any other scale length in the plasma and all relevant frequencies are small compared to the cyclotron frequency. Under these conditions the basic equations for an electron plasma ($q = -e$) are Poisson’s equation,

$$\nabla^2 \phi = 4 \pi e \int f \, dv,$$

(1)

the drift kinetic equation with a collision operator,

$$\frac{\partial f}{\partial t} + v \frac{\partial f}{\partial z} + \frac{e}{m} \frac{\partial f}{\partial \phi} + \frac{c}{B} \times \nabla \phi \cdot \nabla f = C(f),$$

(2)

and the boundary conditions on the conducting walls. Here $v$ is the axial velocity, $C(\cdot)$ is the approximate Fokker–Planck collision operator.\(^{15}\)
II. DERIVATION OF THE TRANSPORT EQUATIONS

We take as our model a cylindrical plasma of length $L$ with flat ends (see Fig. 1). The model thus ignores end effects and is most suitable for long, thin plasmas. This model allows us to replace the actual plasma by an infinitely long plasma with periodicity $2L$. It also allows us to linearize $f$ and $\phi$ as follows:

$$\phi(r, \theta, z, t) = \phi_0(r) + \phi_1(r, \theta, z, t)$$  \hspace{1cm} (5)

and

$$f(r, \theta, z, t) = f_0(r) + f_1(r, \theta, z, t).$$  \hspace{1cm} (6)

Returning these to Eq. (4) and keeping only zeroth order terms gives

$$\frac{\partial f_0}{\partial t} + v \frac{\partial f_0}{\partial z} = C(f_0),$$  \hspace{1cm} (7)

which has the well known solution

$$f_0(r) = \frac{n_0(r)}{\sqrt{2\pi} \bar{\theta}} \exp\left(-\frac{v^2}{2\bar{\theta}^2}\right).$$  \hspace{1cm} (8)

Here $\bar{\theta}$ may also be a function of radius.

A. First order

Keeping terms of first order in Eq. (4) gives

$$\frac{\partial f_1}{\partial t} + v \frac{\partial f_1}{\partial z} + e \frac{\partial \phi_1}{\partial z} \frac{\partial f_0}{\partial \bar{\theta}} + \frac{c}{B} \frac{\partial \phi_1}{\partial \bar{\theta}} \frac{\partial \phi_1}{\partial \bar{\theta}} - \frac{c}{B} \frac{\partial \phi_1}{\partial \bar{\theta}} \frac{\partial f_0}{\partial \bar{\theta}} = C(f_1).$$  \hspace{1cm} (9)

We now take advantage of the various periodicities in the model to write

$$\phi_1(r, \theta, z, t) = \sum_{n,l} \phi_{n,l,0}(r) \cdot \exp\left[ i \left( \frac{n\pi}{L} z + l\theta - \omega t \right) \right]$$  \hspace{1cm} (10)

and

$$f_1(r, \theta, z, t) = \sum_{n,l} f_{n,l,0}(r) \cdot \exp\left[ i \left( \frac{n\pi}{L} z + l\theta - \omega t \right) \right],$$  \hspace{1cm} (11)

where the sums are over both negative and positive values. Note especially that $\omega$ can be positive or negative. A positive $\omega$ corresponds to an asymmetry that rotates in the same direction as the plasma column; a negative $\omega$ asymmetry rotates against the column. The Fourier mode amplitudes are given by

$$\phi_{n,l,0}(r) = \int_{-L2}^{L} dz \int_{0}^{2\pi} d\theta \int_{0}^{r} dt \int_{0}^{t} \tau \times \exp\left[ i \left( \frac{n\pi}{L} z + l\theta - \omega \tau \right) \right] \phi_1(r, \theta, z, t)$$  \hspace{1cm} (12)

and similarly for $f_{n,l,0}$. Here $t$ is the duration of the experiment. Substituting in Eq. (9) and solving for $f_{n,l,0}$ we obtain

$$f_{n,l,0}(r) = \frac{c l}{rB} \frac{\partial f_0}{\partial \bar{\theta}} - \frac{n\pi}{L} \frac{\partial \phi_0}{\partial \bar{\theta}} \phi_{n,l,0}(r).$$  \hspace{1cm} (13)

Here we have noted that $(c/B)(1/l)(\partial \phi_0/\partial r)$ is the azimuthal $E \times B$ rotation frequency $\omega_R$ and have defined an effective collision frequency $\nu_{eff}$,

$$C(f_{n,l,0}) = -\nu_{eff} f_{n,l,0}.$$  \hspace{1cm} (14)

B. Second order

Since our interest is in radial transport we integrate Eq. (4) over $z$, $\theta$, and $v$. Defining

$$N(r, t) = \int_{-L2}^{L} dz \int_{0}^{2\pi} d\theta \int_{0}^{\infty} dv \cdot f(r, \theta, z, t)$$  \hspace{1cm} (15)

and noting $f(z=L) = f(z=-L)$ and $f(v=\pm \infty) = 0$ we obtain

$$\frac{\partial N}{\partial t} + c \frac{\partial N}{\partial r} \int_{-L2}^{L} dz \int_{0}^{2\pi} d\theta \int_{0}^{\infty} dv \frac{\partial \phi}{\partial r} \frac{\partial f}{\partial \bar{\theta}} - \frac{\partial \phi}{\partial \bar{\theta}} \frac{\partial f}{\partial \bar{\theta}} = 0.$$  \hspace{1cm} (16)

We note that the second term can be written as $(\partial \phi/\partial r)[f(\partial \phi/\partial \theta)] - f(\partial \phi/\partial \theta)(\partial \phi/\partial \theta)$ and after integrating by parts obtain

$$\frac{\partial N}{\partial t} + c \frac{1}{B} \frac{\partial N}{\partial r} \int_{-L2}^{L} dz \int_{0}^{2\pi} d\theta \int_{0}^{\infty} dv \frac{\partial \phi}{\partial \bar{\theta}} = 0.$$  \hspace{1cm} (17)

Substituting in from Eqs. (5) and (6) gives

$$\frac{\partial N}{\partial t} + c \frac{1}{B} \frac{\partial N}{\partial r} \int_{-L2}^{L} dz \int_{0}^{2\pi} d\theta \int_{0}^{\infty} dv \frac{\partial f_1}{\partial \bar{\theta}} = 0.$$  \hspace{1cm} (18)

From Eqs. (10) and (11),
\begin{align}
  f_1 \frac{\partial \phi_1}{\partial \theta} &= \left( \sum_{n,l,\omega} f_{n,l,\omega}(r) \cdot \exp \left( i \frac{n \pi}{L} z + l \theta - \omega t \right) \right) \\
  &\cdot \left( \sum_{n',l',\omega'} i l' \phi_{n',l',\omega'}(r) \right) \\
  &\exp \left( i \frac{n' \pi}{L} z + l' \theta - \omega' t \right) .
\end{align}

Eliminating \( f_{n,l,\omega} \) using Eq. (13), the right-hand side of (19) becomes

\begin{align}
  \sum_{n,l,\omega} \sum_{n',l',\omega'} i l' \phi_{n,l,\omega} \phi_{n',l',\omega'} &\cdot \frac{c l}{r B} \frac{\partial f_0}{\partial r} - \frac{n \pi e \partial f_0}{L m \partial v} \\
  &\cdot \exp \left( i \frac{n \pi}{L} z + l \theta - \omega t \right) \times \exp \left( i \frac{n' \pi}{L} z + l' \theta - \omega' t \right) .
\end{align}

Returning this expression to Eq. (18), we perform the \( z \)- and \( \theta \)-integrals and also integrate over the duration of the experiment to obtain

\begin{align}
  \frac{\partial N}{\partial t} &= - \frac{c}{B} \frac{1}{r} \frac{\partial}{\partial r} 2 \pi L \int_{-\infty}^{\infty} du \\
  &\times \sum_{n,l,\omega} \left| \phi_{n,l,\omega} \right|^2 i l n \pi e \partial f_0 \\
  &\cdot \exp \left( i \frac{n \pi}{L} z + l \theta - \omega t \right) .
\end{align}

where \( \langle \partial N/\partial t \rangle = (1/\pi) \int d\theta (\partial N/\partial \theta) \) and we have used the fact that \( \phi_{n,-l,-\omega} = \phi_{n,l,\omega}^* \). Multiplying top and bottom by the complex conjugate of the denominator and keeping the real part of the result (since the physical result must be real) gives

\begin{align}
  \frac{\partial N}{\partial t} &= 2 \pi L \frac{1}{r} \frac{\partial}{\partial r} \sum_{n,l,\omega} \left| c l \phi_{n,l,\omega} \right|^2 \\
  &\times \int_{-\infty}^{\infty} du \left( \frac{n \pi}{L} \frac{v_{\text{eff}}}{v + l \omega R - \omega} \right)^2 \\
  &\times \left( \frac{v_{\text{eff}}^2}{\omega R} \right) \\
  &\times \frac{\partial f_0}{\partial r} - \frac{n \pi e}{L m c l} \frac{\partial f_0}{\partial v} .
\end{align}

Noting that \( N \) is \( 2 \pi L \) times the plasma density and recalling the particle continuity equation, we can identify the average radial particle flux as

\begin{align}
  \Gamma &= - \sum_{n,l,\omega} \frac{c l \phi_{n,l,\omega}}{r B} \int_{-\infty}^{\infty} du \left( \frac{n \pi}{L} \frac{v_{\text{eff}}}{v + l \omega R - \omega} \right)^2 + v_{\text{eff}}^2 \\
  &\times \left( \frac{v_{\text{eff}}}{\omega R} \right) \\
  &\times \left( \frac{c l \phi_{n,l,\omega}}{r B} \right) .
\end{align}

Note that although the flux involves an integral over all velocities this is strongly conditioned by the resonance function \( v_{\text{eff}} \left( [(n \pi L) v + l \omega R - \omega)^2 + v_{\text{eff}}^2 \right)^{-1} \). This function, which peaks at the velocity

\begin{align}
  v_{\text{res}} &= \frac{L}{n \pi} (\omega - l \omega_R) \\
  \text{and has full-width at half-maximum } \Delta v &= \frac{2 L v_{\text{eff}}}{n \pi},
\end{align}

shows that the flux tends to be dominated by particles that move resonantly with the asymmetry mode specified by \( n, l, \) and \( \omega \). If \( \Delta v \ll \bar{v} \) (i.e., if the width of the resonance is small compared to variations in \( f_0 \)), then

\begin{align}
  \frac{\partial f_0}{\partial r} &= \frac{L}{n \pi} \left( v_{\text{eff}} \left( \frac{n \pi v}{L + l \omega R - \omega} \right)^2 + v_{\text{eff}}^2 \right) \\
  \Rightarrow \Gamma &= - \sum_{n,l,\omega} \frac{L}{n \pi} \left( \frac{n \pi v}{L + l \omega R - \omega} \right)^2 \\
  &\times \left( \frac{n \pi v}{L + l \omega R - \omega} \right)^2 + v_{\text{eff}}^2 \right) \\
  &\times \left( \frac{c l \phi_{n,l,\omega}}{r B} \right) .
\end{align}

The velocity integral in Eq. (23) is now easily done. We obtain

\begin{align}
  \Gamma &= - \sum_{n,l,\omega} \frac{n_0}{\sqrt{2 \pi \bar{v}^2}} \frac{L}{n \pi} \left( \frac{n \pi v}{L + l \omega R - \omega} \right)^2 \\
  &\times \left( \frac{n \pi v}{L + l \omega R - \omega} \right)^2 + v_{\text{eff}}^2 \right) \\
  &\times \left( \frac{n \pi v}{L + l \omega R - \omega} \right)^2 + v_{\text{eff}}^2 \right) \\
  &\times \left( \frac{c l \phi_{n,l,\omega}}{r B} \right) .
\end{align}

If we plug in \( f_0 \) of the form of Eq. (8) this becomes

\begin{align}
  \Gamma &= - \sum_{n,l,\omega} \frac{n_0}{\sqrt{2 \pi \bar{v}^2}} \frac{L}{n \pi} \left( \frac{n \pi v}{L + l \omega R - \omega} \right)^2 \\
  &\times \left( \frac{n \pi v}{L + l \omega R - \omega} \right)^2 + v_{\text{eff}}^2 \right) \\
  &\times \left( \frac{n \pi v}{L + l \omega R - \omega} \right)^2 + v_{\text{eff}}^2 \right) \\
  &\times \left( \frac{c l \phi_{n,l,\omega}}{r B} \right) .
\end{align}

where \( x = v/v_{\text{res}} \) and \( \omega_c \) is the cyclotron frequency \( eB/mc \).

It is worth noting several features of this solution. As is typical of plateau regime transport, the flux is independent of collision frequency and proportional to the square of the asymmetry amplitude. The plasma length \( L \) appears explicitly, but is also part of the variable \( x \). Also hidden in this variable is the asymmetry frequency \( \omega \), and we note that \( x \) can be positive or negative as \( \omega \) is greater than or less than \( \omega_R \). Thus, while static field asymmetries \((\omega = 0, x < 0)\) move electrons radially outward, an appropriately chosen asymmetry \((\omega > \omega_R, x > 0)\) can move particles radially inward.

Equation (27) can be heuristically derived to within a numerical factor. Start from the relation

\begin{align}
  \Gamma &= - D \frac{\partial f_0}{\partial H} \Delta v ,
\end{align}

where, for the case of a static asymmetry \((\omega = 0)\), the derivative is taken at constant \( H \) since the Hamiltonian is constant for a particle moving in a static field, the distribution function and diffusion coefficient are evaluated at the resonant velocity, and \( \Delta v \) is the width of the velocity resonance. Since the Hamiltonian for this plasma is \( H = (p_z^2/2m) - e \phi_0 + \mu B \), the Maxwellian distribution function can be written

\begin{align}
  f_0 &= \frac{n_0}{\sqrt{2 \pi T/m}} \exp \left( - \frac{H + e \phi_0 - \mu B}{T} \right) .
\end{align}
Noting that \( n_0, T, \) and \( \phi_0 \) are functions of \( r \), it is straightforward to show that \( (1/\bar{f}_0)(\partial \phi_0/\partial r)_H \) reproduces the curly bracket of Eq. (27) for the case \( \omega_l = 0 \). In order to include the \( \omega \neq 0 \) case, \( (\partial \phi_0/\partial r)_H \) must be generalized to \( (\partial \phi_0/\partial r)_H \), where \( \bar{H} = \bar{H} - (\omega l_l)P_g \) is the Hamiltonian in a frame rotating at frequency \( \omega_0 \) and \( P_g \) is the single particle canonical angular momentum which in the guiding center approximation is equal to \( -\frac{1}{2}m\omega_r \). The diffusion coefficient \( D \) is estimated as the average step size squared divided by the time between collisions: \( D = (\Delta r)^2/\tau \). The average step size \( \Delta r \) is the radial \( \times B \) drift velocity \( v_\perp = (eE_0/B) \) times the time between collisions. Since the relevant collisions occur at the enhanced rate \( v_{\text{eff}} \) we obtain \( D = (1/v_{\text{eff}})(cl\phi_{\text{sub}}/lB)^2 \). Finally the width of the velocity resonance \( \Delta v \) in the plateau regime can be obtained by taking the half-width of the velocity resonance function appearing in Eq. (23). \( \Delta v \) \( \approx \) \( v_{\text{eff}}/\sqrt{n\pi} \). Plugging these estimates into Eq. (28), the factor of \( v_{\text{eff}} \) cancels and the resulting flux is equal to the left-hand side of Eq. (27) divided by \( \pi \).

As stated in the Introduction, the plateau regime corresponds to a collisionality regime where \( v_{\text{eff}} \approx \omega_T \). For our case where the asymmetry varies in both \( z \) and \( \theta \), the trapping frequency is given by

\[
\omega_T^2 = \frac{e}{m} \left( \frac{n \pi}{L} \right)^2 - \frac{c^2 l^2}{r B} \frac{d \omega_e}{d r} \phi_{n,l,w}.
\]

An estimate for \( v_{\text{eff}} \) can be obtained by examining the form of Eqs. (3), (13), and (14). Because of the presence of a velocity resonance, the first term in Eq. (3) will dominate and we can estimate \( v_{\text{eff}} \approx v_{\text{ce}}(v/\Delta v)^2 \). But the velocity resonance has half-width \( \Delta v = v_{\text{eff}}/\sqrt{n\pi} \). Combining these we obtain

\[
v_{\text{eff}} = v_{\text{ce}} \left( \frac{n \pi \bar{v}}{L} \right)^2 \approx v_{\text{ce}} \omega_b^2.
\]

where \( \omega_b \) is the axial bounce frequency \( \pi \bar{v}/L \). Since \( \omega_b \) is large compared to \( \nu_{\pi} \), \( v_{\text{eff}} \) will be larger than \( v_{\text{ce}} \). This is a reflection of the fact that only a small change in the velocity is necessary to knock a particle out of resonance. The collision time for this type of event is much less than for a ninety degree collision.

Similar heuristic arguments can be employed to obtain an approximate expression for \( \Gamma \) in the banana regime where \( \omega_T \approx v_{\text{eff}} \). The basic radial step is now the width of the resonance island which may be estimated as \( \Delta r \approx (v_{\perp}/\omega_T) \) \( = (c l \phi_{n,l,w}/l B \omega_T) \) and thus \( D = v_{\text{eff}}(c l \phi_{n,l,w}/l B \omega_T)^2 \). The width of the velocity resonance, which is broadened in the banana regime, is given by \( \Delta v \approx (L/n \pi)\omega_T \). In this case the collision frequency does not cancel out and we obtain

\[
\Gamma = -\sum n_{l,w} \nu_{\text{ce}} \left( \frac{L}{n \pi} \right)^2 \left( \frac{r \bar{v}}{c} \right)^2 \frac{e \phi_{n,l,w}}{T} \frac{1}{\left( 1 - \frac{L}{n \pi} \right)^2} \frac{d \omega_e}{d r} \frac{1}{\left( 1 - \frac{L}{n \pi} \right)^2} \frac{d \omega_e}{d r} \left( \frac{1}{2} \right) \frac{n_0}{\sqrt{2 \pi}} \frac{1}{v_{\text{eff}} d r} + \frac{1}{4 T} \frac{d}{d r} \left( x^2 - \frac{1}{2} \right) + \sqrt{2} \frac{n \pi r \omega_e}{L} \left( \frac{1}{L} \right) e^{-x^2}.
\]

III. DETERMINATION OF THE ASYMMETRIC POTENTIAL IN THE PLASMA

In order to evaluate the flux we must determine the complex Fourier mode amplitudes \( \phi_{n,l,w}(r) \) produced in the plasma by the applied wall potentials. For perturbed potentials of the form of Eq. (10), Poisson’s Eq. (1) becomes [using Eq. (13) to eliminate \( f_{n,l,w} \)]

\[
\left[ \frac{1}{r} \frac{d}{d r} \frac{d}{d r} - \frac{l^2}{r} - \left( \frac{n \pi}{L} \right)^2 \phi_{n,l,w}(r) \right]
\]

\[
= 4 \pi \epsilon \int dv \frac{c l}{r B} \frac{d \phi_0}{d r} - \frac{v_{\text{eff}}}{L} \frac{d \phi_0}{d r} \phi_{n,l,w}(r).
\]

This equation must be solved subject to the conditions that \( \phi_{n,l,w} \) is finite at \( r = 0 \) and equal to the wall potential at \( r = R \). Although analytical solutions exist18 for special cases (e.g., constant density and temperature), a numerical solution is required for experimental density and temperature profiles. Fortunately, this is quickly and easily done using a modification of the “shooting” technique.

For Maxwellian \( f_0 \), the right hand side of Eq. (32) can be cast in terms of the plasma dispersion function19 \( Z(x) \) and its derivative \( Z'(x) \) for which numerical codes exist. This relieves us of the task of numerically evaluating the integral. The result is18

\[
\left[ \frac{d^2}{d r^2} + \frac{1}{r} \frac{d}{d r} - \frac{l^2}{r} - \left( \frac{n \pi}{L} \right)^2 \right] \phi_{n,l,w} = \phi_{n,l,w}
\]

\[
= \left( \frac{L}{n \pi r \omega_e} \right) Z(x) \frac{d}{d r} \frac{\omega_p^2}{a} - \left( \frac{\omega_p}{a} \right)^2 \times Z'(x) \left[ 1 + \frac{L x}{n \pi r \omega_e} \frac{d a}{d r} \right] \phi_{n,l,w}.
\]

Here \( \omega_p \) is the plasma frequency and \( a = \sqrt{2} \bar{v} \). The radial derivatives on the left hand side are now written as second order central finite difference expressions.20 Equation (33) then becomes

\[
\phi_{j+1} - 2 \phi_j + \phi_{j-1} \left( \frac{\Delta r}{r_{j}} \right) + \frac{1}{r_j} \phi_{j+1} - \phi_{j-1} - \eta_j \phi_j = 0.
\]

Here we have suppressed the \( n, l, \omega \) indices and have divided the space between \( r = 0 \) and \( R \) into intervals of length \( \Delta r \) specified by the index \( j \). The quantity \( \eta_j \) is defined as
standing waves and the strong dip near $\omega_R$ corresponds to Debye shielding. We emphasize, however, that there is continuous variation between these extreme cases and that an accurate determination of the transport flux depends sensitively on this calculation. Note also that, for a given asymmetry frequency, the field might be enhanced or diminished by these collective effects, depending on the details of the plasma parameters.

This variation in the amplitude of $E_\phi$ is accompanied by strong variations in the radial dependence of $\phi_{n,l,w}$. A sampling of this variation is shown in Fig. 3, where we plot the normalized magnitude of $\phi_{n,l,w}$ vs $r$ for the four frequencies indicated in Fig. 2(a) on the $T=1$ eV curve. Again the extreme cases of standing waves (3A) and shielding (3C) are shown along with two intermediate cases (3B, 3D).

### IV. DISCUSSION

It is interesting to compare these results with previous and ongoing experimental work. The presence of asymmetry-induced transport in non-neutral plasmas was first suggested by the discovery\(^4\) of confinement time scaling with $(L/B)^{-2}$. Comparing with our results, it is tempting to seize upon the $(L/L_{ce})^2 \sim L^2/B^2$ in the leading factor of the banana regime flux given in Eq. (31) and contrast this with the $L/B^2$ for the plateau regime [Eq. (27)]. However, $L$ is also hidden in the variable $x$, and the third term in brackets also contains $\omega_{ce}$. Thus, without a knowledge of the spectrum of background asymmetries it is impossible to draw a firm conclusion.

Consistent with this theory, early experiments\(^{6,21}\) found that standing waves could produce enhanced transport. These experiments also reported that modes rotating in the same direction as the plasma column but at a faster rate (i.e., $\omega > \omega_R$) produced inward transport. More recently, Huang and co-workers\(^8\) and Anderegg and co-workers\(^{22}\) have used an asymmetry with $\omega > \omega_R$ to balance the normal background transport and produce a steady-state plasma. The importance of standing waves in enhancing transport is also clear in this latter paper.
In order to study asymmetry-induced transport apart from such collective enhancements, Eggleston\textsuperscript{23} has measured the confinement of low density ($<10^5$ cm$^{-3}$) electrons in a trap where a biased wire running along the axis of the trap replaces the plasma column. Under these conditions the variations in $E_\theta$ shown in Fig. 2 are essentially eliminated. The confinement time in this trap was found to have the same magnitude and $(L/B)^{-2}$ scaling as observed in the higher density experiments. Since the low density and higher temperature of this experiment give an electron–electron collision frequency $\nu_{ee}$ that is much lower than in the plasma experiments, it was argued that the transport could be in this regime due to the small value of $\nu_{ee}$. However, this conundrum is somewhat softened by the fact that the transport regime depends on $\nu_{eff}$ rather than $\nu_{ee}$, and the former has a much weaker dependence on density and temperature. Referring to Eq. (30), we see that although $\nu_{ee}$ goes like $n_0^1/T^{1/2}$, $\nu_{eff}$ goes like $n_0^{1/3}/T^{1/6}$. [Note that, in Eq. (30), $n$ is the axial mode number, not the density $n_0$.] Thus the factor of $10^3$ difference in $\nu_{ee}$ becomes a factor of 6 in $\nu_{eff}$.

Several experiments have measured the amplitude-scaling of asymmetry-induced transport. In experiments with static (i.e., $\omega=0$) asymmetries, Notte and Fajans\textsuperscript{7} observed a confinement time scaling $\tau \propto \phi^4$, with $s = 1.72 - 2.14$, i.e., an amplitude scaling similar to the $\phi^2$ dependence of Eq. (27). Interestingly, this experimental scaling was found to hold for wall voltages up to 40 V, well past the point where the plateau regime theory should apply. In contrast, more recent work with static asymmetries has found a robust linear scaling\textsuperscript{24} while rotating-wall experiments\textsuperscript{22} have given $s = 0.7 - 1.1$. In low density experiments with time-varying asymmetries,\textsuperscript{26} $\phi^4$ scaling was observed for small amplitude asymmetries but at higher amplitudes the scaling was $\phi^{4/3}$. To our knowledge, no one has observed the banana-regime scaling of $\phi^{1/2}$.

Finally, Eggleston\textsuperscript{26} has measured the flux produced by a single asymmetry mode (i.e., a single value of $n$ and $l$) as a function of asymmetry frequency $\omega$. The experiments show a resonance similar to that predicted by our theory and thus seem to confirm that the transport is dominated by resonant particles, but there are also important differences between the experiment and this theory. These results, along with a numerical comparison of the experimental and theoretical flux, will be presented in a subsequent paper.

From this discussion it should be clear that asymmetry-induced transport is far from understood. Several experimenters are currently studying this transport, but the results are not yet in agreement with each other or with any theory. While the theory presented in this paper can certainly stand further refinement (e.g., a more realistic treatment of particle motion at the ends of the plasma), we hope it will contribute to discussions of this phenomena.

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