the open drift orbit region for their model storm period of August 1990. Thus one must indeed be careful in applying the standard model to ring current ions. The dotted line in Figure 1 encircles the region for which we have data and over which we apply the diffusion model. Above L=4.2, we include open and distorted orbits in the model only because the diffusion rate is greater than the loss rate so that the effect of the faster convective transport does not appear to change the solution significantly.

We can estimate the rates for direct comparison. The loss rate includes the contributions from both Coulomb drag and charge exchange processes which can be estimated crudely as follows:

$$\frac{1}{t_L} = \sum_i n_i \sigma_i v \tag{6}$$

where n_i is defined to be the appropriate target plasma density for that loss term (e.g., cold electrons for Coulomb drag), while σ_i is the appropriate cross section (tabulated for both charge exchange and Coulomb drag in paper 1) and v the velocity. A coarse approximation to the radial diffusion timescale is given by [e.g., *Cornwall*, 1968]

$$\frac{1}{t_D} = \frac{2D_{LL}}{(\Delta L)^2} \tag{7}$$

Thus we can directly compare the local diffusion rate with the loss rate, as plotted in figure 10 of paper 1. For L>4 and M>0.01 keV/nT the standard diffusive transport is generally substantially faster than the loss rate, justifying the inclusion of this region in the model.

Modified Resonant Frequencies

The standard diffusion mechanism for trapped particles may be described as a resonance between the perturbation frequency of the field (electric or magnetic) and the drift frequency of the ion [e.g., *Parker*, 1960; *Schulz and Lanzerotti*, 1974]. For example, an abrupt increase in the cross tail electric field would cause ions on the nightside to convect toward the Earth more quickly, but ions on the dayside would convect away from the Earth. If a corotating ion was near dusk at the start of the compression, but had drifted to dawn by the end of the compression, it would be found on an L shell closer to the Earth. Thus fluctuations at the same frequency as the drift frequency will cause the ion at a given energy to spiral either inward or outward, depending on the the ion drift phase in its azimuthal motion. So indeed, the radial direction of this spiral convection is drift phase dependent, but if the fluctuations occur randomly in time, we can average over phase and consider the radial transport to be diffusive. One might possibly surmise that such impulses would merely create structure in the convected signatures rather than the smooth profiles generated by diffusion. However, the subsequent radially dependent drift of the ion population smears out any structure that might have been created by any single impulse. Thus azimuthal drift, though not a stochastic process in itself, amplifies the stochasticity of the external impulses when they resonate with the drift frequency [e.g., *Chen et al.*, 1993] to cause diffusion. Should convection

Stall speed. For ions in the magnetosphere the corota tional electric field drift and the $\vec{\nabla}B$ drift have opposite sig so that there is an energy and location at which the drift spe is negligible. (Although it should be noted that in realist electric and magnetic field models the two drifts are not ge erally antiparallel.) Thus there should be no resonance, a very little standard resonant diffusion at that energy whe the drift velocity (frequency) goes to zero [*Riley and Wo* 1992] (in cycles per hour)

$$0 = \Omega_T = \frac{1}{24}\left(1 - 0.997\frac{E(L)}{E_0}\frac{L}{3}\right)$$

where E_0 is 10 keV/e, or

$$E(L)_{\text{stall}} \cong \frac{30}{L} \quad \text{keV}/e$$

Fälthammer [1965] modeled transport for particles wi high enough energy (E>100 keV) to permit the neglect any static electric fields. These coefficients are not valid f ring current energies at tens of keV and below, particular for energies close to the ion "stall" energy. Thus we inclu the effect of a static electric field in what follows [e. *Mozer*, 1971; *Riley and Wolf*, 1992].

Modified magnetic diffusion. We begin by includi only the effect of the corotation electric field on the i drift resonance. The modification to *Fälthammer*'s [196 linearized equations of motion (his equation 2.1.6) is slig replacing Ω in the $\vec{\nabla}B$ drift frequency with $\Omega_T = \Omega_{\nabla B}$ $\Omega_{E\times B}$ the combined $\vec{\nabla}B$ and $\vec{E} \times \vec{B}$ drift frequencies. T derivation of the magnetic diffusion coefficient remains t same resulting in (his equation 4.11),

$$\langle(\Delta r)^2\rangle = \frac{1}{8}\frac{r^4\Omega_T^2(\mu, r)}{B_d^2(r)} H_1(\Omega_T(\mu, r))$$

where B_d is the dipole field strength at the equator, r is t radial distance in R_E at the equator (equivalent to L), and is the power in the magnetic field fluctuations at a frequen Ω. In keeping with the subsequent literature, we simpli the leading coefficient from his 25/196 to 1/8. The mo important implication of this modified diffusion, is that can become arbitrarily small for a subset of ions in the ener range of our data set.

Since Ω_T can go to zero in our energy range, we mu model the power spectrum down to zero frequency. The retically, we reasonably expect the power to vanish both zero frequency and at infinite frequency. We have expe imental measurements of the magnetic field power spec made with mid- to low-latitude ground based magnetome or induction loops [*Lanzerotti and Wolfe*, 1980; *Maclenn et al.*, 1991]. All show a maximum power at or below a fr quency of 2 cycles/h. If we let Ω_M be the frequency of t maximum power, we can model the power spectrum wi parameterized exponents above and below Ω_M, as follow

$$H_1(\Omega_T) \propto (\Omega_T/\Omega_M)^{-l} \qquad \Omega_T \geq \Omega_M$$
$$\propto (\Omega_T/\Omega_M)^{k} \qquad \Omega_T < \Omega_M$$