Chapter 4

Modes inherent to nonadiabatic/kinetic passing electrons

4.1 Introduction

Having elucidated the effect of kinetic/nonadiabatic electrons on temperature gradient driven modes of thermal ions and trapped electrons we now discuss the modes inherent to the nonadiabatic/kinetic passing electrons. A correct kinetic/nonadiabatic passing electron model should be able to produce modes that are inherent to the nonadiabaticity of passing electrons. Therefore, this chapter is dedicated to the study of temperature and density gradient driven modes of kinetic electrons, namely, the electron temperature gradient (ETG) mode and universal drift mode. We will study the linear ETG mode in the presence of completely kinetic ions and Debye Shielding effect. A comparison of the pure ion temperature gradient (ITG) mode with the pure ETG mode in the presence of fully gyro-kinetic second species (e.g., electrons for ITG and ions for ETG) will be carried out. One will see the breaking of isomorphy of ITG and ETG modes even in the electrostatic limit without incorporating the trapped electrons. A comparison of the electron flux by ETG mode in the presence of nonadiabatic ions with \( \eta_i \) above the ITG threshold and ion flux by ITG mode in the presence of nonadiabatic electrons with \( \eta_e \) above the ETG threshold reveals that these modes are not independent of each other. In
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fact, one mode tends to reduce the transport by the other and vice versa.

Regarding the other mode driven by the density gradient of nonadiabatic passing electron, i.e., the universal drift mode: although the studies have evolved from a simple slab model to the toroidal geometry, most of them, however, are based on the fluid or hybrid kinetic-fluid models. A few gyrokinetic models either exploited the ballooning formulation or a simple geometry. In this chapter we will present a global, fully gyrokinetic linear study of the toroidicity driven universal drift mode considering both ions and electrons to be nonadiabatic. The formulation retains toroidal coupling effects due to both electron and ion $\nabla B$ and curvature drift with no assumption regarding the magnetic drift frequency of the ions and electrons compared to the mode frequency, thereby allowing full magnetic drift resonance by both species. The formulation also keeps the Landau damping term of both electrons and ions. The finite Larmor radius effects are kept up to all orders. More importantly, the present study retains the transit frequency resonance term in the nonadiabatic part of the density perturbation for both species as shown in Eqs. (2.8) and (2.9). It is to be noted that, we use a large aspect ratio, circular geometry for the tokamak, with no Shafranov shift. Though the universal toroidal mode is inherently due to the passing nonadiabatic electrons, effects of trapped electrons and trapped ions are also retained in the formulation. Furthermore, no collisional effect is considered in the formulation.

With this model, various parametric studies of the toroidal branch of universal mode have been carried out. We observe finite mode frequencies and growth rates beyond the critical $\eta = L_n/L_T$ for Ion Temperature Gradient (ITG) and Electron Temperature Gradient (ETG) modes, where $L_n$ and $L_T$ are, respectively, the density and temperature scale lengths. A comparative study of the contribution of magnetic drift resonance as well as Landau resonance from both species toward the stability properties of the mode is performed by a systematic parametric scan. An electromagnetic study of the mode is also carried out that elucidates the effect of finite $\beta$ on the universal drift mode driven by toroidicity. The effect of trapped electrons on the universal mode is studied and growth rates and real frequencies are compared with the ion temperature gradient mode and trapped electron mode.
4.2 Electron Temperature Gradient Driven Mode

Having unravelled the effect of the nonadiabatic passing electrons on the ITG mode [132] and trapped electron coupled ITG mode (ITG-TEM) [133] in the previous chapters, we now proceed to study the mode inherent to the nonadiabatic passing electrons, that is, the ETG mode.

With the formulation discussed in Chapters 2 and 3, the Poisson equation can be written as,

\[ \nabla^2 \tilde{\varphi} = \frac{e}{\varepsilon_0} \sum_{j=e,i} \tilde{n}_j(\mathbf{r}; \omega); \quad (4.1) \]

where, \( j \) stands for charge species, viz., electrons (e) and ions (i). For single charged passing ions and electrons we have

\[ \sum_{k'} \sum_{j=i} \mathcal{M}^j_{k,k'} \tilde{\varphi}_{k'} + \sum_{k} \sum_{j=e} \mathcal{M}^j_{k,k} \tilde{\varphi}_{k} = 0 \]

### 4.2.1 Profiles and Parameters

For the purpose of our study, we consider following profiles and parameters.

**Table 4.1: Profiles and parameters**

<table>
<thead>
<tr>
<th>Parameters</th>
<th>Equilibrium Profiles</th>
</tr>
</thead>
<tbody>
<tr>
<td>- B-field : ( B_0 = 1.0 ) Tesla</td>
<td>- N-profile and T-profile</td>
</tr>
<tr>
<td>- Temperature : ( T_0 = T(s_0) = 7.5 ) keV</td>
<td>( \frac{N(s)}{N_0} = \exp\left( -\frac{a \delta s_n}{L_{n0}} \tanh \left( \frac{s-s_0}{\delta s_n} \right) \right) )</td>
</tr>
<tr>
<td>- Major Radius : ( R = 2.0 ) m</td>
<td>( \frac{T_e(s)}{T_0} = \exp\left( -\frac{a \delta s_T}{L_{T0}} \tanh \left( \frac{s-s_0}{\delta s_T} \right) \right) )</td>
</tr>
<tr>
<td>- Minor Radius : ( a = 0.5 ) m</td>
<td>( \delta s_n = 0.35, \delta s_T = 0.2 ) at ( s = s_0 )</td>
</tr>
<tr>
<td>- radius : ( s = \rho/a, \ 0.01 &lt; s &lt; 1.0, \ s_0 = 0.6 )</td>
<td>- ( q(s) = 1.25 + 0.67 \ s^2 + 2.38 \ s^3 )</td>
</tr>
<tr>
<td>- ( \rho_{n0} = 0.4 \ m, \ L_{T0} = 0.2 \ m )</td>
<td>- (-0.06 \ s^4 ) such that ( q(s = s_0) = 2.0 );</td>
</tr>
<tr>
<td>- ( \eta_{i,e}(s_0) = 2.0, \ \epsilon_n = L_{n0}/R = 0.2 )</td>
<td>shear ( s = s_0, \ \bar{s} = 1. )</td>
</tr>
<tr>
<td>( \tau(s) = T_e(s)/T_i(s) = 1. )</td>
<td></td>
</tr>
</tbody>
</table>

The chosen parameters lead to the value of \( \rho^* \equiv \rho_{Li}(s = s_0)/a \simeq 0.0175 \).
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that $\rho_e = 2.065 \times 10^{-4} \text{m}$ and $\rho_i = 8.848 \times 10^{-3} \text{m}$. Any change in the parameters will be stated wherever necessary.

4.2.2 Pure ETG Mode

At the outset, let us compare the ETG mode results with respect to some known works. To that end we have chosen the local, linear, electrostatic, and kinetic results of Horton et al. [21] for which $k_0 \rho_e \simeq 0.5$, $\epsilon_n = L_n / R = 0.2$ and $\tau = 1.0$. From Figs.3 and 4 of Horton et al. [21] we have extracted manually some points and replotted with the results of our global, linear, electrostatic, and gyrokinetic model for nonadiabatic electrons in Fig. 4.1 for the same parameters. For the ease of comparison for the readers, we have also copied and pasted Figs. 3 and 4 of Horton et al. in Fig. 4.2. It is to be noted that $r_n$ and $v_{ei}$ in the latter case correspond to $L_n$ and $v_{the}$ in our case. Note that the differences in the real frequencies as well as growth rates in both cases can be ascribed to the differences in the two models, namely, local versus global and kinetic versus gyrokinetic.

Effect of Debye Shielding

Now a days, fast wave electron heating (FWEH) experiments [69], to study specifically electron channel transport, use preferentially dominant electron heating such that $T_e$ can take very high values. Furthermore, experiments dedicated to the ETG mode study require separation between the electron and ion channel of transport. This is achieved by reducing the energy exchange between the two species. The conducive environment is achieved with plasma of low density that ensures less amount of collisionality. Thus, the Debye length, which is proportional to $\sqrt{T_e/n}$, can be expected to violate the condition $k \lambda_{de} << 1$. In such a situation, one requires to take in to account the space charge effect, and the Debye shielding effect inevitably comes into the picture.

The Debye shielding effect was taken into account in a number of previous works in the slab and sheared slab [20, 134, 135] geometry and toroidal geometry [22, 136]. In the present case, we produce a toroidal mode number $n$ and $k_0 \rho_e$ scan with and without the Debye shielding effect for the ETG mode that contains no trapped particles and no effect of nonadiabatic ions. Figure. 4.3 presents the real frequency and growth rate versus $k_0 \rho_e$ with and without the Debye shielding.
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Figure 4.1: The normalized real frequency $\omega_r$ (upper panel) and growth rate $\gamma$ (lower panel) for the ETG mode as function of $\eta_e$ for $k_0\rho_e \simeq 0.5$, $\epsilon_n = L_n/R = 0.2$, and $\tau = 1.0$. The lines with squares represent manually extracted points from Horton et al. [21] which uses local kinetic formulation. The lines with open circles depict the results from our global linear gyrokinetic model.
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The corresponding toroidal mode numbers $n$ are shown in the upper axis for both frequency and growth rate of the mode. From the figure it is apparent that the mode frequency for the case without Debye shielding is being slightly higher than that in the case with the Debye shielding. Also, one can infer that the Debye shielding has strong stabilizing effect on the mode. One important point to be noted is that the Debye shielding effect removes the high $k$ tail of the ETG mode. The observed effect of the Debye shielding on the real frequency of the mode is weak as compared to the effect of the same on the growth rate. The other purpose, these figures serve, is that they exhibit the dispersion diagram for the ETG mode with and without Debye shielding showing the dependence of the frequency and growth rate on $k\theta\rho_e$. For the case with Debye shielding the growth rate peaks at $k\theta\rho_e = 0.34$ and for the case without Debye shielding the same peaks at $k\theta\rho_e = 0.5$, and both decrease by substantial fraction as one goes both side from the respective peaks. The toroidal mode numbers corresponding to both cases are $n = 250$ and $n = 360$, respectively. Thus, the Debye shielding not only reduces the growth rate of the mode but also shifts the maximum of the growth rate toward lower $k$, and wipes out the higher $k$ tail. It is interesting to see whether Debye shielding has any effect on the threshold of the onset of the ETG mode; it may be important to take into account the Debye shielding effect while calculating precisely the threshold. Figure 4.4 displays the growth rate of the ETG mode against $\eta_e$ for the two cases with and without the Debye shielding effect. The
Figure 4.3: The normalized real frequency $\omega_r$ (upper panel) and growth rate $\gamma$ (lower panel) for the ETG with and without Debye shielding for $\eta_e(s_0) = 2$. Ions are considered adiabatic.
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![Graph](image)

Figure 4.4: $\eta_e(s_0)$ scan for the growth rate of the ETG mode with and without Debye shielding for $n = 250$ and $380$, respectively.

Debye shielding enhances the threshold of the mode and the mode without the Debye shielding effect is more susceptible to the instability, lowering the threshold of the instability. Thus, one can conclude that Debye shielding affects significantly the ETG mode; it has stabilizing effect on the mode leading to enhanced threshold; shifts the maximum growth rate toward lower $k_0\rho_e$ or $n$; wipes out the high wave number tail of the mode.

**Effect of Nonadiabatic Ions: A Comparison With ITG Mode**

The coupling of ITG mode with trapped electrons has been studied both theoretically, and also observed experimentally. Earlier studies of the ETG mode preferred to proceed with the adiabatic ion model, upon the basis of the assumption that the ions are adiabatic as their Larmor radii are larger than the scale-length of the ETG/streamer/zonal flow/geodesic acoustic mode.

The finite mass effect of ions, nonetheless, was studied earlier in Refs. [136, 137] and was found to have very weak effect on the mode. However, of late, nonlinear simulations [36, 37, 91, 92, 93, 94] have demonstrated the limitations of this adiabatic ion model. The nonadiabatic ions are found to be important for the development of low $k$ tail of the ETG mode. Furthermore, the nonadiabatic ions help getting saturation in the simulation of the electron heat diffusivity. Here, we investigate the impact of the ions on the ETG mode if considered fully nonadia-
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Figure 4.5: Upper panel: The normalized real frequency $\omega_r$ for the ETG mode without Debye shielding $\eta_i(s_0) = 2.0$ with adiabatic ions and nonadiabatic ions. $\eta_i(s_0)$ takes values 2, 4, 6, 8 for nonadiabatic ions. Lower panel: The corresponding growth rates $\gamma$. 
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Figure 4.6: Upper panel: the normalized frequency $\omega_r$ for the ETG mode with Debye shielding for $\eta_e(s_0) = 2.0$ with adiabatic ions and nonadiabatic ions. $\eta_e(s_0)$ takes values 2, 4, 6, 8 for nonadiabatic ions. Lower panel: the corresponding growth rates $\gamma$. 

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Figure 4.7: Upper panel: the normalized frequency $\omega_r$ for the ITG mode for $\eta_i(s_0) = 2.0$ with adiabatic and nonadiabatic electrons. $\eta_i(s_0)$ takes values 2, 4, 6, 8 for nonadiabatic electrons. Debye shielding is not included for these runs. Lower panel: the corresponding normalized growth rates $\gamma$. 

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Figure 4.8: Mode structures for the ITG mode for $\eta_h(s_0) = 2.0$ (left) and ETG mode for $\eta_e(s_0) = 2.0$ with (middle) and without (right) Debye shielding on the poloidal cross section of a tokamak for the maximum growth rates. The other species (electron for ITG, ion for ETG) is considered adiabatic.
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Figure 4.9: Closeup view of Fig. 4.8

Figure 4.10: Poloidal Fourier harmonics for the modes shown in Fig. 4.8
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batically, i.e., taking into account all the kinetic effects of ions from FLR to various resonances and keeping \( \eta_i \) above the threshold of the ITG drive. Figure 4.5 depicts the mode frequency and growth rate of the ETG mode against \( k_0 \rho_i \) (Note that \( k_0 \) is normalized with \( \rho_i \) here.) with (i) adiabatic ions and (ii) nonadiabatic ions for increasing \( \eta_i \) for the case without Debye shielding. The corresponding toroidal mode numbers are displayed on the upper axis. It is observed that the nonadiabatic ions have very weak effect on the growth rates of the mode. In the second case, we have carried out same kind of scan but with Debye Shielding effect included. Figure 4.6 shows the real frequency and growth rate for the case of the ETG mode when Debye shielding is taken in to account. A comparison of the real frequency and growth rate with the case without Debye shielding reveals that the ion dynamics have visible albeit weak effect on the mode frequency as well as growth rate of the mode. The nonadiabatic ions tend to lower the growth rate of the ETG mode as one increases the ITG drive by increasing \( \eta_i \) of ions in the long wavelength side of the ETG mode corresponding to low \( k \) tail of the spectrum. For the purpose of comparison, we present a \( k_0 \rho_i \) scan for the ITG mode (from Chapter 2) including the adiabatic and nonadiabatic electrons, with \( \eta_e \) increasing for the latter case. Figure 4.7 depicts the real frequency as well as corresponding growth rates for the ITG mode. While the ITG mode is influenced strongly by the nonadiabatic electrons with substantial change in the growth rate, the ETG mode, on the other hand, is weakly affected by the nonadiabatic ions. Thus, one can draw a conclusion that even within the electrostatic limit, without taking into account the trapped species, the isomorphism of the ITG and ETG mode breaks, when the other species (electrons for the ITG mode and ions for the ETG mode) is considered fully kinetically. The corresponding mode structures of the ITG and ETG mode with and without the Debye shielding effect are compared in the top panel of Fig. 4.8 for the maximum growth rates of the respective cases. The figure delineates the vast difference in the scale-lengths of the two modes. While the ITG mode prevails over a considerable fraction, ETG mode, on the other hand, is restricted to a small annular ring in the poloidal cross section of the tokamak. The middle panel shows a closeup view of the ITG and ETG mode structures. The poloidal Fourier components for the three cases are displayed in the bottom panel.
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4.2.3 Mixing length estimate of flux

It would be interesting to see, how one mode's scale-dynamics have effect on the transport of the other or in other words how the ITG scale can affect the ETG scale and vice versa. For that purpose, simple mixing length estimate of transport has been evaluated. Before going to the results, we would like to add a line of caveat regarding the mixing-length based estimation of transport. The earlier results showed that the calculation of transport using mixing-length theory gives very low level of transport of electrons. Nonlinearly these modes generate streamers, by means of which, the ETG mode can give rise to experimentally relevant level of transport. Nonetheless, within our scope, we present a qualitative picture of the transport with the mixing-length calculation.

Pure ETG: a comparison with ITG

Figure 4.11 presents the heat diffusivity in the electron gyro-Bohm unit for the pure ETG mode without (top panel) and with (middle panel) the Debye shielding effect for the multiple cases of increasing $\eta_i$ of nonadiabatic ions and with adiabatic ions. It is apparent from the figure that even though the nonadiabatic ions have very weak effect on the growth rate of the ETG mode, they can change the heat diffusivity of the electrons substantially. The diffusivity peaks toward the low $k$ side of the spectrum but not at $k$ where the growth rate peaks, and decreases as the $\eta_i$ increases from 2 to 8, in steps of 2 which are all above the ITG threshold. One can hence conclude that the ion scale drive in the low $k$ regime reduces the high $k$ ETG drive, even if it has weaker effect on the mode frequency and growth rate of the ETG mode. The bottom panel displays the nonadiabatic electron effect on the ITG mode, where $\eta_e$ of the electrons increases gradually from 2 to 8. This leads to the reduced ion diffusivity for the ITG case in the presence of an electron drive with $\eta_e$ being above the ETG threshold. Thus, the comparison of the three cases depicted in Fig. 4.11, leads to the interesting conclusion that the ETG drive on the high $k$ side tends to reduce the ion transport on the low $k$ side, while the ITG drive on the low $k$ side tends to reduce the electron transport on the high $k$ side of the spectrum. One other conclusion is that in spite of exhibiting disparate scales, one mode can have effect on the other mode, putting a caveat to the tacit assumption that one mode is independent of the other because of the
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Figure 4.11: Top panel: mixing length estimate for transport coefficient $D_{ML} = \gamma / < k_{\perp}^2 >$ in electron gyro-Bohm units as a function of $k_{\theta}\rho_i$ for the ETG mode without Debye shielding at $\eta_e(s_0) = 2$ with adiabatic ions and nonadiabatic ions for $\eta_i(s_0) = 2, 4, 6, 8$; Middle panel: mixing length estimate for transport coefficient $D_{ML} = \gamma / < k_{\perp}^2 >$ in electron gyro-Bohm units as a function of $k_{\theta}\rho_i$ for the ETG mode with Debye shielding at $\eta_e(s_0) = 2$ with adiabatic ions and nonadiabatic ions for $\eta_i(s_0) = 2, 4, 6, 8$; Bottom panel: mixing length estimate for transport coefficient $D_{ML} = \gamma / < k_{\perp}^2 >$ in ion gyro-Bohm units as a function of $k_{\theta}\rho_i$; $\eta_i(s_0) = 2$, with adiabatic electrons and nonadiabatic electrons for $\eta_e(s_0) = 2, 4, 6, 8$. 

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Figure 4.12: Left panel: normalized perpendicular wave numbers $k_r$, $k_\theta$, $k_\perp$ vs toroidal mode number $n$ for the ETG mode without Debye shielding. Right panel: normalized perpendicular wave numbers $k_r$, $k_\theta$, $k_\perp$ vs toroidal mode number $n$ for the ETG mode with Debye shielding.

Vast spatio-temporal differences in their respective scales. The fact, in spite of the weak effect on the mode frequency and growth rate of the ETG mode by the nonadiabatic ions, that the mixing length estimates show a considerable reduction in the electron heat diffusivity has its origin in the change in the perpendicular scale lengths brought about by the nonadiabatic ions. Figure 4.12 plots the $k_r \rho_e$, $k_\theta \rho_e$ and $k_\perp \rho_e$ against toroidal mode number $n$. It is apparent that $k_r \rho_e$ of the ETG mode rises in the presence of the nonadiabatic ion dynamics, so does $k_\perp \rho_e$. This leads to the reduction in the mixing length estimates of electron transport.
4.3 Toroidal Universal Drift Instability: A Global Gyrokinetic Study

The confinement of electrons is marred not only by the temperature gradient driven instabilities such as ETG modes described in the preceding section but also by the density gradient driven instabilities. The density gradient also can equally be the source of free energy for instabilities in the case of no temperature gradient or very weak temperature gradient. These instabilities in tokamaks are called the toroidal universal drift instabilities. A brief review of earlier works on this topic can be found in the section 1.2 of Chapter 1.

In the present section, various parametric studies of the toroidal branch of the universal drift mode are carried out. We observe finite mode frequencies and growth rates beyond the critical \( \eta = L_n/L_T \) for the Ion Temperature Gradient (ITG) and Electron Temperature Gradient (ETG) modes, where \( L_n \) and \( L_T \) are respectively, the density and temperature scale lengths. A comparative study of the contribution of the magnetic drift resonance as well as of the Landau resonance from both species towards the stability properties of the mode is performed by a systematic parametric scan. An electromagnetic study of the mode is also carried out that elucidates the effect of finite \( \beta \) on the universal drift mode driven by toroidicity. The effect of trapped electrons on the universal mode is studied and growth rates and real frequencies are compared with the ion temperature gradient and trapped electron modes.

4.3.1 Model equations

The electrostatic formulation has been discussed in Chapters 2 and 3. Here we will elaborate the electromagnetic formulation only. For the electromagnetic case the perturbed density is modified as [48]

\[
\bar{n}_j(r; \omega) = - \left( \frac{q_j N}{T_j} \right) \left[ \tilde{\varphi} + \int d\mathbf{k} \exp (i \mathbf{k} \cdot \mathbf{r}) \times \int d\mathbf{v} \frac{f_{Mj}}{N} (\omega - \omega_j^*) \left( v P_j \right) \left[ \tilde{\varphi}(\mathbf{k}; \omega) - v_\parallel \tilde{A}_\parallel(\mathbf{k}; \omega) J_0^2(x) \right] \right], \quad (4.2)
\]
where $\tilde{A}_\parallel$ is the component parallel to the equilibrium magnetic field of the vector potential associated with the perturbation. The other terms are defined in Chapter 2 and 3. In addition to $\tilde{n}_j$, one has to consider the fluctuation of the parallel current density given by

$$\tilde{j}_\parallel(r; \omega) = -\left(\frac{q_j^2}{T_j}\right) \left[ \int dk \exp(ik \cdot r) \times \right.$$

$$\left. \int v_{\parallel} dv f_{Mj} (\omega - \omega_j^p) (iP_j) \left[ \tilde{\varphi}(k) - v_{\parallel} \tilde{A}_\parallel(k) \right] J_0^2(x_{Lj}) \right] \quad (4.3)$$

Along with the quasineutrality condition Eq. (2.10), Ampere’s law

$$\frac{1}{\mu_0} \nabla^2 \tilde{A}_\parallel = - \sum_j \tilde{j}_\parallel$$

will finally close the set of equations, to give a linear system of equations of the form

$$\sum_{k'} \sum_{j=\pm} \mathcal{M}_{j,k,k'} \left( \frac{\tilde{\varphi}_{k'}}{\tilde{A}_\parallel_{k'}} \right) = 0$$

Simple diagnostics for various physical quantities are computed as averages over the eigenmode. For example mode-averaged $k^2_\theta$ for the electrostatic case is computed as

$$\langle k^2_\theta \rangle = \frac{\int d\rho \sum_m \left| \frac{m}{\rho} \varphi_{(\rho,m)} \right|^2}{\int d\rho \sum_m \left| \varphi_{(\rho,m)} \right|^2}. \quad (4.4)$$

The above shown averaging procedure is suitably extended to the electromagnetic cases by including $\tilde{A}_\parallel$ mode structure averaging as follows:

$$\langle k^2_\theta \rangle = \frac{\int d\rho \sum_m \left| \frac{m}{\rho} \varphi_{(\rho,m)} \right|^2 + \int d\rho \sum_m \left| \frac{m}{\rho} \tilde{A}_{\parallel,(\rho,m)} \right|^2}{\int d\rho \sum_m \left| \varphi_{(\rho,m)} \right|^2 + \int d\rho \sum_m \left| \tilde{A}_{\parallel,(\rho,m)} \right|^2}. \quad (4.5)$$
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4.3.2 Results and Discussion

In the present section, we will delineate the results from the global linear gyrokinetic numerical analysis. It is to be noted that the mode frequencies and growth rates are expressed in units of $v_{thi}/a$ throughout the paper.

![Equilibrium profiles](image)

Figure 4.13: Equilibrium profiles to study the global toroidal universal drift instability mode (for parameters in Table I): (a) normalized density (dots), temperature (circle), $\eta_\perp,\epsilon$ (triangle), (b) Safety factor $q$ (circle) and magnetic shear $\dot{s}$ (dots) profiles as functions of normalized radius $s = r/a$. Note that $q(s_0) = 2.0$, $\dot{s}(s_0) = 0.40$, $\epsilon_n(s_0) = 0.1$, and $\tau(s_0) = 3.0$ for $s_0 = 0.6$.

Profiles and Parameters:

Let us consider the profiles and parameters as displayed in Table 4.2 for a plasma with single charged ions. The equilibrium profiles corresponding to these parameters are shown in Fig. 4.13.

Growth Rate $\gamma$ and Real Frequency $\omega_r$ vs $k_0\rho_{Li}$

The dispersion diagram for the toroidal universal drift instability with real frequency and growth rate plotted versus the normalized poloidal wave number $k_0\rho_{Li}$, is shown in Fig. 4.14. The upper axis presents the corresponding toroidal mode
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Table 4.2: Profiles and parameters

Parameters:
- B-field : \( B_0 = 1.0 \text{ Tesla} \)
- Temperature : \( T_0 = T(s_0) = 7.5 \text{ keV} \)
- Major Radius : \( R = 2.0 \text{ m} \)
- Minor Radius : \( a = 0.5 \text{ m} \)
- radius : \( s = \rho/a, \ 0.01 < s < 1.0, \ s_0 = 0.6 \)
- \( L_{n0} = 0.2 \text{ m} \)
- \( \eta_{h,e}(s_0) = 0.0, \ \epsilon_n = L_{n0}/R = 0.1 \)
- \( \tau(s) = T_e(s)/T_i(s) = 3. \)

Equilibrium Profiles:
- N-profile and T-profile
  \[
  \frac{N(s)}{N_0} = \exp \left( -\frac{a \delta s_n}{L_{n0}} \tanh \left( \frac{s-s_0}{\delta s_n} \right) \right)
  \]
  \[
  \frac{T_i,e(s)}{T_0} = \exp \left( -\frac{a \delta s_T}{L_{T0}} \tanh \left( \frac{s-s_0}{\delta s_T} \right) \right)
  \]
  \( \delta s_n = 0.35, \ \delta s_T = 0 \) at \( s = s_0 \)
- \( q(s) = 1.691 + 0.603 s^2 + 0.705 s^4 \)
  such that \( q(s = s_0) = 2.0; \)
  shear \( s = s_0, \ s = 0.4. \)

![Figure 4.14: Real frequency and growth rate for the electrostatic case corresponding to the parameters in the Table I and profiles shown in Fig. 4.13.](image)

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numbers. The real frequency is in the electron diamagnetic drift direction. The real frequency at very low $k_\theta \rho_{Li}$ increases first and then peaks at the value of $k_\theta \rho_{Li} = 0.4$ corresponding to $n \approx 6$. After this point, it starts falling with $k_\theta \rho_{Li}$ in a monotonic way. The growth rate, on the other hand, initially increases until the point $k_\theta \rho_{Li} \simeq 0.58$ corresponding to the toroidal mode number $n \approx 10$ and is practically constant at larger value of $k_\theta \rho_{Li}$. It is apparent from this observation that the toroidal branch of the universal drift instability spans from the low wave number or longer wavelength regime, where ion dynamics are dominant, to higher wave number or shorter wavelength regime, where usually electron dynamics play the dominant role. This is in contrast to the observation in the slab case [99], which is marginally stable at high wave number regime and damped in the low wave number regime.

**Electrostatic mode structure**

In the present section, we discuss the global electrostatic mode structure of the toroidal branch of the universal drift mode. Figure 4.15 displays the potential contours on a poloidal cross section of the tokamak in the upper left panel for $k_\theta \rho_{Li} = 0.58$ corresponding to the toroidal mode number $n = 10$. The various poloidal components of the potential with coupling brought about by the toroidicity, both in Fourier and real space, are presented in the upper right panel for $k_\theta \rho_{Li} = 0.58$ corresponding to the toroidal mode number $n = 10$. A few important points to be noted in this context are: (1) The mode structure is quite global passing through many mode rational surfaces. (2) It exhibits a weak ballooning character, with a finite amplitude observed at the favourable curvature side (high field side). In the upper right panel, coupling of poloidal components has been shown across the minor radius, with a maximum amplitude at $s=\rho/a=0.6$, where the density gradient peaks. The points in the upper axis, labelled by the corresponding poloidal mode numbers $m$, display the position of the mode rational surfaces where $k_{\parallel}(m,n) = 0$. Corresponding to each of these points, one can see a dip in the potential corresponding to each poloidal mode number. These dips correspond to $k_{\parallel}(m,n) = 0$ surfaces where $|\omega/k_{\parallel}| >> v_\text{th,e}$, $v_\text{th,e}$ being the electron thermal velocity. Thus, the strong effect of the off-resonant electrons is clearly visible from this figure. The convergence in the Fourier space for the considered
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Figure 4.15: Upper panel: (Left) The electrostatic mode structure for toroidal mode number \( n = 10 \), \( k_\theta \rho_{Li} = 0.58 \), corresponding to the parameters in the Table I and profiles shown in Fig. 4.13. (Right) poloidal component of \( \tilde{\phi} \) in (top) radial Fourier representation and (bottom) radial direct space. Lower panel: A closeup view of the mode structure.
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mode is presented in the upper part of the upper right panel of Fig. 4.15. Lower panel presents a closeup view of the mode considered here.

![Graphs showing the effect of electron and ion Landau resonance](image)

Figure 4.16: The effect of electron and ion Landau resonance for the mode $n = 10$ corresponding to $k_0 ho_{Li} = 0.58$. This is done by weighting the Landau resonance term by $\alpha$ and running it from 0 to 1 for one species, and keeping $\alpha = 1.0$ for the other species and vice versa.

**Effect of Landau Resonance**

To investigate the effect of the Landau resonance of electrons and ions on the toroidal universal drift instability, one can artificially put a multiplying factor, say $\alpha$, in front of the $k_\parallel v_\parallel$ term in the denominator of the propagator for both species [see Eq. (2.8)] and decrease it gradually from 1 to 0, once for ions, keeping full electron Landau resonance effect, and vice versa. It is to be noted that only the values 0 and 1 of the artificial factor $\alpha$ are physically meaningful. The other values of $\alpha$ simply represent a fractional weight to the Landau resonance term in the propagator so as to enable us to track numerically the Landau effect continuously. Thus the value 1 will refer to the case of full Landau resonance term taken into account and 0 the complete omission of the Landau resonance term from the propagator. It is clear from Fig. 4.16 that the ion Landau resonance apparently has no significant effect on the growth rate as well as on the mode frequency compared to that of the electrons. For the electron Landau resonance, the growth rate exhibits
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![Graphs showing ion and electron magnetic drift resonance for mode n=10](image)

Figure 4.17: Effect of ion and electron magnetic drift resonance for the mode n=10 corresponding to \( k_\rho \rho_L = 0.58 \). This is done by weighting the magnetic drift term by \( \alpha \) and running it from 0 to 1, and keeping \( \alpha = 1.0 \) for the other species and vice versa.

a nonmonotonic dependence on the electron Landau resonance weighting parameter \( \alpha \). For example, at lower values of \( \alpha \) the growth rate increases and again falls at higher values. Regarding the mode frequency: in contrast to the growth rate, it increases monotonically with \( \alpha \) for the electrons, while it has little variation in the case of ions. Thus, with the complete omission of the electron Landau resonance, the mode may become nonexistent even if one keeps the other destabilizing factors intact.

Effect of Magnetic Drift Resonance

In toroidal geometry, a mode will certainly have magnetic drift resonance if its frequency is of the same order as the magnetic drift frequency. We have looked at the effect of magnetic drift resonance for both species on the toroidal universal mode. This is done in a similar way as for the study of the previous section, by putting a multiplying factor \( \alpha \) in front of the magnetic drift resonance term, \( x_{ij} \) appearing as arguments of the Bessel functions in the numerator of the propagator in Eq. (2.8). To be noted again that, while doing the scan for one species, the full weight \( \alpha = 1 \) is kept for the other species. One notes a destabilizing effect due to the magnetic drift resonance of ions when the multiplying factor \( \alpha \) increases. As shown in Fig. 4.17, the real frequency decreases with the increasing multiplying...
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![Graphs showing shear vs. \(\omega_{a/V_{th}}\) and \(\gamma_{a/V_{th}}\) in the shear vs. shear plane.]

Figure 4.18: Shear scan for the mode \(n = 10\) corresponding to \(k_0 \rho_{Li} = 0.58\) at position \(s = s_0 = 0.6\), where the density gradient peaks. For these scans the safety factor at \(s = s_0\) is kept at the fixed value \(q = 2.0\).

The factor \(\alpha\), while the growth rate increases and starts saturating as one approaches \(\alpha = 1\), i.e., for the full magnetic drift resonance term in the ion propagator. Electron magnetic drift resonance however does not play a significant role for the universal drift instability, as is clear from Fig. 4.17.

**Shear Scan**

Though magnetic shear has a strong stabilizing influence on the universal drift mode, in a slab model, the toroidal resonance effects, intrinsic to toroidal geometry, can reduce the shear damping of the mode. The shear scan is displayed in Fig. 4.18. It is clear that for the parameters chosen in this study, the critical value of shear, beyond which the mode gets stabilized, is of the order of 1. The growth rate and real frequency, however, do not decay monotonically with increasing shear. The growth rate rather increases weakly at low shear and then starts decreasing with increasing shear. Similarly the real frequency of the mode also decays with increasing shear. The stabilization of the mode by shear in the presence of finite toroidicity can be understood as follows. The shear damping of the universal drift mode in the slab geometry is basically due to the convection of energy away from the mode rational surfaces. In the presence of finite toroidicity, the toroidal coupling effect inhibits this convection of wave energy, and thus reduces the shear damping [101]. Similar effects of toroidicity on shear induced stabilization was also
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Figure 4.19: Real frequency and growth rate for $\epsilon_n = L_n/R$ for the parameters and profiles as in the Table I and Fig. 4.13 in the case of mode $n = 10$ corresponding to $k_0 \rho_{Li} = 0.58$. Note that $a$, $L_n$, $Rq$ and $nq$ are kept constant in this scan.

observed for high-$n$ toroidal universal drift instabilities.

Toroidicity Scan

The real frequency and growth rate of the mode versus $\epsilon_n = L_n/R$ is presented in Fig. 4.19. The toroidicity scan is done by varying $R$, but keeping $Rq$, $L_n$, $nq$, and $a$ constant. While the real frequency decreases almost monotonically, the growth rate, on the other hand increases first with toroidicity, peaks at around $\epsilon_n \simeq 0.1$, and then starts falling for larger values of $\epsilon_n$. Since, $\epsilon_n \rightarrow 0$ implies $R \rightarrow \infty$, i.e., the cylindrical limit, the toroidal driving term becomes weak at low $\epsilon_n$. Since the mode is basically driven by the magnetic drift resonance, the growth rate increases with $\epsilon_n$ for low values. However, for large enough values of $\epsilon_n$ the mode becomes off-resonant with respect to the magnetic drift frequency, which increases with $\epsilon_n$ and growth rate falls down. Also, increasing $\epsilon_n$ implies decreasing $R$, which means reduced connection length $\sim Rq$ between the favorable and unfavorable magnetic field. This nonmonotonic dependency of growth rate on $\epsilon_n$ for toroidal universal drift modes has not been reported before.
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\[ \tau = T_e/T_i \] Scan

The dependence of the mode frequency and growth rate for the toroidal universal drift instability on the temperature of the species is elucidated in the present section. Figure 4.20 displays the plots of the mode frequency and growth rate corresponding to \( k_\theta \rho_{Li} = 0.58 \) (\( n = 10 \)) as a function of the ratio of temperatures of electrons and ions, i.e., \( T_e/T_i = \tau \). To be noted that in this scan the ion temperature \( T_i \) is kept constant, while varying only the electron temperature \( T_e \). The real frequency increases monotonically with the magnitude of \( \tau \), i.e., with increasing electron temperature. The growth, on the other hand, exhibits a nonmonotonic character: increases at first with \( \tau \), peaks at around \( \tau = 5.0 \) and then starts decaying with increasing \( \tau \). One may correlate this result with the role of electron Landau resonance on the universal mode, as the electron distribution in the vicinity of parallel resonant velocity, which is strongly dependent on the thermal velocity and so the temperature of the electrons, is the key factor in determining Landau damping or inverse Landau damping of the mode.

**Effect of Temperature Gradient**

Thus far, the entire analysis has been carried out considering flat temperature profiles, that is, zero temperature gradients by putting \( \delta s_T = 0 \) in the profiles.
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Figure 4.21: Real frequency and growth rate in the case of a temperature gradient scan for the parameters and profiles of Table I and Fig. 4.13 for the mode \( n = 10 \) corresponding to \( k_d \rho_{Li} = 0.58 \). The flat temperature profile in Table I, has been replaced by one with \( \delta s_T = 0.2 \) instead of 0 for the previous cases.

displayed in Table I. In the present section, we incorporate profile variation to the temperatures of both ions and electrons. This is achieved by using a finite value for \( \delta s_T \), which is chosen as 0.2 in this case. Since most tokamaks contain temperature gradients in the pressure profile, it is thus necessary to look at the effect of the temperature gradient on the toroidal universal drift instability. This is done by evaluating the real frequency and growth rate against \( \eta_{e,i} = L_n/L_T \), keeping \( L_n \) constant and varying \( L_T \). Three cases are considered here: (1) the temperature gradient scale lengths for both electrons and ions are increased simultaneously, (2) the temperature gradient of only ions is increased, keeping that of the electrons zero, and (3) the temperature gradient of only electrons is increased, keeping that of the ions zero. The last two options may be relevant to experimental situations with preferential ion heating [e.g., ion cyclotron resonance heating (ICRH)] or electron heating [e.g., electron cyclotron resonance heating (ECRH)], respectively. The results for all the cases are presented in Fig. 4.21. The real frequency is reduced with increasing temperature gradient for the cases 1 and 2, while it increases in the case 3. The growth rate, on the other hand, decreases in all three cases. However, it decays more slowly when the electron temperature profile is flat, as apparent from case 2. It is clear from case 3 that the mode exhibits a finite growth rate for values of \( \eta_e \) beyond the critical value for ETG instability (\( \eta_e \simeq 1.0 \)). Thus,
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Figure 4.22: Real frequency and growth rate for the electrostatic (dashed curve) and electromagnetic case (solid curve) for the parameters in Table I and profiles as shown in Fig. 4.13. The value of $\beta$ considered here is 0.001.

The point to be noted is that, even in the presence of finite $\eta$ above the critical value for the temperature gradient driven modes to get destabilized, the universal drift instability preserves finite growth rate. It is observed from case 2 that the ion temperature gradient has weaker effect on the mode. It has finite growth rate even after the critical value of $\eta_i$ for the ion temperature gradient driven mode ($\eta_i \approx 1.0$). Thus, one may conclude that, in some situations as delineated in the last two cases, temperature gradient driven modes and universal drift mode driven by the density gradient can coexist. One other important point to be noted here is that the toroidal universal drift mode is unstable in the $\tau$ domain from 1 to more than 10, as evident from Fig. 4.20, while ETG modes are stable at higher values of $\tau$ even with finite $\eta_e$. So electron transport at high $k_{\perp}\rho_{Li} > 1$ with larger values of $\tau$ may have contributions from the toroidal universal drift mode as well. Similarly, at low $k_{\perp}\rho_{Li}$, where ITG is dominant, the electron transport can be due to this mode, as it appears to be unstable in regions where ITG is pronounced.

Electromagnetic Effects

$k_{\theta}\rho_{Li}$ scan

The electrostatic assumption is justified in a low $\beta$ plasma. However, considering the higher $\beta$ environment in the present day devices, it is of interest to study the effect of the electromagnetic fluctuation on the toroidal branch of the universal
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mode. In this section, we perform a \( k_\theta \rho_{Li} \) scan for the growth rate and mode frequency in the presence of transverse magnetic perturbations with \( \beta = 0.001 \) in the zero temperature gradient limit. The corresponding results are plotted in Fig. 4.22, with the upper axis representing the respective toroidal mode numbers \( n \). For the purpose of comparison, we also plot the purely electrostatic values for the real frequency and growth rate of the mode (same results as in Fig. 4.14).

At very low \( k_\theta \rho_{Li} \), the real frequency increases first and then peaks at a value of \( k_\theta \rho_{Li} \simeq 0.4 \) corresponding to \( n = 6 \). Beyond this point, the frequency starts falling monotonically with \( k_\theta \rho_{Li} \) (or \( n \)). The growth rate however increases until \( k_\theta \rho_{Li} \simeq 0.58 \) (or \( n = 10 \)) and then becomes practically constant. It is apparent from the figure for the real frequency that the effect of finite \( \beta \) is more pronounced at lower \( k_\theta \rho_{Li} \), and the real frequency is reduced in this region. Going towards the higher \( k_\theta \rho_{Li} \) the effect of \( \beta \) seems to be weaker on the mode frequency. The growth rate, on the other hand, is substantially reduced by finite plasma \( \beta \). A \( \beta \) of value 0.001 brings almost 20% reduction in the growth rate as compared to the electrostatic case. A complete \( \beta \) scan is presented in the following section clearly illustrating the stabilizing effect of \( \beta \).

A global mode structure for the electromagnetic case for \( n = 10 \) and \( \beta = 0.001 \) corresponding to \( k_\theta \rho_{Li} \simeq 0.58 \) is shown in Figs. 4.23 and 4.24. The electrostatic part \( \hat{\phi} \) (Fig. 4.23) is very similar to the purely electrostatic mode in Fig. 4.15. The \( \hat{A}_\| \) component (Fig. 4.24), on the contrary, apparently shows a weak antiballooning character, being weaker at the outboard side than the inboard side. The convergence in the radial and poloidal Fourier space for the mode has been depicted in the upper part of the right panel in Fig. 4.24. The lower panel presents the radial dependence of various poloidal mode numbers \( m \). It retains the effect of nonresonant electrons at \( k_\| (m, n) = 0 \) surfaces. The antiballooning character of the \( \hat{A}_\| \) mode structure and the stabilization of the mode in the presence of finite \( \beta \) are all related to the inherent electrostatic nature of the toroidal universal drift instability.

**\( \beta \) Scan**

A complete \( \beta \) scan for the mode with \( k_\theta \rho_{Li} = 0.58 \) (\( n = 10 \)) is displayed in Fig. 4.25. Both real frequency and growth rate are reduced with increasing \( \beta \). The complete stabilization occurs at \( \beta \simeq 1.1\% \). This is in contrast to earlier investigations in slab geometry, where the value of critical \( \beta \) was much higher.
Figure 4.23: (Upper panel) The global mode structure for the $\tilde{\phi}$ component in the poloidal cross section in the electromagnetic case for $n = 10$, $k_0 \rho_{Li} = 0.58$, and $\beta = 0.001$. (Lower panel) Poloidal component of $\tilde{\phi}$ in (top) radial Fourier representation and (bottom) radial direct space.
Figure 4.24: (Upper panel) The global mode structure for the $\tilde{A}_||$ component in the poloidal cross section in the electromagnetic case for $n = 10$, $k_B \rho_{Li} = 0.58$, and $\beta = 0.001$. (Lower panel) Poloidal component of $\tilde{\phi}$ in (top) radial Fourier representation and (bottom) radial direct space.
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Figure 4.25: $\beta$ scan for the mode frequency and growth rate for the parameters and profiles as in Table I and Fig. 4.13 for the mode $n = 10$ corresponding to $k_\theta \rho \approx 0.58$.

[105]. The observed stabilization is perhaps due to the coupling of the wave with the Alfvén perturbation.

The relative strength of the electromagnetic to the electrostatic character is shown

Figure 4.26: Electromagnetic ratio with increasing function of $\beta$ for the parameters and profiles as in Table I and Fig. 4.13 for the mode $n = 10$ corresponding to $k_\theta \rho \approx 0.58$.

in Fig. 4.26, measured as the ratio of flux surface averaged squared $\tilde{A}_\parallel$ to $\tilde{\phi}$ with increasing value of $\beta$ expressed in percentage. It is clear from this plot that the strength of the magnetic fluctuation in comparison to the electrostatic fluctuation
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increases almost linearly with increasing plasma \( \beta \).

Effect of trapped electron on the electrostatic mode:

As clearly demonstrated by the above results, the global toroidal universal drift instability is triggered by purely passing electron dynamics. However, as a toroidal device is bound to have some fraction of trapped electrons, it would be interesting to obtain the effect of trapped electrons on the purely universal drift mode studied in the previous sections. To this end, an extensive \( \eta \) scan is performed with and without trapped electrons for the electrostatic case. To identify the most unstable mode, the other drift instabilities such as ITG with trapped electrons as well as the TEM branches are computed together with the universal mode. The combined data is plotted in Fig. 4.27. There are several interesting points to be noted: (1) The trapped electron coupled universal drift mode and pure universal drift mode have distinct real frequencies and growth rates. (2) In contrast to the pure universal drift mode studied in previous sections, whose growth rate was shown to decrease with increasing \( \eta \), the trapped electron coupled universal drift mode appears to be more unstable with increasing \( \eta \). This study indicates that in the presence of trapped electrons, the nature of the universal drift mode is predominantly “trapped electron like”. (3) To make a better quantitative comparison, growth rates and real frequencies of the pure trapped electron mode (TEM), the ion temperature gradient mode with trapped electrons (ITG-TE) for the same equilibrium parameters are also plotted. For the parameters studied here, it appears that the trapped electron coupled universal drift modes in the presence of \( \eta \geq 1 \) have growth rates comparable to ITG-TE or TEMs and could contribute substantially to the overall transport.

4.4 Conclusions

In the present work, we have presented some features of the electron temperature gradient driven (ETG) mode using a linear gyrokinetic model in toroidal geometry that treats both species, namely, ions and electrons fully gyro-kinetically, taking into account all the kinetic effects. The effect of Debye Shielding, breaking of isomorphism of ITG and ETG modes even in the electrostatic limit when the
Figure 4.27: The real frequency and growth rate vs temperature gradient for different unstable modes in the presence of trapped electrons in the same regime defined by the parameters and profiles of Table I and Fig. 4.13 for the mode \( n = 10 \) corresponding to \( k_B \rho_L l_i = 0.58 \). The flat temperature profile in Table I, has been replaced by one with \( \delta_T = 0.2 \) instead of 0 for the previous cases. The three dashed curves (circle, square, diamond) are for universal mode without trapped electrons (same as Fig. 4.21), three solid lines (circle, square, diamond) are for universal mode in the presence of trapped electrons (UNV-TE), the dashed curve with triangles is for ion temperature gradient mode with trapped electrons (ITG-TE) and the solid curve with stars is for pure trapped electron mode (TEM).
other species is considered fully gyrokinetically are revealed one by one. In the following we summarize the results obtained.

- Debye Shielding is stabilizing to the ETG mode, enhances the threshold in $\eta_e$, and wipes out the high $k$ tail of the ETG spectrum.

- Nonadiabatic ions have very weak effect on the growth rate and mode frequency of the pure ETG modes. In contrast, nonadiabatic electrons affect the mode frequency and growth rate of the ITG mode strongly. It breaks the supposed isomorphy between the two modes even in the electrostatic limit.

- We have presented an estimation of the transport of ions and electrons on the basis of mixing length theory. Results reveal that, drive for the ion channel tends to reduce the transport via the electron channel and vice versa. This means that a low $k$ mode can have strong effect on a high $k$ mode and vice versa. The assumption of adiabatic particles fails to interpret this result.

Also, we have performed a global linear gyrokinetic study of the toroidal universal drift mode driven by the density gradient in the presence of finite toroidicity on the intermediate scale $k_{\perp} \rho_{Li}$. The model considers both passing electrons and ions to be fully nonadiabatic, incorporating toroidal coupling effects, magnetic drift resonances, Landau resonance effects, transit harmonic resonances, finite Larmor radius to all orders, and orbit width effect for both species. The effect of finite $\beta$ is also studied in the frame of an electromagnetic model that retains the transverse magnetic perturbation. However, effects of collisions and Shafranov shift have been dropped. Furthermore, the model considers large aspect ratio circular cross section for the tokamak plasma. The major results are as follows

- The growth rate increases at lower $k_\theta \rho_{Li}$ until $k_\theta \rho_{Li} \approx 0.58$ and starts saturating thereafter. The real frequency too increases at lower $k_\theta \rho_{Li}$ and then decays monotonically with $k_\theta \rho_{Li}$ at larger $k_\theta \rho_{Li}$.

- The electrostatic mode structure is global and exhibiting structure at mode rational surfaces.

- Studying the effect of Landau resonance for both electrons and ions shows weak dependence of the frequency and growth rate on ion Landau damp-
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...ing and a strong dependence on electron Landau damping preserving finite growth rate in both cases.

- Both electron and ion magnetic drift resonance terms are considered in the formulation. The ion magnetic drift resonance does play a significant role in making the toroidal branch of the universal drift mode unstable, while electron magnetic drift resonance has a weak effect on the stability property of the mode.

- As reported in earlier works for high $n$ modes ($k_\theta \rho_L > 1$), the toroidicity driven universal drift mode is found to be stable beyond a shear value $s \approx 1$, even for low $n$ modes.

- The nonmonotonic dependency of the growth rate on $\epsilon_n = L_n/R$ in a toroidicity scan (varying $R$ and keeping $L_n$, $a$, $Rq$ and $nq$ constant) is demonstrated here for the first time.

- The mode is unstable in a fairly large domain of $\tau = T_e/T_i$ ranging from 1 to more than 10, thus clearly showing that in regions of $\tau$ where the electron temperature gradient (ETG) mode is believed to be stable, electron transport can be due to this toroidal universal drift mode.

- The $\eta$ scan for both ions and electrons shows that the universal drift mode driven by toroidicity can coexist with the temperature gradient driven modes. Therefore, electron transport at low $k_\theta \rho_L$ may have contributions from the mode under investigation. Similarly at higher $k_\theta \rho_L$, where ETG is thought to be the main driving mechanism for electron transport, this mode may also contribute.

- The electromagnetic effect is found to be strongly stabilizing in the present case. The $A_\parallel$ component of the mode structure exhibits anti-ballooning character. The mode gets stabilized at $\beta \approx 1.1\%$. The relative magnetic fluctuation amplitude $<A_\parallel^2>/<\dot{\phi}^2>$ varies almost linearly with the magnitude of $\beta$.

- Trapped electrons enhance the growth rate of the universal mode. However, the universal mode changes its character regarding its dependence on
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the density and temperature gradients. While in the absence of trapped electrons the universal mode decays with the temperature gradient, trapped electrons, on the contrary, enhances the growth rate of the mode. The universal mode with trapped electrons exhibits, qualitatively, the same character as the trapped electron mode. Also, it has a comparable growth rate to the trapped electron coupled ion temperature gradient mode in the parameter range considered in this study.