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UNSTABLE $m = 1$ ECCENTRIC MODES IN THERMALLY COOLING, SELF GRAVITATING DISKS

Adam M. Dempsey & Yoram Lithwick
Center for Interdisciplinary Exploration and Research in Astrophysics (CIERA) and Department of Physics and Astronomy
Northwestern University 2145 Sheridan Road Evanston, IL 60208 USA

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ABSTRACT

1. INTRODUCTION

2. EQUATIONS OF MOTION

The equations of motion for a two dimensional, viscous disk are

$$(\partial_t + \mathbf{v}\cdot\nabla)\mathbf{v} = -\frac{\nabla P}{\Sigma} - \nabla\Phi + \frac{1}{\Sigma}\nabla\cdot(\nu\Sigma\mathbf{D}) \quad (1)$$

$$(\partial_t + \mathbf{v}\cdot\nabla)\Sigma = -\Sigma\nabla\cdot\mathbf{v} \quad (2)$$

$$T(\partial_t + \mathbf{v}\cdot\nabla)s = -\frac{T-\bar{T}}{t_{\rm cool}} \quad (3)$$

Where all quantities are assumed to be a vertical average of their three dimensional counterparts. Here, Σ is the surface density, s is the entropy of the fluid, T is the disk temperature, and $\mathbf{D}$ is the viscous stress tensor. These equations describe a viscous fluid that is thermally heated and or cooled to some "background" temperature profile $\bar{T}$ over a timescale $t_{\rm cool}$. To close the system of equations we assume the fluid is ideal and perfect, so that $P = \mathcal{R}\Sigma T$ and $s \equiv \frac{\mathcal{R}}{\gamma-1}\ln(P\Sigma^{-\gamma})$; hereafter we set $\mathcal{R} = 1$ without loss of generality. Following Ostriker, Shu, & Adams (2), we can connect the two dimension adiabatic index γ to the true three dimensional adiabatic index Γ by assuming that the gas is vertically isothermal and gaussian distributed.

We consider a disk with a prescribed azimuthally averaged background profile given by $(\bar{v}_r, \bar{v}_\phi, \bar{\Sigma}, \bar{T}) = (0, r\Omega(r), \bar{\Sigma}\propto r^\mu, \bar{T}\propto r^\delta)$. We look at linear perturbations to this background state. Denoted linear quantities with a prime and background quantities with a bar, the linearized equations of motion become

$$(\partial_t + \Omega\partial_\phi)v_r' - 2\Omega v_\phi' = -\frac{1}{\bar{\Sigma}}\partial_r P' + \frac{\Sigma'}{\bar{\Sigma}^2}\frac{d\bar{P}}{dr} - \partial_r\Phi' + \text{visc.} \quad (4)$$

$$(\partial_t + \Omega\partial_\phi)v_\phi' + \frac{\kappa^2}{2\Omega}v_r' = -\frac{1}{r\bar{\Sigma}}\partial_\phi P' - \frac{1}{r}\partial_\phi\Phi' + \text{visc.} \quad (5)$$

$$(\partial_t + \Omega\partial_\phi)\Sigma' + v_r'\frac{d\bar{\Sigma}}{dr} = -\bar{\Sigma}\left(\frac{1}{r}\partial_r(rv_r') + \frac{1}{r}\partial_\phi(v_\phi')\right) \quad (6)$$

$$(\partial_t + \Omega\partial_\phi)s' + v_r'\frac{d\bar{s}}{dr} = -\beta\Omega\frac{T'}{\bar{T}} \quad (7)$$

Where we've parametrized the cooling time as $t_{\rm cool}^{-1} = \beta\Omega$ (Gammie 2001). We now Fourier Transform the linear variables in the azimuthal direction as well as in time and write $(v_r', v_\phi', \Sigma', P', s', T', \Phi') = \sum_m (u, v, \sigma, p, s, T, \phi)e^{im(\phi-\Omega_p t)}$. We further work under the assumption that the different m modes only couple to the $m = 0$ mode (i.e the background) and not to other $m \neq 0$ modes. The linearized equations of motion are now

$$im(\Omega-\Omega_p)u - 2\Omega v = -\frac{1}{\bar{\Sigma}}\frac{dp}{dr} + \frac{\sigma}{\bar{\Sigma}^2}\frac{d\bar{P}}{dr} - \frac{d\phi}{dr} + \text{visc.} \quad (8)$$

$$im(\Omega-\Omega_p)v + \frac{\kappa^2}{2\Omega}u = -\frac{imp}{r\bar{\Sigma}} - \frac{im\phi}{r} + \text{visc.} \quad (9)$$

$$im(\Omega-\Omega_p)\frac{\sigma}{\bar{\Sigma}} + u\frac{d\ln\bar{\Sigma}}{dr} = -\left(\frac{1}{r}\frac{d}{dr}(ru) + \frac{imv}{r}\right) \quad (10)$$

$$im\left[\left(1 - i\frac{\beta(\gamma-1)}{m}\right)\Omega - \Omega_p\right]\frac{T'}{\bar{T}} + u\frac{d\ln\bar{T}}{dr} = -(\gamma-1)\left(\frac{1}{r}\frac{d}{dr}(ru) + \frac{imv}{r}\right) \quad (11)$$

Where we've replaced the entropy in favor of the temperature through the relations

$$(\gamma-1)s' = \frac{T'}{\bar{T}} - (\gamma-1)\frac{\sigma}{\bar{\Sigma}} \quad (12)$$

$$(\gamma-1)\frac{d\bar{s}}{dr} = \frac{d\ln\bar{T}}{dr} - (\gamma-1)\frac{d\ln\bar{\Sigma}}{dr} \quad (13)$$

3. $m = 1$ SLOW MODES

We specialize to $m = 1$ modes and to the low pattern speed limit $\Omega_p \ll \Omega$. We can simplify the 4 dimensional set of equations to just one equation describing the global structure and evolution of an eccentricity vector, $e(r)$ (5; 4; 3). The eccentricity is defined through the Lagrangian displacement ξ_r as

$$e \equiv \frac{\xi_r}{r} \quad (14)$$

Following (3) we can write an evolution equation for the disk eccentricity by an appropriate linear combination of the two momenta equations and using

$$\frac{\sigma}{\bar{\Sigma}} = -\frac{d\ln\bar{\Sigma}}{d\ln r}e - r\frac{de}{dr} \quad (15)$$

$$(1 - i\beta(\gamma-1))\frac{T'}{\bar{T}} = -\frac{d\ln\bar{T}}{d\ln r}e - (\gamma-1)r\frac{de}{dr} \quad (16)$$

Which follow from (10) & (11) to zeroth order in Ω_p and with $\Omega \approx \Omega_K$. We obtain the evolution equation

$$\Omega_p e = Ae + Br\frac{de}{dr} + Cr^2\frac{d^2e}{dr^2} \quad (17)$$

adamdempsey2012@northwestern.edu

Where the coefficients are,

$$A = \omega_p + \mathcal{H}\left[2\mu + \mu' + \eta(\delta' + \delta(2+\mu+\delta)\right] \quad (18)$$

$$B = \mathcal{H}\left[3 + \mu + \eta\left(\gamma\delta + (\gamma-1)(3+\mu)\right)\right] \quad (19)$$

$$C = \mathcal{H}\left[1 + (\gamma-1)\eta\right] \quad (20)$$

Where,

$$\eta \equiv \frac{1 + i\beta(\gamma-1)}{1+\beta^2(\gamma-1)^2} \quad (21)$$

$$\mathcal{H} \equiv \frac{\bar{T}}{2\Omega r^2} \quad (22)$$

$$\mu' = \frac{d^2\bar{\Sigma}}{d\ln r^2} \qquad \delta' = \frac{d^2\bar{T}}{d\ln r^2} \quad (23)$$

We solve (17) as a generalized eigenvalue problem by discretizing the eccentricity on a logarithmic grid in radius.

$$\Omega_p \mathbf{Q}\mathbf{e} = \mathbf{M}\mathbf{e} \quad (24)$$

The matrices $\mathbf{Q}$, and $\mathbf{M}$ are the boundary condition matrix and the coefficient matrix, respectively. We solve for the spectrum of normal eigenmodes, $\mathbf{e}_n$ and their associated (possibly complex) eigenvalues $\Omega_{p,n}$.

3.1. *Boundary Condition*

We adopt a zero Lagrangian pressure boundary condition at both the inner and outer boundaries. Using the equations of motion it can be shown that this is equivalent to satisfying

$$r\frac{de(r)}{dr}\bigg|_{r_i,r_o} = \left[\frac{1}{1+\left(\frac{\eta}{1-\eta}\right)\gamma}\right]\frac{d\ln\bar{T}}{d\ln r}e(r)\bigg|_{r_i,r_o} \quad (25)$$

at each boundary in the slow mode approximation. In the isothermal ($\beta \to \infty$, $\eta \to 0$) and adiabatic ($\beta \to 0$, $\eta \to 1$) limits the boundary condition reduces to,

$$r\frac{de(r)}{dr}\bigg|_{r_i,r_o} = \begin{cases} \frac{d\ln\bar{T}}{d\ln r}e(r)\Big|_{r_i,r_o} & \text{Isothermal} \\ 0 & \text{Adiabatic} \end{cases} \quad (26)$$

4. METHODS

5. RESULTS

We studied the growth rates of unstable modes in the slow mode ($\Omega_p \ll \Omega_k$) approximation. The classical WKB dispersion relation for a self gravitating and isothermal disk (see e.g Binney & Tremaine (1)) is

$$m^2(\Omega - \Omega_p)^2 = \kappa^2 - 2\pi G\Sigma|k_r| + c^2k_r^2 \quad (27)$$

In the $m = 1$ slow mode limit this simplifies to (3)

$$2\Omega(\omega_p - \Omega_p) = -2\pi G\Sigma|k_r| + c^2k_r^2 \quad (28)$$

Where $\omega_p = \Omega - \kappa$ is the orbital precession frequency of a fluid element. From the dispersion relation it is clear that pressure creates retrograde modes ($\Omega_p < 0$) and self gravity tends to create prograde modes ($\Omega_p > 0$). Additionally, global modes (i.e low k_r) also tend to create prograde modes (3).

5.1. *Lowest Order Mode with No Self Gravity*

The mode with zero nodes is perhaps the simplest to understand.

5.2. *Self Gravity*

6. DISCUSSION

7. CONCLUSIONS

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