Observational Signatures of Planets in Protoplanetary Disks: Temperature Structures in Spiral Arms

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Abstract
High-resolution imaging of protoplanetary disks has unveiled a rich diversity of spiral structure, some of which may arise from disk-planet interaction. Using 3D hydrodynamics with $\beta$ cooling to a vertically stratified background, as well as radiative-transfer modeling, we investigate the temperature rise in planet-driven spirals. In rapidly cooling disks, the temperature rise is dominated by a contribution from stellar irradiation, $0.3\%–3\%$ inside the planet radius but always $<0.5\%$ outside. When cooling time equals or exceeds dynamical time, however, this is overwhelmed by hydrodynamic $\frac{\partial U}{\partial t}$ work, which introduces a $\sim10\%–20\%$ perturbation within a factor of $\sim2$ from the planet’s orbital radius. We devise an empirical fit of the spiral amplitude $\Delta(T) = (M_p/\dot{M}_d)^{1/2}(\frac{\partial U}{\partial t}v_{\text{tor}}/c_{\text{s}} + f_{\text{rad}})$ to take into account both effects. Where cooling is slow, we find also that temperature perturbations from buoyancy spirals—a strictly 3D, nonisothermal phenomenon—become nearly as strong as those from Lindblad spirals, which are amenable to 2D and isothermal studies. Our findings may help explain observed thermal features in disks like TW Hydrae and CQ Tauri, and underscore that 3D effects have a qualitatively important effect on disk structure.

Unified Astronomy Thesaurus concepts: Protoplanetary disks (1300); Hydrodynamical simulations (767); Exoplanet detection methods (489)

1. Introduction
High-resolution, near-infrared scattered-light imaging instruments such as VLT/SPHERE and Gemini/GPI have revealed spiral features in a number of disks, including MWC 758 (Grady et al. 2013; Benisty et al. 2015) and SAO 206462 (Muto et al. 2012; Garufi et al. 2013; Stolker et al. 2016), among several others. More recently, spiral structures have been found in ALMA 12CO observations of TW Hydrae (Teague et al. 2019) and CQ Tauri (Wölfer et al. 2020). Various explanations have been proposed for these features (Dong et al. 2018), including, e.g., gravitational instability (Toomre 1964; Hall et al. 2019), while their ubiquity suggests that at least some of them may have a planet-driven origin. So far, however, direct imaging of planets in protoplanetary disks has proven challenging (e.g., Kepler et al. 2018; Müller et al. 2018), and connections between observed spirals and planets have been circumstantial and inconclusive (e.g., Ren et al. 2018; Wagner et al. 2019; Xie et al. 2020).

How, exactly, might planets generate spiral structure in their natal disks? One mechanism, extensively investigated in decades of analytical (e.g., Goldreich & Tremaine 1978, 1979; Lubow & Ogilvie 1998; Ogilvie & Lubow 2002) and numerical work (e.g., Kley 1999; Dong et al. 2011a; Zhu et al. 2015; Dong & Fung 2017; Hord et al. 2017; Zhang & Zhu 2020; Ziampras et al. 2020), is wave excitation at Lindblad resonances, where the local orbital frequency is a multiple of the Doppler-shifted forcing frequency of the planet. Goodman & Rafikov (2001) studied wake structure and angular-momentum transport in the low-amplitude, linear spirals excited by low-mass planets. Rafikov (2016) and Arzamasskiy & Rafikov (2018) generalized this picture to the arms created by superthermal mass companions. The superposition of the waves launched at Lindblad resonances constructively interfere to produce multiple arms, as studied by Bae & Zhu (2018a, 2018b) and Miranda & Rafikov (2019).

Less well studied are the arms launched at buoyancy resonances (e.g., Zhu et al. 2012; Lubow & Zhu 2014; McNally et al. 2020; Bae et al. 2021), where the Doppler-shifted planetary frequency commensurates with the local Brunt-Väisälä frequency. This is inherently a 3D effect requiring the solution of an energy equation, and is thus unaccounted for in commonly used 2D (Zhang & Zhu 2020; Ziampras et al. 2020) or 3D isothermal (Dong & Fung 2017) simulations. Zhu et al. (2012) and Zhu et al. (2015)’s 3D adiabatic simulations found weak buoyancy arms, but given that their temperature structure was vertically unstratified (unlike e.g., Juhász & Rosotti 2018), this finding is not necessarily applicable to real disks.

Owing to their isothermality, past hydrodynamical simulations of spiral arms have necessarily focused on density perturbations. But while visible in scattered light (Dong et al. 2015), density structures are far less pronounced in gas tracers like $^{12}$CO, which are optically thick for typical disk surface densities. In the present work, we therefore concentrate on the corresponding temperature perturbations, which may be more readily apparent in ALMA observations of gas emission.

Gas temperature structure in disks is shaped by two mechanisms. The first is hydrodynamics. As gas passes through spiral arms, compression and expansion subject it to $\frac{\partial U}{\partial t}$ work; this effect is suppressed when the gas-cooling timescale is much shorter than the arm-crossing time, but saturates in the adiabatic limit. Rafikov (2016) finds that for spiral density perturbations up to order unity, the $\frac{\partial U}{\partial t}$ work done on the gas is mostly reversible; any residual accumulation of heat over successive orbits is due to nonlinear/shock heating.
(Lyra et al. 2016) and is reflected in the gap-opening process, which is not the focus of our work.

Stellar irradiation also has an important role to play. The atmospheres of passively heated disks experience direct stellar illumination and thus reach high temperatures. Midplanes by contrast, are heated only indirectly and are therefore cooler. This gives rise to a vertically stratified equilibrium temperature structure. If a column of disk (in e.g., a spiral) at a given radial-azimuthal location has a greater scale height than its surroundings, it intercepts more starlight and acquires a greater temperature in vertical average, and vice versa.

In what follows, we use 3D adiabatic hydrodynamics with cooling, as well as radiative-transfer simulations, to better understand the temperature structures in spiral arms driven by planets in disks with a realistically vertically stratified temperature structure.

2. Methods

We conduct 3D simulations of disk-planet interaction using the GPU-accelerated, Lagrangian-remap hydrodynamics code PENQUIN (Fung 2015), which uses the third-order piecewise-parabolic method (Colella & Woodward 1984) to reconstruct quantities in solving the Riemann problem at cell boundaries. PENQUIN solves the viscous, compressible Navier–Stokes equations:

\[
\frac{D\rho}{Dt} = -\rho(\nabla \cdot \mathbf{v})
\]

(1)

\[
\frac{D\mathbf{v}}{Dt} = -\frac{\nabla p}{\rho} + \frac{1}{\rho} \nabla \cdot \mathbf{T} - \nabla \Phi
\]

(2)

\[
\frac{De}{Dt} = -\frac{p}{\rho} \nabla \cdot \mathbf{v} - \Lambda
\]

(3)

where \(\rho\) is the mass density, \(\mathbf{v}\) the velocity field, \(p\) the pressure, \(\mathbf{T}\) the viscous stress tensor, \(\Phi\) the gravitational potential, \(e \equiv p/(\gamma - 1)\rho\) the internal energy per unit mass, and \(\Lambda\) the cooling rate.

We adopt an adiabatic equation of state, \(p \equiv c_s^2 \rho / \gamma\), where \(c_s\) is the adiabatic sound speed and \(\gamma \equiv \partial \ln p / \partial \ln \rho = 1.4\) is the adiabatic index. \(\mathbf{T}\) is proportional to the kinematic viscosity \(\nu = \alpha c_s^2 \Omega^{-1}\), where \(\Omega\) is the local orbital frequency and \(\alpha\) the Shakura & Sunyaev (1973) parameter. We choose an \(\alpha = 10^{-3}\) in order to prevent the growth of large-scale disk instabilities, such as the Rossby wave instability. This relatively low level of viscosity is also motivated by observations of disks that have generally revealed a low level of turbulence (e.g., Flaherty et al. 2015; Teague et al. 2018).

Our beta-cooling prescription \(\Lambda\) relaxes the local sound speed to that of a fixed, vertically stratified, location-dependent background \(c_{s,0}(x)\) on a characteristic dimensionless timescale of \(t_c \equiv t_{\text{cool, physical}} / \Omega\):

\[
\Lambda = \frac{t_c^{-1} \Omega}{\gamma - 1} \left( \frac{p}{\rho} - \frac{c_{s,0}^2(x)}{\gamma} \right).
\]

(4)

The gravitational potential \(\Phi\) is given by Dong & Fung (2017)

\[
\Phi = -\frac{GM}{1 + q} \left[ \frac{1}{r} + \frac{q}{\sqrt{r^2 + R_p^2 - 2RR_p \cos \phi' + R_p^2}} \right]
\]

\[
-\frac{q R \cos \phi'}{R_p^2}
\]

(5)

where \(G\) is the gravitational constant, \(q\) the planet-star mass ratio, \(r\) the distance from the origin, \(R \equiv r \sin \theta\) the cylindrical radius, \(R_p\) the radius of the planet, and \(\phi' \equiv \phi - \phi_p\) the azimuthal separation from the planetary location. Because our simulations are in 3D, our smoothing length \(r_s = 0.01875 R_p\) is chosen solely to avoid singularity. The planet is fixed on a circular and coplanar orbit with \(R_p = 1\) in code units. A moderate inclination that keeps the planet within the disk \((i \lesssim h\), where \(h\) is the disk aspect ratio\) is unlikely to enhance spiral temperature structures as strongly as it does scattered light (e.g., Kloster & Flock 2019), although more substantial inclination would diminish the planet’s ability to transmit angular momentum to the disk and consequently reduce spiral visibility. We leave a detailed study on the effect of inclination to future work.

2.1. Background Temperature Profile

We obtain an axisymmetric background temperature profile, stratified in both radial and vertical directions, using the radiative-transfer code HOCHUNK3D (Whitney et al. 2013). Because radiative transfer is not scale-free, we set fiducial parameters loosely inspired by TW Hya; the planetary location at \(R_p = 1\) is scaled to 37 au (where one of the gaps in the system is located; see Tsukagoshi et al. 2016), and the disk mass set to 30MJ with a dust-to-gas ratio of 0.01 (all interstellar medium grains from Kim et al. 1994; well coupled to gas). We set the stellar radius to 2.09 \(R_p\) and temperature to 4000 K, typical of T Tauri stars. To ensure that the vertical distribution of disk material reflects the temperature profile, we enable the HSEQ mode of HOCHUNK3D (Whitney et al. 2013), which iterates the vertical disk profile after each RT iteration until convergence is reached.

We fit the resulting HOCHUNK3D temperature as described in Appendix A, and plot the result in Figure 1. From here, we find the sound speed as \(c_{s,0}(x) = \gamma k_B T(x) / \mu m_H\), taking the mean molecular weight \(\mu = 2.34\) as in the minimum-mass solar nebula. At the location of the planet, the effective aspect ratio \(h_{eq} \equiv c_{s,0}(x)/(\gamma^{1/2} R_p \Omega) \approx 0.07\). For a vertically isothermal disk in hydrostatic equilibrium, this would equal the true aspect ratio \(h\), but in a stratified disk we must define it explicitly as

\[
h(R) \equiv \left( \frac{1}{R^2 \Sigma(R)} \int_{-\infty}^{\infty} z^2 \rho(z) dz \right)^{1/2}.
\]

(6)

2.2. Disk Setup and Tests

We initialize the disk with a surface density

\[
\Sigma = \Sigma_0 \left( \frac{R}{R_p} \right)^{-1}.
\]

(7)
Figure 1. The stratified background temperature of our disk (no planet) obtained from a radiative transfer calculation assuming hydrostatic equilibrium, clearly rising with increasing distance above the midplane. The aspect ratio at a given cylindrical radius, $h(R)$, is plotted as a solid black line. We indicate the location of the planet $R_p = 1 = 37$ au with a vertical dotted line.

Our science simulations use a resolution of $267(r) \times 683(\phi) \times 42(\theta)$, spanning a radial range of $r = \{0.4, 5\}$, a polar range of $\theta = \{\pi/2, \pi/2 - 0.42\}$ (covering $\sim 6$ scale heights from the midplane at the planet location), and the full $2\pi$ in azimuth. Cells are spaced logarithmically in the radial direction, but uniformly in the polar and azimuthal directions. This yields a resolution of roughly seven cells per effective scale height at the location of the planet. We use periodic boundary conditions in the azimuthal direction, and reflective boundary conditions in the polar direction; in the radial direction, we use outflow boundaries because appropriate fixed boundary conditions for density, velocity, etc for our stratified temperature structure lack an analytic expression.

We verify our implementation of beta cooling with tests at $t_c = 10^{-4}$ and $10^{-6}$; these demonstrate typical fractional azimuthal temperature perturbations of order $\sim 10^{-5}$ and $\sim 10^{-14}$, respectively (not shown). In the vertical average, however, we see fluctuations in temperature deviating from the background at percent-level. This is because, even though temperature in any given grid cell deviates little from its background value, mass is redistributed upward and downward in the disk by the spiral density wave, changing the column-averaged temperature. This effect would not be captured in 2D simulations, so we cover it in more detail in the following section.

In the opposite limit of long cooling times, it is not immediately clear whether our simulations would reach any sort of steady state. To test whether they do, we run $t_c = 10^4$ models for each planet mass, and find that they yield essentially identical temperature and density structures to our $t_c = 10^{-2}$ runs. As an additional test, we run a $M_p = 50M_{\oplus}$, $t_c = 10^4$ simulation out to 400 orbits, and find that once spiral morphology and amplitude is established at $t \approx 10$ orbits, it remains in steady state (up to the gap opening) throughout. This is because planet-induced spiral arms are patterns that gas enters (to be compressed and heated) and leaves (to expand and cool) over the course of a single orbit, rather than persistent accumulations of gas.

Throughout this work, we define the relative strength along the spine of a Lindblad spiral at a given cylindrical radius ($\Delta T/T_{\text{spiral}}$) as:

$$\frac{\Delta T}{T_{\text{spiral}}}(R) = \frac{1}{2h_{\text{eff,p}}} \times \int_{\phi_{\text{peak}}(R) - 2h_{\text{eff,p}}}^{\phi_{\text{peak}}(R) + 2h_{\text{eff,p}}} d\phi \left( \frac{\Delta T}{\langle T \rangle}(R, \phi) \right)$$

where $\phi_{\text{peak}}(R)$ is the azimuthal location of the spiral density peak and $h_{\text{eff,planet}} = 0.07$ is the effective aspect ratio at the location of the planet, and $(\Delta T/\langle T \rangle)(R, \phi)$ is the mass-weighted, fractional azimuthal perturbation in vertically averaged temperature at a given 2D location:

$$\left( \frac{\Delta T}{\langle T \rangle}(R, \phi) \right) = \frac{1}{\Sigma(R, \phi)} \int_0^{\infty} d\rho \rho T = 1.$$
maps for one representative simulation—the $M_p = 50M_⊕$, $t_c = 10^2$ run used in our resolution test—in Figure 3.

3. Results

3.1. Lindblad Spirals

Lindblad spirals are formed by constructive interference of the Fourier modes excited in the disk by the planetary potential. They have a characteristic width of $\sim hR$ and a pattern speed of $\Omega(R_p)$, with waves launched at roughly $R_p \pm hR_p$, so one can estimate in the rotating frame that it takes gas parcels at most one dynamical time $t_{\text{dyn}} = \Omega^{-1}(R_p)$ to cross through. For the shortest cooling time in our model grid ($t_c = 10^{-4}$), therefore, $PdV$ work by the Lindblad spiral on a gas parcel is dissipated much faster than it is performed, so compression and expansion are nearly isothermal processes. For $t_c \gtrsim 1$, $PdV$ work is retained through an arm crossing, making the compression and expansion effectively adiabatic. We emphasize that, even when the simulation is in “steady state,” only the spiral pattern is fixed; the underlying gas is not in vertical hydrostatic equilibrium (Figure 3 of Dong & Fung 2017), as gas in the inner spiral arms has significant vertical motion (Zhu et al. 2015). In Figure 4, we plot $(\Delta T/T)_{\text{spiral}}$ at $r = 1.5$ in our Lindblad spirals for all 12 simulations in our model grid.

For $t_c = 10^{-2}$, therefore, the temperature in any given grid cell remains essentially unchanged from the prescribed background value; any nonaxisymmetry in temperature arises because kinematic effects redistribute material vertically. In the region of the Lindblad arm where a gas parcel is being compressed azimuthally—in anticipation of the spiral density peak—it becomes vertically inflated and thus heats up in vertical average. Following the density peak, gas expands azimuthally and shrinks to the midplane, cooling down in vertical average. This effect is stronger for inner spiral arms (0.3%–3%) than for the outer arms (0.1%–0.5%) whose amplitudes appear in Figure 5. Results from our test runs at $t_c = 10^{-4}$, not shown here, show quantitatively similar results.
By contrast, when $t_c > 1$, the $PdV$ work done by a Lindblad spiral is no longer dissipated before a gas parcel fully crosses through the arm. As a result, the high-density central spine of the arm becomes hotter adiabatically, while the lower-density regions before and after expand and cool. With the rising cooling time, this effect grows to dominate the overall temperature perturbation, causing it to follow the density (rather than its rate of change in azimuth, as in $t_c = 10^{-2}$ runs). As shown in the lower panel of Figure 5, this picture holds at all altitudes in the disk, although increasing sound speed and distance from the planet widen and somewhat weaken the spiral perturbation.

Motivated by these considerations, we present the following fit for the vertically averaged spiral temperature amplitudes as a function of both $M_p$ and $t_c$, plotted in the upper panel of Figure 4 for fiducial radius $R/R_p = 1.5$:

$$\frac{(\Delta T/T)_{\text{spiral,fit}}}{M_p} = \left(\frac{M_p}{M_{\text{th}}}ight)^{1.5} \left[ f_{\text{PdV}} e^{-t_{\text{arm}}/t_c} + f_{\text{rad}} \right]$$

(11)

where $c = 0.6$ is the power-law exponent of the temperature curves, $f_{\text{PdV}} = 0.06$ is the temperature perturbation arising from $PdV$ work in the adiabatic limit, $f_{\text{rad}} = 0.002$ that from vertical material redistribution in the stratified temperature structure, and $t_{\text{arm}} = 0.25$ a characteristic timescale for compression and expansion of material in the spiral arms. Differences between fitted and simulated amplitudes are typically $\lesssim 0.05 \times (\Delta T/T)_{\text{spiral,fit}}$. We note that a more extensive parameter survey in the future may help further refine this.

At any given planet mass, the temperature perturbation at any given planet mass is weakest at $t_c = 10^{-2}$ and strongest at $t_c = 10^2$. As visible in our Figure 6, however, density perturbation is nonmonotonic in regions at least several scale heights from the planet—strongest for $t_c = 10^{-2}$, but weakening at $t_c = 1$ before regaining some strength at $t_c = 10^2$ (Ziampras et al. 2020; Zhang & Zhu 2020). The analytical study of Miranda & Rafikov (2020) investigates this in detail, finding (in the linear limit) a radial separation between wavelocets set by

$$\text{Re}(k_m) \approx \frac{\Omega^2 - \bar{\omega}^2}{c_s}$$

(12)

and an amplitude damping rate

$$\text{Im}(k_m) = \frac{(\gamma - 1)\Omega \bar{\omega} t_c}{2(\Omega^2 + \bar{\omega}^2 t_c^2)} \text{Re}(k_m)$$

(13)

where $\bar{\omega} \equiv \pm m(\Omega - \Omega_p)$ is the Doppler-shifted forcing frequency of the planet, times the azimuthal wavenumber $m$. 

Physically, Im(\(k_m\)) represents the energy lost as beta cooling erodes the temperature component of the wave over the course of each oscillation. At all Lindblad resonances \(\Omega_m = \Omega_p(1 \mp 1/m)\), these losses are maximized for a \(t_{\text{crit}} \approx \gamma^{-1/2} \approx 1\), meaning waves are suppressed at their launching points. Cooling times longer and shorter than this, however, allow Lindblad waves to propagate more freely through the disk.

In Appendix B, we present vertically averaged galleries of 2D \(\Delta T/T\) (Figure 11) and \(\Delta \Sigma/\Sigma\) maps (Figure 12), as well as \((r, \theta)\) cuts of both temperature and density normalized to the azimuthal average (Figure 13). These provide a visual and intuitive understanding of the spiral perturbation for all simulations in our model grid.

3.2. Buoyancy-resonance Spirals

Buoyancy waves become apparent in temperature plots for \(t_c \gtrsim 1\) and saturate in strength at \(t_c = 10^2\). In each simulation, the vertically averaged temperature structure reveals at least three sets of buoyancy spirals, corresponding to different azimuthal tidal forcing wavenumber \(m\) (Zhu et al. 2012, 2015). Whereas for low planet masses and short cooling times these perturbations are weak (~1%–2.5%), they become stronger and substantially distorted in the opposite limit, wrapping around the full \(2\pi\) in azimuth and intersecting the Lindblad-resonance arm. In all cases with \(t_c \gtrsim 1\), we find buoyancy perturbation strengths several orders of magnitude stronger than in Zhu et al. (2015), who did not use a stratified background temperature.

The Lindblad spiral is strongest close to the midplane, widening and weakening in the disk atmosphere where sound speed is higher. Buoyancy spirals, however, are weak near the midplane but stronger in the disk atmosphere, the region typically probed by gas tracers such as \(^{12}\)CO. This is made quantitatively clear in the lower panel of Figure 5. Buoyancy causes hot material to rise and expand while pushing cold material to fall and contract, leaving pressure constant (at fixed altitude); consequently, in the long-\(t_c\) limit the ratio between temperature and density perturbations is \(\approx -1\) for buoyancy spirals, whereas for Lindblad spirals (an adiabatic phenomenon) the ratio is \(\approx (\gamma - 1) = 0.4\).

Previous work (e.g., Zhu et al. 2015; McNally et al. 2020; Bae et al. 2021) has established that buoyancy resonances occur where the Doppler-shifted forcing frequency of the planet commensurates with the local Brunt–Väisälä frequency:

\[
\pm m(\Omega - \Omega_p) = N
\]

(14)

where in a protoplanetary disk with \(g = \Omega^2z\),

\[
N^2 = \Omega^2 \left[ \frac{z}{T} \frac{\partial T}{\partial z} - \frac{1}{\gamma} \frac{z}{P} \frac{\partial P}{\partial z} \right].
\]

(15)

Because our disks are vertically stratified, the first (temperature) term is always present, resulting in very weak buoyancy arms close to corotation even with a \(t_c = 10^{-2}\). For longer cooling times (equivalently, the effective \(\gamma\) rising from 1 to its notional value of 1.4), the second (pressure) term makes resonances diagonal in the upper disk atmosphere, while the first creates an additional bend at the interface between disk...
midplane and atmosphere, where temperature is changing rapidly. For superthermal planets, buoyant temperature perturbations are sufficiently strong as to change $N$ and alter the resonance location, creating coupling between different buoyancy modes that manifests as merging and splitting of buoyant arms.

We stress that while the Brunt–Väisälä frequency is well behaved as $\gamma \to 1$ from above, it is undefined for a truly isothermal equation of state, for which $\gamma = 1$ and the Navier–Stokes energy equation is degenerate. Thus there are no buoyant perturbations in simulations like Juhász & Rosotti (2016), even though the background temperature is vertically stratified.

3.3. Observational Diagnostics

In Figure 7, we plot temperature cuts in the disk at $\theta = 1, 2$, and $3H$; this is intended to qualitatively mimic the different layers of the disk probed by different gas tracers. As before, we normalize and subtract away the average in each azimuthal ring; however, we extend the radial range of the plots to $0.5, 5$ to emphasize that $(\Delta T/[T])_{\text{spir}}$ is only $0.025$ within a factor of $2\sim3$ orbital radii from the perturber. We present results from the $t_c = 1, 50 M_\odot$ and $400 M_\odot$ models as representative cases. In Figure 8 we compare our simulated results (scaled to physical units) with those of Teague et al. (2019).

For both planet masses, the primary Lindblad spiral widens, and a secondary spiral emerges (Bae & Zhu 2018a, 2018b), with increasing altitude in the disk. Buoyancy spirals likewise grow in strength and number with increasing altitude. Both Lindblad spirals wind more tightly with increasing distance from the planet, as expected from Equation (12)—which, in the limit $r \gg r_p$, reduces to $\text{Re}(k_n) \approx (m/\gamma h_{\text{eq}} R)(\Omega_p/\Omega)$. Convolved with a beam, these two arms emanating from a single point may resemble observations of the CQ Tau (Wöfler et al. 2020). On the other hand, buoyancy spirals are tightly wound even close to the planet. We find, in line with previous isothermal simulations (e.g., Fung et al. 2014; Fung & Chiang 2016), that planet-carved gaps are circular in the midplane, and widen with increasing altitude; but, in a departure from previous work, at higher altitudes our nonisothermal gaps exhibit a substantial “break” at the planet location.

For a better observational understanding, we postprocess our hydrodynamical simulations in radiative transfer with HOCHUNK3D, using $10^9$ photons to obtain background temperatures and scattered-light images at face on for each of our $t_c = 1$ runs. In these simulations, $Pd\nu$ contributions from the hydrodynamics are obviated, and only explicit irradiation effects on $\Delta T$ are included. We plot both in Figure 9.

For our temperature panels, we plot the $\theta = 0.21 \approx 3H$ layer studied by gas observations, as in Figure 7. Taking a cut in this fashion allows us to isolate the effects of spiral shadowing from those of vertical redistribution of material (see Section 3.1, paragraph 2, for a more detailed discussion). In the inner disk, we find azimuthal asymmetries that trace the primary and secondary Lindblad arms, ranging from $1\%$ for our $50 M_\odot$ case to $8\%$ for the $400 M_\odot$ case whose arms strongly perturb the disk surface. The outer spirals cause no shadowing effect on temperature; but, for superthermal planets, the reduction of disk scale height near the planet location exposes the outer disk at $\phi = \pi$ to increased stellar irradiation, giving the impression of a radial “arm” in temperature. We note that this effect is expected to be transient and disappear once a gap is opened.

While spiral shadowing can be noticeable at high altitudes, it has a markedly lower impact on the disk as a whole. In mass-weighted vertical average, the azimuthal temperature perturbation obtained with HOCHUNK3D is typically $<1\% \sim 3\%$ interior to the planet, but $<1\%$ in the outer disk. As these differences are consistent with those plotted in Figures 4 and 6 for our $t_c = 10^{-2}$ runs, we surmise that they arise primarily from
vertical redistribution of disk material in the preexisting stratified temperature structure, already accounted for in our hydrodynamics. In any case, $PdV$ work in Lindblad spirals (at high altitudes and large planet masses, buoyancy spirals as well) overwhelms any radiative effects.

In our near-infrared scattered-light images, the inner Lindblad spiral becomes prominent for the superthermal $M_p = 200, 400 M_\oplus$ runs, and the outer arms are less clearly observable. This aligns with expectations that Lindblad spirals ought to be visible only when planets substantially alter the disk scattering surface, and with the simulations of Dong et al. (2016), who test substantially higher thermal masses (albeit in isothermal disks with temperature constant along cylinders). In scattered light, the inner primary and secondary Lindblad arms are nearly as strong as each other, but the outer secondary Lindblad and buoyancy arms are barely visible. As in the temperature, there is a radial pseudo “arm” in scattered light for superthermal planets that reshape disk structure around them.

4. Conclusions

Using 3D hydrodynamic simulations with an adiabatic equation of state and beta cooling as well as radiative transfer pre and postprocessing, we quantitatively investigate the spiral temperature perturbations driven by planets in their natal disks, and provide an empirical fit (Equation (11)). Our simulations employ a stratified background temperature, as appropriate for a passively heated disk in hydrostatic equilibrium. In the short-$t_c$ (dimensionless cooling timescale) limit, equivalent to isothermal equation of state, we observe azimuthal temperature perturbations of <0.5% for exterior Lindblad spirals once vertically averaged (0.3%–3% for interior, depending on planet mass). This effect is mainly caused by a vertical redistribution of material in a stratified background. Azimuthal temperature perturbations in our hydrodynamical simulations, both vertically averaged and as a function of height above the midplane, can be found in Figures 11–13 in our Appendix.
Typical protoplanetary disks are optically thick. We find, with radiative-transfer postprocessing, that a secondary irradiation effect—shadowing from the spirals themselves, on top of the stratified background—has only a minor impact on vertically averaged disk temperature. However, in the tenuous upper layers of the disk probed by $^{12}$CO mapping ($\rho(\theta = 0.21 \approx 3H/\rho_{\text{mid}} \approx 0.01$), shadowing has an effect on inner spirals—from 1% at 50 $M_\odot$ to a potentially observable 8% at 400 $M_\odot$, assuming $t_c = 1$—but remains unimportant for outer spirals.

For longer cooling times, the radiative effects remain, but are overwhelmed by local $PdV$ heating, which produces a vertically averaged temperature perturbation ranging from 3%–11% for $t_c = 1$ and 4%–14% for $t_c = 100$, as mass increases from 50–400 $M_\odot$. Lindblad spirals are strongest in the midplane, whereas buoyancy spirals are strongest in the low-density disk atmosphere, and so are somewhat weaker in vertical average. Our superthermal, long-cooling models show azimuthal temperature perturbations comparable to those observed in $^{12}$CO in TW Hydrae (Teague et al. 2019) and CQ Tauri (Wölfer et al. 2020), and provide a starting point for simulations with a more realistic, radius-dependent cooling time (e.g., Miranda & Rafikov 2020).

Our investigations reveal that gas temperature and observations tracing material distributions complement each other as signatures of planet-driven spirals. Hydrodynamical scale-height perturbations are stronger for inner Lindblad spirals (e.g., Fung et al. 2015; Dong & Fung 2017) and for short cooling times (Ziampras et al. 2020), enhancing their visibility in scattered light. By contrast, outer Lindblad spirals are geometrically larger, and spirals with longer cooling times experience $PdV$ heating, making them more prominent in temperature. Spiral temperature structure typically persists for 2–3 windings inside and outside the planet’s orbital radius before fading to an azimuthal contrast of $\sim 2.5\%$, even for superthermal mass planets. This, conversely, allows us to use observed arms to constrain the location of the perturber.

Buoyancy waves, an inherently 3D phenomenon requiring a nonisothermal equation of state, have historically been found to be weak (Zhu et al. 2012; Lubow & Zhu 2014; Zhu et al. 2015) in simulations where temperature is constant along cylinders. We find that a realistic, stratified temperature structure amplifies the effect of buoyancy (see also Bae et al. 2021, whose greater vertical temperature gradient leads to even stronger buoyancy resonances), which fundamentally is a process that causes hot material to rise and expand and cold material to sink and contract. Relatively unimportant near the midplane, buoyancy spirals strengthen in the hot disk atmosphere, becoming comparable in $\Delta T/T$ to the more extensively studied Lindblad spirals despite a weaker density perturbation.

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Appendix A
Temperature Fit

We fit the radially and vertically stratified temperature profile from HOCHUNK3D for a passively heated disk in hydrostatic equilibrium (see Section 2.1) as the following:

$$T_{\text{fit}}(r, \theta) = (T_{\text{fit}}(r) - T_{\text{mid}}(r))(1 - f_2(r, \theta) - f_4(r, \theta)) + T_{\text{mid}}(r)$$  \hspace{1cm} (A1)

where

$$T_{\text{mid}}(r) = T_{\text{mid},0} \left( \frac{r}{R_p} \right)^{\beta_{\text{mid}}}$$  \hspace{1cm} (A2)

is the midplane temperature,

$$T_{\text{atm}}(r) = T_{\text{atm},0} \left( \frac{r}{R_p} \right)^{\beta_{\text{atm}}}$$  \hspace{1cm} (A3)

is the temperature in the optically thin upper part of the disk atmosphere, and

$$f_2 = (1 - f_{4,0}(r))\exp(-\theta^2/\theta_2^2(r))$$  \hspace{1cm} (A4)

$$f_4 = f_{4,0}(r)\exp(-\theta^2/\theta_4^2(r))$$  \hspace{1cm} (A5)

are interpolations between the two. We obtain that $T_{\text{mid},0} = 31.522K$, $\beta_{\text{mid}} = -0.384$, $T_{\text{atm},0} = 106.655K$, and $\beta_{\text{atm}} = -0.422$. We fit $f_{4,0}$, $\theta_2$, and $\theta_4$ as second-order polynomials in the logarithm of radial position $r$,

$$g(r) = a_0 + a_1 \log(r/R_p)$$

For $\theta_2$, $a_0 = 0.13127$, $a_1 = -0.28745$, and $a_2 = 1.63323$; for $\theta_4$ we have $a_0 = 0.30463$, $a_1 = 0.16974$, $a_2 = -0.06711$; and, for $f_{4,0}$, $a_0 = 0.91957$, $a_1 = -0.079279$, $a_2 = 0.20998$. In Figure 10, we plot the relative error $\Delta T/T_{\text{fit}}(r, \theta)$, where in this case

$$\Delta T/T_{\text{fit}}(r, \theta) = \frac{1}{2\pi} \int_0^{2\pi} d\phi \left( \frac{T_{\text{actual}}(r, \theta, \phi)}{T_{\text{fit}}(r, \theta)} - 1 \right)$$  \hspace{1cm} (A6)
Appendix B
Temperature and Density Plots

Figure 11. Azimuthal perturbation in vertically averaged temperature obtained from PENGUIN hydrodynamics simulations of our model grid. Perturbation amplitude, for both the Lindblad and buoyancy spirals, increase with planetary mass and cooling times. The vertical line at $r = 1.5$ in our $M_p = 50M_\oplus$, $t_c = 10^2$ run indicates the cut we take for our temperature plot in the upper panel of Figure 5. All runs were taken at 25 orbits.
Figure 12. Azimuthal perturbation of surface density in our simulations. These are clearly strongest in the isothermal limit, weakening somewhat with increasing cooling time at a given planet mass. The straight line at the planet radius counterbalances concentration of material at the planet location, which is strongest for high planet masses and short cooling times (as in, e.g., Fung et al. 2019). The vertical line at $r = 1.5$ in our $M_p = 50M_⊕$, $t_c = 10^2$ run indicates the cut we take for our temperature plot in the upper panel of Figure 5.
Figure 13. Plot of the temperature perturbation $\Delta T_{\text{cut}}$ (above) and density perturbation $\Delta \rho_{\text{cut}}$ (below) in a cut at $\phi = -\pi/2$ relative to the planet, relative to the azimuthal averages $\langle T \rangle_{\text{ring}}$ and $\langle \rho \rangle_{\text{ring}}$ at each $(r, \theta)$. Spherical radius $r$ is on the x-axis and polar angle $\theta$ on y. Lindblad spirals bend outward because the vertical increase of temperature shifts the effective resonance location outward (Artymowicz 1993). Buoyancy resonances are also distorted, especially for superthermal planets that themselves meaningfully alter the disk temperature structure.
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