X-Ray Ionization of Planet-Opened Gaps in Protostellar Disks

S. Y. Kim\textsuperscript{1,2}, N. J. Turner\textsuperscript{3}

\textbf{ABSTRACT}

Young planets with masses approaching Jupiter’s have tides strong enough to clear gaps around their orbits in the protostellar disk. Gas flow through the gaps regulates the planets’ further growth and governs the disks’ evolution. Magnetic forces may drive that flow if the gas is sufficiently ionized to couple to the fields. We compute the ionizing effects of the X-rays from the central young star, using Monte Carlo radiative transfer calculations to find the spectrum of Compton-scattered photons reaching the planet’s vicinity. The scattered X-rays ionize the gas at rates similar to or greater than the interstellar cosmic ray rate near planets the mass of Saturn and of Jupiter, located at 5 au and at 10 au, in disks with the interstellar mass fraction of sub-micron dust and with the dust depleted a factor 100. Solving a simplified gas-grain recombination reaction network yields charged particle populations whose ability to carry currents is sufficient to partly couple the magnetic fields to the gas around the planet. However the material near the planet’s orbit has ambipolar and/or Hall diffusivity so large in all the cases we examine that the non-ideal terms dominate the magnetic field’s evolution. Thus the flow of gas in the gaps opened by the giant planets depends critically on the finite conductivity and its spatial variation.

\textit{Subject headings:} protoplanetary disks — X-rays — radiative transfer — magnetic fields

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1. INTRODUCTION

A protoplanet that grows past the mass of Saturn has gravity strong enough for its tides to clear an annular gap of low surface density around its orbit in the protostellar disk, as reviewed by Baruteau et al. (2014). The planet can continue to grow only if material from the gap walls is able to reach its Hill sphere. Furthermore, the planet’s subsequent orbital evolution is governed by the distribution of gas across the gap, especially the amount of material located near orbital resonances. The orbit’s evolution also depends on the rate at which disk material crosses the gap. In this contribution we investigate whether magnetic forces can act on the gap. We explore whether the gas in the gap is ionized well enough to couple to magnetic fields, so that the fields can displace the material near the planet.

Young stars emit ionizing X-rays with temperatures of thousands of electron volts (Feigelson & Montmerle 1999), able to penetrate the circumstellar gas and dust to columns of order $10 \text{ g cm}^{-2}$ (Glassgold et al. 1997; Ercolano & Glassgold 2013). By comparison, the minimum-mass Solar nebula has a surface density at 5 au of about $150 \text{ g cm}^{-3}$ (Weiden- schilling 1977; Hayashi 1981). Few stellar X-ray photons thus reach the midplane before planets have grown. However once a gap opens in the disk, the X-ray flux at the midplane can increase. The gas making up the gap’s inner rim will forward-scatter some of the photons arriving from the star, deflecting a fraction down to the planet’s vicinity. The gap’s outer rim, exposed directly to starlight, heats and expands vertically (Turner et al. 2012), intercepting extra X-ray photons, some of which will be scattered backward and down into the gap. A planet opening a gap thus receives more X-rays than a non-gap-opening planet at the same location.

Here we examine whether a gap can increase the X-ray flux enough to contribute significantly to the ionization of gas near the planet. We compute the X-ray intensity in the planet’s vicinity using a Monte Carlo radiative transfer approach, estimate the resulting ionization state by integrating a simple ionization-recombination reaction network to equilibrium, and compute the plasma’s magnetic diffusivity. We compare against the threshold diffusivities required for the operation of three mechanisms proposed to drive the accretion flow: magneto-rotational turbulence, Hall-shear instability, and magneto-centrifugal winds. We extend recent work by Keith & Wardle (2015) in carrying out X-ray transfer calculations rather than using results from a gapless disk, while making relatively simple assumptions about the gap structure and the strength and geometry of the magnetic fields.

The paper is laid out as follows. The model star and disk are described in §2, the X-ray transfer methods in §3. The ionization-recombination chemistry and how we translate the charged species’ abundances into diffusivities and magnetic stresses are set out in §4. The resulting X-ray spectra, ionization levels, and magnetic coupling are shown in §5. Discussion
2. STAR AND DISK

The young star is of Solar mass and twice Solar radius, and emits a blackbody spectrum with effective temperature 4500 K. The resulting luminosity of $L_*=1.5$ times Solar is in the range indicated by stellar evolution modeling at ages 1–2 Myr (D’Antona & Mazzitelli 1994; Siess et al. 2000). The star is surrounded by an axisymmetric disk laid down on a radiative transfer grid logarithmically-spaced in radius with 120 cells from $10^{-1.4}$ to $10^{1.6}$ au (approximately 0.04 to 40 au). The vertical structure is resolved by dividing each annulus into 60 cells spaced uniformly from the midplane up to 6 initial pressure scale heights.

We adopt disk surface density profiles from a one-dimensional, semi-analytic model constructed by Lubow & D’Angelo (2006), in which the planet’s gap-opening tidal torques balance the disk’s gap-closing viscous stresses (Lin & Papaloizou 1986). The planet follows a fixed circular orbit and the gap is treated separately from the rest of the disk. The torques exerted by the planet on the disk are assumed confined to the gap, so that planet-generated waves propagating away into the rest of the disk are ignored. The accretion stresses within the disk are modeled using the Shakura-Sunyaev viscous prescription with $\alpha = 0.005$, and the flow is assumed to be steady-state.

Before any planet is added, the disk’s surface density varies inversely with radius, and is 280 g cm$^{-2}$ at 1 au. The planet partly dams the flow and accretes much of the material reaching its orbit. The ratio of the planet’s accretion rate to the flow rate at the same place in the planet-free disk is $E = 6$. This is lower than the $E = 8$ fiducial case of Lubow &
D’Angelo (2006) because our models have lower aspect ratios $H/r \approx 0.04$ near 5 and 10 au. Other parameters take the fiducial values Lubow & D’Angelo (2006) obtained by matching their semi-analytic solution to results from numerical hydrodynamical calculations.

We place the planet at either $r_p = 5$ or 10 au, approximating the orbits of Jupiter and Saturn, respectively. To test how magnetic forces can be expected to vary as the planet grows, at each location we consider bodies of both Saturn and Jupiter masses, using planet-to-star mass ratios $q = 3 \times 10^{-4}$ and $10^{-3}$, respectively. As a first guess at the midplane temperature profile we adopt $T(r) = 124(L_*/L_\odot)^{1/4}(r/\text{au})^{-1/2}$ K, based on similar radiative transfer calculations without a gap Turner et al. (2012). This is cooler than the minimum-mass Solar nebula because it includes the effects of the disk’s large optical depth. The resulting surface density profiles are shown in figure 1. The mass of material on the grid is 0.18 $M_\odot$ with the planet at 5 au, and 0.08 $M_\odot$ with the planet at 10 au, whether the planet has the mass of Jupiter or Saturn. These disk masses are a few to ten times more than the minimum-mass Solar nebula. We verified that the model disks are not susceptible to fragmentation under their own self-gravity, having the Toomre $Q$ parameter $c_s\Omega/(\pi G \Sigma)$ greater than unity everywhere. Here $c_s$ is the sound speed, $\Omega$ the orbital frequency, $G$ the gravitational constant, and $\Sigma$ the surface mass density. The mass accretion rate in all models is within 20% of $3 \times 10^{-8} M_\odot \text{yr}^{-1}$ outside the planet’s orbit, and $5 \times 10^{-9} M_\odot \text{yr}^{-1}$ inside.

The gap modifies the disk’s temperature profile from our initial guess. Evacuating the gap allows starlight to directly strike the top of the gap’s outer wall. The starlight heats the wall, increasing its internal gas pressure, so in hydrostatic equilibrium the wall becomes taller and intercepts yet more starlight. The tall wall is likely to intercept more of the stellar X-rays too. We therefore include these effects, using an iterative procedure similar to Turner et al. (2012). Given a density distribution, we obtain new temperatures under radiative balance with the starlight using Monte Carlo transfer with the Bjorkman & Wood (2001) relaxation method. We then displace gas up or down to restore vertical hydrostatic balance, holding fixed the variation of temperature with column. The new density distribution serves as the input for the next iteration. We quit after five iterations when the structure no longer changes significantly.

The disk’s opacity to the starlight and reprocessed infrared radiation comes from dust grains. We take dust opacities from Preibisch et al. (1993), where the particle size distribution

$$n(a) \propto a^{-p}, \quad a_{\text{min}} < a < a_{\text{max}},$$

with $p = 3.5$. At temperatures below 125 K, the grains are composed of a silicate core and an icy mantle whose radius is 14% of the core’s. The mantle is polluted with tiny amorphous carbon grains ($a_{\text{min}} = 0.007 \mu\text{m}, a_{\text{max}} = 0.03 \mu\text{m}$). Above 125 K, the icy mantle sublimes, baring a silicate grain ($a_{\text{min}} = 0.04 \mu\text{m}, a_{\text{max}} = 1 \mu\text{m}$) and releasing the amorphous carbon
grains; the silicates sublimate at 1500 K, the carbon grains at 2000 K. The grains are well-mixed in the gas, and we assume the scattering that makes up part of the starlight opacity is isotropic. To model small grains’ incorporation into larger bodies, which is likely to have occurred by the time planets grow to Saturn or Jupiter mass, we also compute models with the dust opacities multiplied by a factor $\epsilon = 10^{-2}$. In the dust-depleted disks, the starlight first reaches unit optical depth nearer to the midplane, so less of the stellar luminosity is intercepted, the midplane temperature is lower, and hydrostatic equilibrium requires that the density scale height is less. Considering two values each for the planet location, the planet mass, and the dust abundance, we compute the eight model disks listed in table 1. Each model is given a name whose first digits are the planet location in au, followed by a letter indicating the planet has Saturn (S) or Jupiter (J) mass, and a final digit that is the logarithm of the factor by which the dust abundance is reduced.

After carrying out the transfer calculations, we noticed that the very low surface densities near the four Jupiter-mass planets imply radial flow speeds $v_r = \dot{M}/(2\pi r \Sigma)$ exceeding the sound speed, which seems unlikely given that the models are otherwise near hydrostatic equilibrium, with flows driven in part by gas pressure. For the coupling calculations described below in section 4.3, we therefore increase the surface density uniformly across the gap so the flow is subsonic everywhere. Despite increasing by factors $5 \times 10^4$ at 10 au and $4 \times 10^8$ at 5 au over the values listed in table 1, the surface density remains low enough that the gap is quite optically-thin to all three relevant wavebands — the visible starlight, the disk’s own infrared emission, and the X-rays. Therefore redoing the transfer calculations would yield basically identical results.

<table>
<thead>
<tr>
<th>Model</th>
<th>Radius $r_p$, au</th>
<th>Mass $q$</th>
<th>Dust Depletion $\epsilon$</th>
<th>Column $\zeta(r_p, z=0)$, s$^{-1}$</th>
<th>Ratio to IG99 fit</th>
</tr>
</thead>
<tbody>
<tr>
<td>10S0</td>
<td>10</td>
<td>$3 \times 10^{-4}$</td>
<td>1</td>
<td>$5.7 \times 10^{-19}$</td>
<td>0.034</td>
</tr>
<tr>
<td>10S2</td>
<td>10</td>
<td>$3 \times 10^{-4}$</td>
<td>0.01</td>
<td>$4.9 \times 10^{-18}$</td>
<td>0.29</td>
</tr>
<tr>
<td>10J0</td>
<td>10</td>
<td>$10^{-3}$</td>
<td>1</td>
<td>$3.1 \times 10^{-16}$</td>
<td>6.0</td>
</tr>
<tr>
<td>10J2</td>
<td>10</td>
<td>$10^{-3}$</td>
<td>0.01</td>
<td>$3.6 \times 10^{-16}$</td>
<td>6.9</td>
</tr>
<tr>
<td>5S0</td>
<td>5</td>
<td>$3 \times 10^{-4}$</td>
<td>1</td>
<td>$8.0 \times 10^{-18}$</td>
<td>0.10</td>
</tr>
<tr>
<td>5S2</td>
<td>5</td>
<td>$3 \times 10^{-4}$</td>
<td>0.01</td>
<td>$3.5 \times 10^{-17}$</td>
<td>0.46</td>
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<td>$10^{-3}$</td>
<td>1</td>
<td>$2 \times 10^{-10}$</td>
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</tr>
<tr>
<td>5J2</td>
<td>5</td>
<td>$10^{-3}$</td>
<td>0.01</td>
<td>$2 \times 10^{-10}$</td>
<td>6.7</td>
</tr>
</tbody>
</table>
3. X-RAY IONIZATION

We compute X-ray ionization rates in each model disk by sampling the X-ray mean intensity using Monte Carlo techniques, and converting the absorbed X-ray energies into ionization rates. In most respects we follow Igea & Glassgold (1999, hereafter IG99). In the sections below we sketch the IG99 approach and note our points of departure.

3.1. X-Ray Source

Magnetic reconnection heats plasma in the young star’s corona to temperatures exceeding \( kT_X = 1 \) keV. The plasma emits X-rays by thermal bremsstrahlung, whose spectral luminosity at energy \( E \) we approximate by

\[
L(E) = \frac{L_X}{kT_X} \exp \left( -\frac{E}{kT_X} \right),
\]

where \( L_X \) is the total X-ray luminosity and \( T_X \) is the source temperature. A less approximate treatment would involve quantum electrodynamic corrections. An electron-velocity-distribution-averaged Gaunt factor that has been set to unity in deriving eq. 2 would then decline by a factor four across the energy range from 1 to 30 keV, tilting the spectrum toward low energies (Karzas & Latter 1961). This would yield higher ionization rates in the disk’s surface layers, and lower rates in the interior. However, we follow IG99 in using eq. 2.

Also like IG99, we locate the X-ray source in a helmet streamer taking the shape of a ring centered on the rotation axis, whose radius and height above the disk midplane are each ten times the stellar radius. We adopt a source temperature \( T_X = 5 \) keV and an X-ray luminosity \( L_X = 2 \times 10^{30} \) erg s\(^{-1} \), the median for Solar-mass stars in the Orion Nebula Cluster (Garmire et al. 2000). We consider X-rays below 1 keV to be attenuated in the star’s magnetosphere and inner wind. We thus compute the transfer of X-rays with energies from 1 up to 30 keV.

3.2. X-Ray Opacities

IG99 used a simple power-law energy dependence for the X-ray absorption cross section, and carried out their Monte Carlo calculations assuming all heavy elements were segregated from the gas. We adopt a more detailed fit by Bethell & Bergin (2011a), who for solar elemental abundances find a cross section of the form

\[
\sigma_{\text{tot}} = \sigma_{\text{gas}} + \epsilon f_b(E) \sigma_{\text{dust}},
\]

where \( \sigma_{\text{gas}} \), \( \sigma_{\text{dust}} \), and \( \epsilon f_b(E) \) are the gas, dust, and electron-velocity-averaged Gaunt factor contributions, respectively.
where $\epsilon$ and $f_b$ are the dust settling and grain growth parameters, respectively, and the cross sections of the gas and dust are each given by

$$\sigma(E) = 10^{-24} \text{cm}^2 \times (c_0 + c_1 E + c_2 E^2) E^{-3}, \quad (4)$$

with different fitting coefficients $c_i$ for dust and gas. The coefficients are piecewise linear functions of energy, to reproduce discontinuous increases in the absorption opacity due to K-shell photoelectric absorption by various metals. Bethell & Bergin (2011a) provide fits up to 10 keV; we extrapolate to higher energies. Consistent with the optical and infrared transfer, we consider both $\epsilon = 1$ and $\epsilon = 0.01$ to allow for the depletion of small grains. The dust growth parameter $f_b$ is unity for all our calculations. Like IG99 we take the scattering cross section from the Klein-Nishina formula. Cross sections are converted to opacities using a mean gas molecular weight of 2.3, corresponding to Solar composition with the hydrogen in molecular form. The energy dependence of the opacities is shown in Figure 2. The scattering cross section is almost constant across the energy range shown, falling off very slightly toward the top end. The absorption falls off much more steeply with photon energy. Thus the albedo $\alpha = \sigma / (\kappa + \sigma)$, where $\kappa$ is the absorption and $\sigma$ is the scattering opacity, increases dramatically with energy. The highest-energy photons are more likely to be scattered than absorbed. This contributes to the harder X-rays penetrating deeper into the disk.

### 3.3. X-Ray Transfer and Ionization Rate

We sample the X-rays’ mean intensities on the radiative transfer grid using a Monte Carlo approach. We follow each packet of X-ray photons from its emission at the stellar source into the disk and through as many scatterings as needed till the packet either escapes the domain or is absorbed. We assume that all absorbed X-ray energy is converted into ionization. A scattered X-ray packet’s new direction is chosen randomly from the Klein-Nishina phase function and its energy is reduced by Compton losses.

X-ray sources with different energies produce mutually independent intensities in the disk. We therefore perform the Monte Carlo procedure separately for each monochromatic source energy. To construct the desired source spectrum we linearly combine the intensities produced by the monochromatic sources. The resulting mean X-ray intensity is converted to an ionization rate by dividing by the average energy required to produce an ion pair, for which we follow IG99 and adopt the value $\Delta E = 37$ eV.

We differ from IG99 in how we compute the mean intensity. They average over the projected area met by each packet at each radiative transfer grid cell. We instead follow
Fig. 2.— X-ray scattering and absorption opacities. The nearly-horizontal line denotes the scattering opacity derived from the Klein-Nishina formula. The absorption opacities, from top to bottom, are (1) IG99’s power-law fit assuming solar abundances, (2) Bethell & Bergin (2011a)’s more detailed fit with dust undepleted and (3) fully depleted, and (4) IG99’s power-law fit for solar abundance but with heavy elements depleted onto grains which have been removed by settling and growth. We use the Bethell & Bergin opacities.

As shown in figure 3 top panel, the source spectrum is adequately represented with 30 energy bins. Thus, for each of 30 monochromatic X-ray source energies, uniformly spaced from 1 to 30 keV, we send $10^6$ photon packets into our model disks, then construct a weighted sum of the monochromatic results to determine ionization rates for a source with a thermal bremsstrahlung spectrum at a temperature of 5 keV.
4. MAGNETIC COUPLING

4.1. Ionization State

We compute the ionization state by balancing the X-ray ionization with a recombination reaction network including grain surface reactions and simplified gas-phase chemistry (model 4 of Ilgner & Nelson 2006). This yields the equilibrium abundances of seven charged species: electrons, a representative molecular ion (HCO$^+$), a representative metal ion (Mg$^+$), and grains charged by one and two electrons either side of neutral.

The magnesium has abundance $3.39 \times 10^{-7}$ atoms per hydrogen nucleus, 1% of the Solar value, since most of the magnesium is locked up in minerals and only a minority can react on the grain surfaces or in the gas phase. The fraction of magnesium in available form has little impact on the ionization balance, since temperatures are low enough that most magnesium atoms stay adsorbed on the grains. The grains’ abundances in the reaction network match those in the corresponding starlight and X-ray radiative transfer calculations, in the following approximate sense. The reaction network treats grains of a single radius, which we set to 0.1 µm. When combined with a material density of 2 g cm$^{-3}$ this nearly matches the geometric cross-section per unit mass in the optical opacity model.

4.2. Magnetic Field Strengths

We evaluate three scenarios for whether magnetic torques can drive the accretion flow: (1) magneto-rotational turbulence (henceforth MRT), (2) the Hall-shear instability (HSI), and (3) a magneto-centrifugal wind launched from the disk surface. The MRT sustains a tangled field with strong azimuthal and moderate radial components, on a weaker vertical background field. Hall-shear instability couples the Hall term’s rotation of toroidal into radial field, together with the orbital shear’s stretching out the radial component to generate fresh toroidal field. The resulting fields have a dominant toroidal component. The magneto-centrifugal winds we consider have all three field components comparable near the disk surface.

For each scenario to produce the accretion rates present in the model disk, the field must reach a certain strength, as follows. Taking the equation of motion in cylindrical coordinates $(r, \phi, z)$, reducing to the case of a near-axisymmetric, near-Keplerian disk with time-steady mean internal flows, and averaging over the disk thickness $2h$, yields

$$\frac{\dot{M} \Omega}{2r} = \frac{2h}{r} \langle -B_r B_\phi \rangle + \frac{\partial}{\partial r} \left[ h \langle -B_r B_\phi \rangle \right] - (B_z B_\phi)_s,$$

(5)
where angle brackets denote averages through the disk, and the subscript \( s \) marks fields measured on the disk’s top and bottom surfaces (Wardle 2007). The first two terms on the right come from magnetic stresses in the disk interior, for example due to MRT or HSI, while the final term is the wind’s back-reaction on the disk via magnetic torques. As is common, we assume the mean stress varies only over length scales at least comparable to the radius, so that the second term is comparable to or smaller than the first, and can be neglected.

In the MRT scenario, angular momentum conservation thus links the mass flow rate \( \dot{M} \) carried by the disk to the magnetic accretion stress by

\[
\langle -B_r B_\phi \rangle = \frac{\dot{M} \Omega}{4h}. \tag{6}
\]

The stress from the field’s radial and azimuthal components is about one-quarter the squared magnitude of the magnetic field, which in turn is about 20 times the mean squared vertical field, based on direct numerical calculations (Hawley et al. 1995; Sano et al. 2004). We can thus solve eq. 6 for the RMS vertical field.

Eq. 6 also gives the stress in the HSI scenario. To find the mean vertical field strength, we observe that the toroidal component is about 50 and 200 times the radial and vertical components in Lesur et al. (2014, run 1-OHA-5).

In the wind scenario, eq. 5 becomes

\[
(-B_z B_\phi)_s = \frac{\dot{M} \Omega}{2r}, \tag{7}
\]

where the \( s \) subscript indicates values measured at the disk surface. We estimate the vertical field using the fact that magneto-centrifugal wind solutions typically have the three field components roughly equal at the disk surface. The surface connects the interior, with its vertical magnetic field, to the wind, where the field is angled away from the rotation axis.

4.3. Coupling Criteria

With the field strength in hand from eq. 6 or 7, and the charged species’ abundances from sec. 4.1, we can find the magnetic diffusivities and determine whether the field in fact couples to the gas well enough to drive the accretion flow.

We compute the diffusivities \( \eta_O \), \( \eta_H \) and \( \eta_A \) that are the coefficients of the Ohmic, Hall and ambipolar terms in the induction equation, including the contributions from all charged species, following Wardle (2007) eqs. 21-31. The diffusivities depend on the field strength,
which governs whether each charged species mostly gyrates around the field lines under the Lorentz force, or mostly random walks by colliding with neutrals.

Each of the three magnetic scenarios works only when the diffusivities and field strengths meet a set of requirements. For MRT, disturbances with the linear magneto-rotational instability’s fastest-growing wavelength must diffuse away more slowly than they grow. The wavelength $\nu_{Az}/\Omega$ depends on the Alfven speed $\nu_{Az}$ along the vertical magnetic field, the growth rate is close to $\Omega$, and the relevant diffusivity is the sum of the Ohmic and ambipolar values, called the Pedersen diffusivity $\eta_P = \eta_O + \eta_A$. Thus MRT requires

$$\eta_P < \nu_{Az}^2/\Omega \quad (8)$$

(Sano & Inutsuka 2001; Turner et al. 2007; Keith & Wardle 2015). In addition the wavelength must fit within the disk thickness, corresponding to a vertical magnetic field with pressure less than that of the gas by a factor

$$\beta_z > 8\pi^2 \quad (9)$$

(Okuzumi & Hirose 2011). At the same time, the Hall term must be small enough to not much modify the character of the turbulence. Lesur et al. (2014) find that the Hall term dominates if the Hall length exceeds 20% of the gas scale height. The Hall length is $|\eta_H|/\nu_A$, thus standard MRT requires approximately

$$|\eta_H| < \beta_z^{1/2}\nu_{Az}^2/\Omega \quad (10)$$

For the Hall-shear instability in contrast, a strong Hall term is required: the Hall diffusivity must exceed the right-hand side of eq. 10. Furthermore a large Hall term appears able to drive instability even in the face of significant Ohmic and ambipolar diffusion (Lesur et al. 2014). We therefore place no upper limit on the Pedersen diffusivity under which the Hall-shear process can operate. If the fastest-growing wavelength were the determining factor, a too-strong magnetic field would again be disqualifying, and eq. 9 would also be a requirement for Hall-shear instability (Keith & Wardle 2015). However we note that the spectrum of unstable linear modes can extend to wavelengths shorter than the ideal-MHD cutoff if the Hall term is strong (Sano & Stone 2002), and that the Hall-shear instability operates with a midplane plasma beta near unity in Lesur et al. (2014) run 1-OHA-5. For these reasons we do not limit the field strengths at which Hall-shear instability is allowed.

Finally, to launch a magneto-centrifugal wind, the disk must be able to sustain vertical gradients in the field’s horizontal components. That is, the fields may not diffuse through the disk thickness in less than one orbit. This boils down to

$$\eta_{max} < c_s^2/\Omega \quad (11)$$
where $\eta_{\text{max}}$ is whichever is larger, $\eta_P$ or $\eta_H$ (Keith & Wardle 2015).

To summarize, each magnetic scenario implies a field strength, which together with the charged species’ populations yields the diffusivities. The MRT scenario is viable if the field and diffusivities satisfy eqs. 8, 9 and 10. The Hall-shear scenario is viable if the negation of eq. 10 holds. The magnetized wind is viable given eq. 11. To express these requirements compactly, below we use the definitions $\eta_{v_{Az}} \equiv v_{Az}^2/\Omega$ and $\eta_{cs} \equiv c_s^2/\Omega$.

5. RESULTS

5.1. X-ray Spectra

In all cases, the X-rays reaching the midplane have a much lower flux than those striking the disk atmosphere, and their spectrum favors higher energies. In figure 3 are the spectra observed at four heights above the planet in model 10S0. The lower-energy direct X-rays are absorbed high in the atmosphere, while higher-energy photons’ intensity declines with depth, as Compton down-scattering converts them to lower energies where they are more easily absorbed. Few or no photons reach the planet with energies below 10 keV, while the flux at energies above 10 keV is around 0.1% of that expected in the absence of the intervening disk material. Almost all photons reaching the midplane were emitted from the source with energies above 18 keV. The 8 keV dip in the spectrum at intermediate heights comes from the iron absorption edge near 7 keV, visible in figure 2.

The cleaner gaps opened by the Jupiters let more photons with energies below 10 keV reach the planet’s vicinity (figure 4). Still most of the ionization comes from photons emitted with energies above 10 keV, many of which scatter more than once because their single-scattering albedos exceed one-half.

5.2. Ionization Rates

The ionization rates for all models are plotted in figure 5 versus the column perpendicular to the midplane. In all cases the rates are column-independent in the optically-thin upper layers, where they are fixed by the flux of the numerous low energy X-rays. At greater columns, the disk is optically-thick to the softer X-rays, precipitating a sharp decline in the ionization rate. A scattering shoulder appears at columns greater than $N_H = 10^{23}$ cm$^{-2}$, where the contribution from scattered harder X-rays extends ionization into the disk interior. The ionization rate asymptotes at the highest columns in the Jupiter cases, because the total
Fig. 3.— X-ray spectra in model 10S0 at four locations directly over the planet. The top panel is at the disk surface, the next two are 1.25 and 0.81 au above the planet, and the bottom panel is at the midplane, on the planet’s orbit. Each is labeled with the vertical optical depth at 5 keV. Red curves show the photons received directly from the source, and blue curves the scattered photons. Each curve is surrounded by lighter shading marking the 1/\sqrt{N} uncertainty in the Monte Carlo results. Direct photons are thoroughly absorbed for vertical optical depths unity and greater. Only X-rays with energies over 10 keV reach the midplane in significant numbers, and all these have been scattered.

X-ray optical depth is low, so all points near the midplane see the remaining scattered X-rays.

At greater distances from the star, X-rays are less important because of the inverse-square falloff in their flux. It is thus worth considering whether the X-rays are competitive with the interstellar cosmic rays. These are attenuated by the protostar’s wind (Cleeves et al. [2013]), which however was likely funneled by the disk and its wind into a bipolar configuration. While the disk at 10 au could receive some cosmic rays focused by the disk wind or entering near the equatorial plane, molecular abundances observed at tens of au in at least one disk are well fit by a model without cosmic rays (Cleeves et al. [2015]). Thus the cosmic ray ionization rate inside 10 au is likely well below the interstellar value of about 10^{-17} s^{-1}. Thermal ionization and radionuclide decay ionization are orders of magnitude weaker still at 5 and 10 au.

X-ray ionization rates at the planet exceed 10^{-18} s^{-1} in all the models we consider except the one with a Saturn-mass planet at 10 au and no dust settling. Specifically, the
Fig. 4.— Paths in the \((r, z)\) plane of photon packets chosen randomly from among those reaching the planet’s vicinity in model 5J0, which has a Jupiter-mass planet at 5 au in a dusty disk. The X-ray source is marked by the cross in a circle near \((0.1, 0.1)\) au. Each path is red before the first scattering, and blue thereafter. Filled circles mark scattering points. The thick grey line shows the surface of unit vertical optical depth to 5-keV photons, while thin grey lines denote optical depths spaced by factors of ten. The first panel shows photons with energies up to 10 keV, for which the single-scattering albedo is below 50%. Most photons reaching the planet are scattered just once off the gap rim. The second panel is for higher-energy, higher-albedo photons, which often scatter repeatedly off the gap walls. Some scatter first in the disk surface layers interior to the gap, and a few pass near the rotation axis as they cross from one side of the disk to the other. The number of packets in each panel is in proportion to the energy band’s contribution to ionization near the planet.

X-ray ionization rates listed in the next-to-last column of table 1 are of order \(10^{-19}\) and \(10^{-16}\) s\(^{-1}\) for a Saturn- and Jupiter-mass planet, respectively, at 10 au in a disk with dust-to-gas mass ratio \(10^{-2}\). For a Saturn-mass planet, reducing the dust abundance by a factor 100 or moving the planet inward to 5 au increases the ionization rate an order of magnitude. For a Jupiter-mass planet, the low gas densities in the gap mean a low optical depth for X-rays traveling vertically. Reducing the dust abundance and associated X-ray optical depth therefore has little effect, increasing the ionization rate by less than 20%.

The final column in table 1 shows the ratio of the ionization rate near the planet to that at the same radius and mass column in a fit to the IG99 results (Turner & Sano 2008), the same fit considered by Keith & Wardle (2015). In several of our models with optically-thin gaps, the ratio exceeds unity, even though the column of gas overlying the planet would place it high in the atmosphere in the planet-free IG99 disk. This is possible both because the
Fig. 5.— Ionization rates vs. vertical column of hydrogen nuclei, above planets with the mass of Saturn (left) and Jupiter (right). The X-rays ionize faster than $10^{-18}$ s$^{-1}$ in all models except that with a Saturn-mass planet at 10 au in the dusty disk. For easy comparison between the Jupiter results, the models with the Jupiter at 10 au are shifted left, to column depths $10^6$ times lower. Also, the models with Saturn at 10 au are shifted left by half a decade.

Heavy elements yield an absorption opacity greater than the depleted case IG99 considered, as shown in figure 2, and because we include photons emitted near the star with energies up to 30 keV, while IG99 ended their calculations at 20 keV.

### 5.3. Magnetic Coupling

Inserting the ionization rates in the chemical reaction network yields equilibrium ionization states. For the cases with Saturn-mass planets and the dust depleted, the most abundant charged species at the planet’s location, by wide margins, are electrons and molecular ions. In contrast, the models with the full complement of dust have the positively- and negatively-charged grains more abundant at the planet than either electrons or ions. All the models have temperatures low enough near the planet that most metal atoms are adsorbed on the grains.

The charged species’ movements, and thus the diffusivities, depend on the magnetic field strength as discussed in sec. 4.3. In addition, the coupling criteria eqs. 8 to 10 are explicit functions of the field’s vertical component. Thus in table 2 we list the total and vertical field strengths just inside and outside the planet’s orbit, obtained using eqs. 6 and 7 from the corresponding mass flow rates of $5 \times 10^{-9}$ and $3 \times 10^{-8}$ $M_\odot$ yr$^{-1}$ respectively.
Table 2: Magnetic Field Strengths (mG) Yielding Disk Models’ Mass Flow Rates

<table>
<thead>
<tr>
<th>Model</th>
<th>$M$</th>
<th>MRT</th>
<th>HSI</th>
<th>MCW</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>$10^{-8} M_\odot$ yr$^{-1}$</td>
<td>$B$</td>
<td>$B_z$</td>
<td>$B$</td>
</tr>
<tr>
<td>10S0</td>
<td>0.5</td>
<td>18</td>
<td>4.1</td>
<td>64</td>
</tr>
<tr>
<td></td>
<td>3</td>
<td>45</td>
<td>10</td>
<td>160</td>
</tr>
<tr>
<td>10S2</td>
<td>0.5</td>
<td>16</td>
<td>3.7</td>
<td>58</td>
</tr>
<tr>
<td></td>
<td>3</td>
<td>40</td>
<td>9</td>
<td>140</td>
</tr>
<tr>
<td>10J0</td>
<td>0.5</td>
<td>15</td>
<td>3.5</td>
<td>55</td>
</tr>
<tr>
<td></td>
<td>3</td>
<td>38</td>
<td>8.5</td>
<td>130</td>
</tr>
<tr>
<td>10J2</td>
<td>0.5</td>
<td>15</td>
<td>3.4</td>
<td>54</td>
</tr>
<tr>
<td></td>
<td>3</td>
<td>38</td>
<td>8.4</td>
<td>130</td>
</tr>
<tr>
<td>5S0</td>
<td>0.5</td>
<td>49</td>
<td>11</td>
<td>170</td>
</tr>
<tr>
<td></td>
<td>3</td>
<td>120</td>
<td>27</td>
<td>420</td>
</tr>
<tr>
<td>5S2</td>
<td>0.5</td>
<td>44</td>
<td>9.9</td>
<td>160</td>
</tr>
<tr>
<td></td>
<td>3</td>
<td>110</td>
<td>24</td>
<td>380</td>
</tr>
<tr>
<td>5J0</td>
<td>0.5</td>
<td>41</td>
<td>9.2</td>
<td>150</td>
</tr>
<tr>
<td></td>
<td>3</td>
<td>100</td>
<td>23</td>
<td>360</td>
</tr>
<tr>
<td>5J2</td>
<td>0.5</td>
<td>41</td>
<td>9.2</td>
<td>140</td>
</tr>
<tr>
<td></td>
<td>3</td>
<td>100</td>
<td>22</td>
<td>360</td>
</tr>
</tbody>
</table>

We combine the charged particle populations with the field strengths in each of the three magnetic scenarios, to obtain the Ohmic, Hall, and ambipolar diffusivities. The coupling criteria are then fully specified and we can evaluate whether each scenario is viable. The results for locations just inside and outside the planet’s orbit appear in table 3. Checks mark scenarios that meet the conditions according to the dimensionless numbers in columns 6 to 9. These in turn come from eqs. 8 to 11 with magnetic fields from the corresponding scenarios. The dimensionless numbers are rounded to one significant figure.

In all cases, the MRT scenario yields a large ambipolar contribution to the Pedersen diffusivity, which is strong enough to invalidate the assumption that magneto-rotational turbulence is present. In some cases, the MRT scenario is also ruled out because it requires fields that are too strong for the magneto-rotational wavelength to fit into the disk thickness. In no case is the MRT prevented solely by the Hall diffusivity; thus we do not list the MRT scenario’s Hall number in table 3. In stark contrast to MRT, the HSI scenario implies magnetic fields such that the Hall length is great enough for the Hall-shear instability to operate in all models. In several of the cases that are better-ionized and also dense enough to avoid strong ambipolar diffusion, the maximum diffusivity is low enough for the gas near
Table 3: Magnetic Scenarios’ Viability Just Inside and Outside the Planet’s Orbit

| Model | $M/10^{-8} M_\odot$ yr$^{-1}$ | MRT | HSI | Wind | $\eta_{vA_x}/\eta_P$ | $\beta_z$ | $k^2/\eta_{vA_x}/|\eta_H|$ | Wind |
|-------|-------------------------------|-----|-----|------|-----------------------|---------|-----------------------------|------|
| 10S0  | 0.5                           | —   | ✓   | —    | 8e-7                  | 3e+3    | 1e-5                        | 2e-3 |
| 10S2  | 0.5                           | —   | ✓   | ✓    | 1e-2                  | 5e+3    | 3e-3                        | 1e+1 |
| 10J0  | 0.5                           | —   | ✓   | ✓    | 1e-2                  | 8e+2    | 3e-3                        | 6e+0 |
| 10J2  | 0.5                           | —   | ✓   | ✓    | 2e-2                  | 5e+1    | 9e-4                        | 3e+0 |
| 5S0   | 0.5                           | —   | ✓   | —    | 2e-2                  | 8e+0    | 7e-4                        | 5e-1 |
| 5S2   | 0.5                           | —   | ✓   | ✓    | 8e-5                  | 1e+2    | 4e-5                        | 6e-3 |
| 5J0   | 0.5                           | —   | ✓   | ✓    | 2e-3                  | 7e-1    | 8e-5                        | 2e-2 |
| 5J2   | 0.5                           | —   | ✓   | —    | 3e-2                  | 7e-1    | 5e-2                        | 2e-1 |
| 3     |                               | —   | ✓   | ✓    | 3e-2                  | 1e-1    | 8e-3                        | 3e-2 |

the gap’s midplane to act as the base for a magneto-centrifugal wind.

Profiles of the magnetic coupling through the column of material above the planet are shown in figure 6. The two cases plotted here bracket the range of the full set of models. The first, 10S0, has conditions least favorable for ionization: the planet lies far from the X-ray source, the full dust abundance means both greater X-ray optical depth and rapid recombination on grain surfaces, and the relatively large gas column in the gap likewise makes recombination quick. The other case in figure 6, model 5J2, is at the opposite extreme in all these respects, and has the highest midplane ionization fraction of all our models.

Because the X-ray intensity rises with height above the planet, while the recombination rate declines as the square of the density, the ionization fraction rises with height. In model 10S0, this means a fairly well-coupled layer at intermediate heights, while still further up, the declining density makes the ambipolar diffusivity dominant. In model 5J2, lower gas densities throughout mean the ambipolar term is strong even at the midplane. These trends in the three non-ideal magnetic diffusivities with mass column are similar to those in modeling of protostellar disks without gaps, for example by Wardle (2007) and Bai (2011, 2014).
The coupling conditions shown in figure 6 have the following implications for the three magnetic scenarios. From the top left panel, we see that in the 10S0 model under the turbulent scenario, the magnetic field near the midplane is weak enough to be consistent with magneto-rotational turbulence (red shading), but the diffusivity is too large for the instability to operate (the blue and green curves are below unity). Conversely in the atmosphere, the Hall diffusivity is low enough for turbulence (blue shading), but the field is too strong to allow it. Ambipolar diffusion is an obstacle to turbulence everywhere (green curve). The situation is still less favorable for magneto-rotational turbulence in the 5J2 model (top right panel), as the magnetic field is too strong even at the midplane. The middle row of panels shows that the criterion for Hall-shear instability is satisfied throughout both models (blue shading). The panels in the bottom row show that a magneto-centrifugal wind can be launched from a surface layer in the 10S0 case. The layer’s shallowness suggests the wind will not reach escape speed before ambipolar diffusion decouples it from the fields, ending magnetic acceleration and collimation. However, high in the atmosphere where the direct stellar far-ultraviolet photons are absorbed, much greater ionization fractions are expected (Perez-Becker & Chiang 2011; Bethell & Bergin 2011b). Since we do not treat the FUV photons, our calculations are not valid above the base of this layer. Determining the fate of the wind would furthermore require non-ideal MHD calculations spanning the gap and the nearby disk, where launching conditions may be quite different from the gap.

A note of caution is in order regarding timescales in the 5J2 model, with the Jupiter-mass planet at 5 au in the dust-depleted disk. The gas flows across the gap in just 10 years, while the reaction network takes about 300 years to reach equilibrium at the midplane. The ionization state in the gap will therefore reflect conditions upstream in its outer rim, rather than the local equilibrium values used for table 3 and figure 6. Since the denser rim material has slower ionization and faster recombination, this will mean the gap is even more poorly coupled to the magnetic fields. In measuring the chemical timescale, we have defined equilibrium as reached when all species’ abundances fall within a factor two of their asymptotic values. The equilibration timescales are sensitive to the density in the gap, which is quite uncertain and only crudely represented in our simple disk models, as discussed in section 2.

6. DISCUSSION AND CONCLUSIONS

We have carried out transfer calculations for X-rays emitted from the corona of a young star into a surrounding protostellar disk containing an embedded Saturn- or Jupiter-mass planet. The distribution of the material through which the X-rays pass is determined by the
planet’s tides, which open a gap in the disk, and by the starlight, whose heating determines the disk thickness under our assumption of vertical hydrostatic equilibrium. Some of the X-rays scattered in the disk atmosphere reach and ionize the gas in the planet’s vicinity. The ionization rates are comparable to or greater than those produced in the interstellar medium by cosmic rays. Ultraviolet radiation, which we do not treat, could more strongly ionize the planet’s vicinity in the Jupiter cases, if the planet opens a gap as clean as in the simple model we use. However the Saturn cases’ greater column density, which far exceeds ultraviolet photons’ absorption depth, means there are few options for raising the ionization rate near the planet above that provided by the X-rays.

The ionization rates depend mostly on the shape of the gap and the column of material within. The magnetic coupling results summarized next depend also on the recombination chemistry and magnetic field strength and orientation, for which the ranges of possibilities are wide. The coupling results should therefore be considered less certain.

We compute the equilibrium ionization state of the material near the planet using a simplified chemical reaction network including representative molecular and metal ions, and recombination on grain surfaces. From the populations of charged species, we compute the Ohmic, Hall, and ambipolar diffusivities under three scenarios for the magnetic fields: the accretion flow is driven by either (1) magneto-rotational turbulence, (2) Hall-shear instability, or (3) a magneto-centrifugal wind.

In all cases, the diffusivity is too high for magneto-rotational turbulence to transport angular momentum in the gas near the planet. Ambipolar diffusion decouples the magnetic fields from the disk’s neutral component over the length and time scales that would be associated with the turbulence. Hall-shear instability, in contrast, can drive the accretion flow in all cases: the field strengths needed for HSI to produce the assumed mass flow rates yield Hall diffusivities great enough for the HSI to operate.

A scenario with a magneto-centrifugal wind also appears viable in the less-diffusive cases we considered, especially if ultraviolet ionization sustains coupling at heights above where the X-rays are important. The least-diffusive models have Saturn-mass planets, which due to their lower mass only partly clear their gaps, yielding midplane gas densities great enough for ion-neutral collisions to reduce the rate of ambipolar diffusion. Planets growing from Saturn to Jupiter mass could thus experience a change in the ambient magnetic field configuration, and the pattern of gas flow in the gap, as they transition from a wind-plus-Hall-shear geometry to one involving the Hall shear alone. Only the Hall-shear magnetic scenario appears capable of driving the accretion flow in our models with Jupiter opening a gap at 5 au.
We conclude that the gas dynamics in gaps in the giant-planet-forming region of the Solar nebula and other similar protostellar disks depend critically on the non-ideal terms in the induction equation. While the X-rays yield gas-magnetic coupling too weak for magneto-rotational turbulence, material near the planet is ionized well enough that magnetic forces will alter its distribution, affecting the planet’s growth and orbital migration.

Although we have neglected the gradient term in the mass flow–stress relation (eq. 5) when estimating the field strengths needed to drive the flows, this term could dominate around planet-opened gaps, where the diffusivities and hence the magnetic fields may change over distances comparable to the density scale height. In the simple model disk we used, the mass flow rate changes across the planet’s orbit. If conditions vary sharply near the planet, the diffusivity will likewise have steep gradients. Magnetic fields have been shown to evolve toward a more nearly uniform radial profile than the gas in ideal-MHD, unstratified MRT calculations (Zhu et al. 2013). Evaluating the magnetic gradients in a more diffusive environment with radial and vertical structure in the ionization state will require further detailed MHD calculations. While 3-D MHD models of disks with planet-opened gaps have been constructed including the Ohmic diffusivity (Gressel et al. 2013), it now seems clear that the ambipolar and Hall terms are more important still.

We are grateful to Barbara Ercolano for advice on the X-ray opacities, and to Satoshi Okuzumi for providing his subroutine implementing the Bethell & Bergin cross-sections. Lynne Hillenbrand’s sponsorship of SYK at Caltech enabled this work to begin. The research was carried out in part at the Jet Propulsion Laboratory, California Institute of Technology, under a contract with the National Aeronautics and Space Administration, and with the support of the NASA Origins of Solar Systems program through grant 13-OSS13-0114. Copyright 2018 California Institute of Technology. Government sponsorship acknowledged.

REFERENCES

Fig. 6.— Viability of the magneto-rotational turbulence (MRT), Hall-shear instability (HSI), and magneto-centrifugal wind (MCW) scenarios (top to bottom), in the models with a Saturn in a dusty disk (10S0, left column), and a Jupiter in a dust-depleted disk (5J2, right column). Solid curves are for the mass flow rate found just outside the planet’s orbit, dashed curves for the lower rate just inside the orbit. Shading marks the portion of each solid curve that meets the constraints of eqs. 8 (green), 9 (red), 10 (blue), and 11 (purple). Darker sections of the blue curves show where the Hall diffusivity is positive. The MRT scenario meets its constraints nowhere, while the HSI scenario is viable everywhere. MCW yields diffusivities consistent with its own operation only in the Saturn case, where the gas over the planet is relatively dense. Conditions for the MCW are then met in a layer of the disk atmosphere.