Characterizing thermal sweeping: a rapid disc dispersal mechanism

James E. Owen,1,* Mathias Hudoba de Badyn,1,2 Cathie J. Clarke3 and Luke Robins3

1Canadian Institute for Theoretical Astrophysics, 60 St George Street, Toronto, M5S 3H8, Canada
2Department of Physics and Astronomy, University of British Columbia, Vancouver, V6T 1Z1 Canada
3Institute of Astronomy, Madingley Road, University of Cambridge, Cambridge CB3 0HA, UK

Accepted 2013 September 2. Received 2013 September 2; in original form 2013 July 22

ABSTRACT
We consider the properties of protoplanetary discs that are undergoing inside-out clearing by photoevaporation. In particular, we aim to characterize the conditions under which a protoplanetary disc may undergo ‘thermal sweeping’, a rapid (≲10^4 years) disc destruction mechanism proposed to occur when a clearing disc reaches sufficiently low surface density at its inner edge and where the disc is unstable to runaway penetration by the X-rays. We use a large suite of 1D radiation-hydrodynamic simulations to probe the observable parameter space, which is unfeasible in higher dimensions. These models allow us to determine the surface density at which thermal sweeping will take over the disc’s evolution and to evaluate this critical surface density as a function of X-ray luminosity, stellar mass and inner hole radius. We find that this critical surface density scales linearly with X-ray luminosity, increases with inner hole radius and decreases with stellar mass, and we develop an analytic model that reproduces these results. This surface density criterion is then used to determine the evolutionary state of protoplanetary discs at the point that they become unstable to destruction by thermal sweeping. We find that transition discs created by photoevaporation will undergo thermal sweeping when their inner holes reach 20–40 au, implying that transition discs with large holes and no accretion (which were previously a predicted outcome of the later stages of all flavours of the photoevaporation model) will not form. Thermal sweeping thus avoids the production of large numbers of large, non-accreting holes (which are not observed) and implies that the majority of holes created by photoevaporation should still be accreting. We emphasize that the surface density criteria that we have developed apply to all situations where the disc develops an inner hole that is optically thin to X-rays. It thus applies not only to the case of holes originally created by photoevaporation but also to holes formed, for example, by the tidal influence of planets.

Key words: protoplanetary discs – planetary systems – stars: pre-main-sequence.

1 INTRODUCTION
Understanding the dispersal of protoplanetary discs remains a key problem in star and planet formation. In particular, the time at which the disc disperses sets the time in which (gas) planets must form. Additionally, the method and time-scale on which disc clearing operates strongly influence the dust and gas physics in the disc, as well as the evolution of currently forming planets.

Observationally, optically thick primordial discs are known to survive for several Myr before being dispersed in both dust (Haisch et al. 2001; Hernández et al. 2007; Mamajek 2009) and gas (Kennedy & Kenyon 2009), leaving behind a pre-main-sequence star possibly surrounded by an optically thin debris disc, containing second generation dust particles (Wyatt 2008). A small fraction of young stars show evidence for a cleared gap or hole in their dust discs, indicated by a lack of opacity at near-infrared (NIR) wavelengths, but values comparable with an optically thick disc at mid-infrared (MIR) wavelengths. These ‘transition’ discs (Strom et al. 1989; Skrutskie et al. 1990; Calvet et al. 2002) have been interpreted as discs caught at a stage which is intermediate between disc bearing and discless. Furthermore, the lack of objects observed in regions of parameter space that represent discs that are no longer optically thick at infrared wavelengths indicates that the time-scale for disc dispersal is short compared to the disc’s lifetime (Kenyon & Hartmann 1995). Comparing the ratio of observed transition discs to primordial discs indicates that the clearing time-scale for the inner disc is roughly 10 per cent of the disc’s lifetime (Luhman et al.)
Characterizing thermal sweeping

2 BASICS OF THERMAL SWEEPING

In this section, we describe the theoretical basis for how X-ray irradiation of a disc with an inner hole leads to rapid, dynamical disc dispersal. We also present the back of the envelope calculation performed by Owen et al. (2012) as well as pointing out the assumptions that may break down.

2.1 Understanding thermal sweeping

In order to explain the origin of thermal sweeping, we show in Fig. 1 a schematic depiction of the flow structure that results from the X-ray irradiation of a disc with an inner hole, based on the simulations of Owen et al. (2010, 2011, 2012).

The X-rays heat up the inner disc and drive a subsonic wind radially inwards. The simulations of Owen et al. (2012) showed that the gas becomes significantly decelerated by centrifugal effects at a radius of roughly $R_{\text{hole}}/2$ and at this point is deflected vertically by pressure gradients, attaining a sonic transition at a height $z \sim R$.

Owen et al. (2012) showed that the mass flux in the wind is set at the sonic point and that the structure of the radially inflowing thermal wind region (see Fig. 1) can be understood in terms of a one-dimensional angular momentum conserving flow subject to this boundary condition on the mass flux. The total radial column density in this thermal wind region is given by

$$N \sim n R_{\text{hole}} = \frac{M_{\text{w}}}{2\pi \mu \kappa v R (H/R)} \times \left( \frac{R_{\text{hole}}}{10 \text{ au}} \right)^{-1} \left( \frac{H/R}{0.1} \right)^{-1}$$

Recall that viscous effects are negligible in the dynamical wind flow and thus the gas in this region conserves its specific angular momentum.
where \( N \) is the column, \( n \) is the number density, \( M_\nu \) the mass-loss rate, \( \mu \) the mean particle mass, \( v \) the gas velocity and \( H \) the scale-height. This is significantly lower than the column required to absorb the X-rays \((\sim 10^{22} \text{ cm}^{-2})\) – Ercolano et al. 2008, 2009. Thus, in this flow topology the material that dominates the absorption of the X-rays is not actually in the thermal wind, but is still bound to the disc, essentially in hydrostatic equilibrium with temperatures \(<1000 \text{ K}\).

Owen et al. (2012) argued that for this bound X-ray heated region to remain in dynamical equilibrium, it must have a radial pressure scale \((\Delta \equiv dP/d \log P)\) smaller than its vertical pressure scale \((d_z/d \log P \equiv H = c_s/\Omega)\). Since in order to obtain dynamical equilibrium, the disc must be able to adjust itself radially to changes in the vertical structure on a time-scale considerably shorter than the dynamical time-scale of the disc \((\sim 1/\Omega)\). In situations where \( \Delta \sim H \) (and where the time-scales for attaining equilibrium in the radial and vertical directions are comparable), this bound X-ray heated region is unstable to progressive penetration by the X-rays as follows. In the case of a small vertical expansion of the flow, the disc cannot regain equilibrium in the radial direction on the vertical expansion time-scale and so the radial column density at the mid-plane (and the local volume density) is reduced. Consequently, the temperature of the X-ray heated flow rises and the vertical expansion is accelerated. The further reduction in mid-plane column density results in further penetration of the X-rays and the process of vertical evacuation of the disc is repeated for a fresh layer of previously unheated disc material. The key difference between this behaviour and ‘normal’ photoevaporative clearing from a holed disc is that vertical expansion leads to runaway heating of the disc’s inner rim.

### 2.2 Column-limited estimate for thermal sweeping

Here, we repeat the simple analytical argument presented by Owen et al. (2012), but emphasize the assumptions made in its derivation. Since the X-ray bound heated region must be in approximate pressure equilibrium with the passively heated (by dust) disc at large radii, one can write

\[
k_b n_X T_X = P^\text{dust}_{\mu} \tag{2}
\]

where \( k_b \), \( n_X \) and \( T_X \) are the number density and temperature in the bound X-ray-heated region and \( P^\text{dust}_{\mu} \) the gas pressure in the passively heated disc. Writing \( \Delta \) as approximately \( N_X/n_X \), and using the fact that thermal sweeping begins when \( \Delta = H \), we can recast equation (2) in terms as a criterion for thermal sweeping as

\[
p^\text{dust}_{\mu} \leq \frac{N_X k_b T_X}{H} \tag{3}
\]

where \( N_X \) is the radial column density corresponding to complete absorption of the X-rays \((\sim 10^{22} \text{ cm}^{-2})\). In the case of a disc in vertical hydrostatic equilibrium, this can be recast in terms of a critical surface density for the inner rim of the dust-heated disc:

\[
\Sigma_{\text{TS}} = 0.4 \text{ g cm}^{-2} \left( \frac{T_X}{400 \text{ K}} \right)^{1/2} \left( \frac{T_{\text{dust}}}{20 \text{ K}} \right)^{-1/2}. \tag{4}
\]

Given \( T_{\text{dust}} \propto R^{-1/2} \), this resulted in a weak \( R^{1/2} \) scaling for the critical surface density with radius, with no explicit dependence on stellar mass or X-ray luminosity.

However, this derivation presented by Owen et al. (2012) makes two strong assumptions: (i) that the bound X-ray heated region always represents a column density of \( 10^{22} \text{ cm}^{-2} \) and (ii) the surface density at the transition between the X-ray heated disc and passively heated disc is representative of the peak surface density of the disc’s profile.

Assumption (i) may break down if the gas density is so high that X-ray heating is insufficient to heat the gas above the dust temperature even if it is optically thin to the X-rays. This typically occurs at ionization parameters \( \xi_{\text{min}} \lesssim 10^{-7} \text{ erg cm s}^{-1} \) (where \( \xi = L_X/n r^2 \)). In the case that the boundary of the X-ray heated region is set by the criterion \( \xi = \xi_{\text{min}} \), we can estimate the column density in the bound X-ray heated region as \( N = n \Delta \). At the point of onset of thermal sweeping, this can be written as \( N = (H/R)n R \), which (from the definition of the ionization parameter) becomes

\[
N = 10^{22} \text{ cm}^{-2} \left( \frac{H}{R} \right) \left( \frac{L_X}{10^{31} \text{ erg s}^{-1}} \right)^{-1} \left( \frac{\xi_{\text{min}}}{10^{-7} \text{ erg cm s}^{-1}} \right)^{-1}. \tag{5}
\]

Therefore, since we expect thermal sweeping to take place when the hole radius is large \((R_{\text{hole}} \gtrsim 10 \text{ au})\), then we need to consider cases where the total column density in the bound X-ray heated region is \( <10^{22} \text{ cm}^{-2} \). This means that equation (4) will not represent how the critical surface density for thermal sweeping varies in certain regions of the parameter space. In particular, if the criterion in some regions of parameter space is set by a minimum ionization parameter rather than total column density, then it should depend on the X-ray luminosity as well as inner hole radius.

In this paper, we assume without further discussion that \( \Delta = H \) is a necessary and sufficient condition for initiating thermal sweeping (an assumption that we will revisit in future two-dimensional studies motivated by the results of this paper.) Here, we use 1D radiation-hydrodynamic models to perform a parameter space study, calculating how the critical surface density \( \Sigma_{\text{TS}} \) – defined as the peak surface density at the inner edge of the dust-heated disc for which \( \Delta = H \) varies with stellar mass, X-ray luminosity and inner hole radius. The advantage of a one-dimensional treatment is that it is possible to perform a large suite of high-resolution simulations that are beyond the scope of 2D simulations and that can inform the choice of parameters and resolution requirements of future 2D studies. The justification for a one-dimensional treatment is that – prior to the onset of thermal sweeping – the flow is to a good approximation one-dimensional in the vicinity of the inner edge.

### 3 NUMERICAL METHOD

In order to calculate the inner disc structures, we must make use of numerical hydrodynamic calculations. We use the ZEUS code (Stone & Norman 1992; Hayes et al. 2006), which has been modified by Owen et al. (2010, 2012), to include the heating of circumstellar material by stellar X-ray photons. We use this modified code to calculate the mid-plane inner disc structure; therefore, we run the code radially in one dimension (although we also evolve the azimuthal velocity\(^2\)). In the absence of X-ray heating, the disc’s temperature is set to the mid-plane temperature expected for a passively irradiated protoplanetary disc (e.g. Chiang & Goldreich 1997; D’Alessio, Calvet & Hartmann 2001):

\[
T_{\text{mid}} = \max \left( T_{\text{in}}, \left( \frac{R}{1 \text{ au}} \right)^{-1/2}, 10 \text{ K} \right). \tag{6}
\]

\(^2\) Often referred to as 1.5D simulations.
where $T_{1\text{au}}$ is the mid-plane temperature at 1 au, and we adopt typical values of 100 K for a 0.7 $M_\odot$ star and 50 K for a 0.1 $M_\odot$ star (D' Alessio et al. 2001). As initial conditions, we adopt a $\Sigma \propto R^{-1}$ profile, where in order to convert between surface density and mid-plane gas density, we assume that the disc is vertically isothermal and use

$$
\Sigma(R) = \sqrt{2 \pi H(R) \rho(R)}.
$$

(7)

This mid-plane density profile is setup in hydrostatic equilibrium including the necessary modification to the Keplerian rotation due to the radial pressure gradient.

In order for the simulations to obtain stable configuration, the boundary conditions require careful consideration. The outer boundary condition is implemented in order to maintain hydrostatic equilibrium across the outer boundary throughout the simulation, so subsonic radial flows are not allowed to interact with the boundary. Explicitly, we do this by maintaining a $\Sigma \propto R^{-1}$ power law across the boundary and adjusting the azimuthal velocity in the boundary cells in order to maintain this equilibrium. In X-ray heated winds, the mass flux is determined explicitly by the conditions at the sonic surface (Owen & Jackson 2012; Owen et al. 2012). It is not possible to follow the flow to the sonic point in a one-dimensional simulation since this occurs at a height $Z \sim R$ above the disc. Therefore, we adjust the inner boundary condition in order to match the flow properties (density, temperature and velocity) with the results of an on-axis streamline taken from the multidimensional inner hole simulations performed by Owen et al. (2012). Such a method is sufficient provided the inner boundary lies between the centrifugal radius and inner hole radius (so the problem remains essentially 1D, see Fig. 1).

The radial grid is non-uniform, spanning $[R_i, 6R_i]$ – where $R_i$ is the radius of the inner boundary – with a high resolution of 200 cells; typically, this gives ~10–15 cells per radial pressure scaleheight in the inner hole region. In general, the structures are spatially resolved with a lower resolution, typically 100 cells; however, the higher resolution reduces the scatter in properties calculated (particularly $\Delta$). Under this setup, the inner disc is then irradiated, with the X-rays driving a flow off the inner disc with a mass flux specified by the conditions at the inner boundary; this flow slowly erodes material from the disc resulting in an inner hole radius that increases slowly with simulation time. The evolution is slow enough that the structure maintains a quasi-dynamical equilibrium at each time step. We stop the simulation once the inner hole moves to such a radius that the inner boundary is no longer outside the centrifugal radius and the 1D approximation is no longer appropriate. In Figs 2 and 3, we show an example of the quasi-steady state flow structure obtained from one of the simulations, for a disc around a 0.1 $M_\odot$ star with an X-ray luminosity of $2 \times 10^{30}$ erg s$^{-1}$.

Using 1D simulations allows us to probe a large range of the observable parameter space of disc models at high resolution. We perform roughly 5000 simulations covering a range of initial surface densities and radial ranges for two stellar masses. For the 0.1 $M_\odot$ star, we consider X-ray luminosities$^3$ of $2 \times 10^{29}$ and $2 \times 10^{30}$ erg s$^{-1}$ and hole sizes in the range ~1–20. For the 0.7 $M_\odot$ star, we consider X-ray luminosities of $2 \times 10^{29}$, $2 \times 10^{30}$ and $2 \times 10^{31}$ erg s$^{-1}$ with inner hole sizes in the range ~10–70 au. For each simulated radial range, we vary the initial surface density

$$
\rho \propto \frac{2\pi R^2 G M}{(R_c^2 - R^2)^{1/2}}
$$

normalization at intervals of 0.15 dex over the range expected in clearing discs. These simulations are then used to compute how the ratio of $\Delta/H$ varies with surface density; thus, extracting the critical surface density for thermal sweeping and how it varies with stellar mass, X-ray luminosity and inner hole radius.

### 4 RESULTS

Our suite of simulations cover a large range of parameter space and provide the mid-plane structure of discs with large inner holes. For a given inner hole radius and surface density normalization in the disc, we use the simulation to solve for the ratio of $\Delta/H$ and can hence determine how this ratio varies as a function of the maximum surface density in the dust-heated disc. An example of how $\Delta/H$ varies with the peak surface density in a disc is shown in Fig. 4.

![Figure 2. Density structure for a simulation of a disc with an inner hole around a 0.1 $M_\odot$ star with an X-ray luminosity of $2 \times 10^{30}$ erg s$^{-1}$](https://example.com/figure2.png)

Figure 2. Density structure for a simulation of a disc with an inner hole around a 0.1 $M_\odot$ star with an X-ray luminosity of $2 \times 10^{30}$ erg s$^{-1}$. We note that the outer boundary is not shown and is at a radius of 38 au in this calculation.

![Figure 3. Temperature structure for the same simulation shown in Fig. 2.](https://example.com/figure3.png)

Figure 3. Temperature structure for the same simulation shown in Fig. 2.

3 Although low-mass stars may have lower X-ray luminosities, the simulated range is sufficient to calibrate the model discussed in Section 5, which can be applied to the full range of X-ray luminosities.
Critical surface density for thermal sweeping ($\Delta = H$) shown as a function of inner hole radius around a 0.7 $M_\odot$ mass star for three X-ray luminosities: $2 \times 10^{31}$ erg s$^{-1}$ (triangles), $2 \times 10^{30}$ erg s$^{-1}$ (circles) and $2 \times 10^{29}$ erg s$^{-1}$ (squares). The error bars show the uncertainty in the calculation of the critical surface density introduced by the finite grid resolution (see the text and Fig. 4). We also show a $\Sigma \propto R^{1/4}$ fit (performed at 10 au) as the dotted line, to represent the column-limited model (equation 4), which clearly fails to reproduce the simulations.

**4.1 Variation with inner hole radius and stellar mass**

The resulting variation in the critical surface density for thermal sweeping is shown as a function of inner hole radius and X-ray luminosity in Figs 5 and 6 for discs around 0.7 and 0.1 $M_\odot$ solar mass stars, respectively. The finite grid resolution that results in the scatter in Fig. 4 gives rise to an uncertainty in the surface density of $\pm 0.05$ dex which we represent as error bars in the figures. We emphasize that this scatter does not mean the simulations are not spatially resolved, in fact the simulations are well resolved (we perform several resolution tests and note the simulations are not spatially resolved, in fact the simulations are well resolved). However, we do not find the shallow $R^{1/4}$ dependence expected from the Owen et al. (2012) model, but a steeper variation with inner hole radius. Furthermore, we find an explicit, roughly linear, dependence on the X-ray luminosity not predicted by Owen et al. (2012). In Section 2, we outlined the two assumptions that underpinned the previous calculation: that the total column density in the bound X-ray heated layer is $\sim 10^{23}$ cm$^{-2}$ and that the surface density at the inner edge of the passively heated disc may not be representative of the peak surface density in the disc. We find that while the second clause certainly plays a role, it is not the major reason why the simulations do not show the expected dependence. The main reason is that the total column density in the bound X-ray heated layer is not always $10^{23}$ cm$^{-2}$ and in the majority of realistic cases is lower than this value and a function of radius.

**5 DISCUSSION AND OBSERVATIONAL IMPLICATIONS**

In the previous section, we have shown how the critical surface density for thermal sweeping varies with stellar mass, inner hole radius and X-ray luminosity. As expected, we find that thermal sweeping will proceed once the remaining mass in the disc is low $\lesssim 1 M_\text{ult}$. However, we do not find the shallow $R^{1/4}$ dependence expected from the Owen et al. (2012) model, but a steeper variation with inner hole radius. Furthermore, we find an explicit, roughly linear, dependence on the X-ray luminosity not predicted by Owen et al. (2012). In Section 2, we outlined the two assumptions that underpinned the previous calculation: that the total column density in the bound X-ray heated layer is $\sim 10^{23}$ cm$^{-2}$ and that the surface density at the inner edge of the passively heated disc may not be representative of the peak surface density in the disc. We find that while the second clause certainly plays a role, it is not the major reason why the simulations do not show the expected dependence. The main reason is that the total column density in the bound X-ray heated layer is not always $10^{23}$ cm$^{-2}$ and in the majority of realistic cases is lower than this value and a function of radius.

For example, we show in Fig. 7 how the total column density and the ionization parameter at the transition from X-ray-bound to dust-heated layer varies with X-ray luminosity. At the highest X-ray luminosity of $2 \times 10^{31}$ erg s$^{-1}$, we find that the surface density is roughly constant at small radius ($R_\text{hole} \lesssim 10$ au). This figure clearly shows that, at about 10 au the X-ray heated region makes a transition from being column limited to being ionization parameter limited or more correctly X-ray temperature limited.

**5.1 An improved model for thermal sweeping**

We can now use the above result: that the properties of the warm X-ray heated bound portion of the disc are set by a minimum...
ionization parameter rather than total column condition. Since the gas velocity in the bound portion of the disc (both dust and X-ray heated) is extremely subsonic, we can treat these regions of the disc as being approximately hydrostatic. Inside a radius $R_c$ (i.e. radius in the disc at which the source of material for the wind is in centrifugal balance), the specific angular momentum is constant and the effective gravity in the radial direction is given by

\[ g_{\text{eff}} = -\frac{GM_s}{R^2} \left( 1 - \frac{R_c}{R} \right) = \Omega^2 \delta(R), \]  

where $\delta(R) = R_c - R$ and $\Omega$ is the local Keplerian velocity. Now, using the definition that $\Delta$ is the radial pressure scalelength in the X-ray heated gas, we see that

\[ \Delta \equiv \frac{c_X}{g_{\text{eff}}} = \frac{H_X}{\delta}, \]  

where $c_X$ is the sound speed at the outer edge of the X-ray heated disc, and we have applied the condition for hydrostatic equilibrium normal to the disc plane, i.e. $c_X = H_X \Omega$. So, on the point of thermal sweeping ($\Delta = H_X$), equation (9) tells us that $\delta = H_X$ so that $\delta \ll R_c$. Thus, we proceed in the limit $\delta \ll R_c$ and retain terms in $g_{\text{eff}}$ to first order in $\delta/R$: we solve for radial hydrostatic equilibrium in the dust-heated disc, where the approximation $\delta \ll R_c$ ensures that the dust-heated disc is approximately isothermal at a temperature $T_c$. (We note that while $T_c$ varies with radius on a scalelength $\sim R_c$, making the isothermal approximation is formally equivalent to only retaining terms first order in $\delta/R$.) This then yields a Gaussian profile for the dust-heated disc such that

\[ P(\delta) = P_c \exp \left[ -\left( \frac{1}{2} \frac{\delta}{H_c} \right)^2 \right], \]  

where $P_c$ is the pressure at $R_c$ and $H_c$ is the hydrostatic scaleheight of the disc in the vertical direction at $R = R_c$. The innermost extent of the dust-heated disc is thus set by the condition of pressure balance with the X-ray heated region (pressure $P_X$) – which we set as $R_{\text{hole}}$ for convenience – and is given by

\[ \delta(R_{\text{hole}}) = 2H_c \sqrt{\log \left( \frac{P_c}{P_X} \right)^2}. \]  

Applying the condition for thermal sweeping ($\Delta = H_X$) allows equation (10) to be rewritten as

\[ \frac{P_c}{P_X} = \exp \left[ \frac{1}{2} \left( \frac{H_X}{H_c} \right)^2 \right] = \exp \left( \frac{T_X}{2T_c} \right). \]

We can qualitatively understand this result by considering a case where $T_X/T_c$ is fixed. If $P_c/P_X$ is high, then the X-ray heated gas can only achieve pressure equilibrium with the cold gas at a point that is relatively far from $R_c$, so the effective gravity is large (equation 8) and the radial scalelength correspondingly small (equation 9). However, as $P_c/P_X$ is reduced, the X-ray heated region is pushed closer to $R_c$, where $g_{\text{eff}}$ is low and the radial scalelength becomes larger. When $P_c/P_X$ is reduced sufficiently, the thermal sweeping criterion ($\Delta = H_X$) is activated.

This minimum value of $P_c$ can be readily converted into a minimum surface density criterion since in hydrostatic equilibrium $P_c \propto \Sigma_c \Omega$, where $c_c$ is the sound speed at $R_c$. Thus, we have

\[ \Sigma_{\text{TS}} = \sqrt{\frac{2\pi}{c_c}} \frac{P_c}{c_c \Omega} \exp \left( \frac{T_X}{2T_c} \right). \]

However, equation (13) allows us to only evaluate $\Sigma_{\text{TS}}$ (for a given radius in the disc and corresponding temperature of the dust-heated disc) if one also knows $T_X$ and $P_X$. It is now necessary to consider the form of the relationship between $\xi$ and $T$ which we use to assign X-ray temperatures. In the dense conditions at the base of the X-ray heated flow (i.e. at low ionization parameter), the relationship is nearly flat at a temperature of $\sim 100$ K and then steepens at $\xi = \xi_{\text{min}}$, with pressure then declining mildly as the temperature is reduced below $\sim 100$ K, where the structure of the temperature–ionization parameter relation is determined by the falling X-ray heating efficiency with ionization fraction (Xu & McCray 1991), in combination with cooling by lines (Owen et al. 2010). The number density on the cold side of the interface is such that the X-ray temperature at this point is equal to $T_c$; since in this case the X-rays would be unable to heat this region above $T_c$. This condition sets the pressure at the interface; on the X-ray heated side of the interface, the gas shares the same pressure but has a temperature of $\sim 100$ K, since this places it in the nearly isothermal portion of the $T$–$\xi$ relationship. We can see from Fig. 8 that a line of constant pressure ($T \propto \xi$) that originates on the $\xi$–$T$ curve at a typical dust temperature of $10$–$100$ K crosses the curve at a value of $\xi_{\text{min}}$, as indicated by our simulations (see Fig. 7). We can therefore simply set $P_X$ and $T_X$ by requiring $\xi = \xi_{\text{min}}$ at this point, where comparing
with Fig. 8 and our simulations indicates a sensible choice for 
\[ \xi_{\text{min}} = 3 \times 10^{-7} \text{ erg s}^{-1} \text{ cm}^{-1}. \]
Substituting this condition in equation (13), we obtain
\[ \Sigma_{\text{TS}} = \sqrt{\Sigma_0 T_X \Omega_{\xi_{\text{min}} R_{\text{hole}}}} \exp \left( \frac{T_X}{2T_c} \right) \]
\[ = 0.14 \text{ g cm}^{-2} \left( \frac{L_X}{10^{30} \text{ erg s}^{-1}} \right) \left( \frac{T_{1\text{au}}}{100 \text{ K}} \right)^{-1/2} \]
\[ \times \left( \frac{M_*}{0.7 \text{ M}_\odot} \right)^{-1/2} \left( \frac{R_{\text{hole}}}{10 \text{ au}} \right)^{-1/4} \]
\[ \times \exp \left[ \left( \frac{R_{\text{hole}}}{10 \text{ au}} \right)^{1/2} \left( \frac{T_{1\text{au}}}{100 \text{ K}} \right)^{-1/2} \right]. \]
Equation (15)

Note that in constructing equation (15) from equation (13), we have assumed that the temperature in the dust-heated disc follows the evolution of a group of discs evolving under photoevaporation. The resulting radius distribution of maximum cavity radius for transition discs created by photoevaporation. The resulting radius distribution is shown in Fig. 10, where we plot the maximum inner hole radius reached before thermal sweeping takes over for discs evolving around a 0.7 M_\odot star with different X-ray luminosities.

Furthermore, the general form (equation 14) could be applied to cases where an optically thin (to the X-rays) gap is created by a planet (e.g. Rice et al. 2006; Zhu et al. 2012) or by a combination of photoevaporation and a planet (Rosotti et al. 2013) to see if planet formation could trigger rapid disc dispersal through thermal sweeping.

5.2 Properties of observed transition discs
Confident that our model presented above can accurately predict the surface density of discs that are unstable to rapid clearing by thermal sweeping, we can now make observational predictions for the inner hole radius at which thermal sweeping clears the disc. We use equation (15) to re-analyse the disc population synthesis model performed by Owen et al. (2011), which was designed to match the evolution of disc fraction as a function of time. The model well reproduced a large fraction of transition discs with small holes and low accretion rates and fully explains the population of transition discs with low-mm fluxes that Owen & Clarke (2012) identified as discs most likely to be transitioning from disc bearing to a diskless state. However, in doing so the model also predicted a large population of ‘relic’ discs with large holes and no accretion. This population is yet to be observed, and Owen et al. (2012) noted that this problem would be alleviated by the thermal sweeping mechanism.

Thus, we take the synthetic disc population run by Owen et al. (2011) and assume that the disc is dispersed instantaneously by thermal sweeping once the disc’s surface density drops below \( \Sigma_{\text{TS}} \). We determine the distribution of maximum cavity radius for transition discs created by photoevaporation. The resulting radius distribution is shown in Fig. 10, where we plot the maximum inner hole radius reached before thermal sweeping takes over for discs evolving around a 0.7 M_\odot star with different X-ray luminosities.

Fig. 10 shows that thermal sweeping prevents the formation of relic discs (i.e. those with large holes and no accretion), explaining the lack of a significant population of transition discs observed to have large holes with no accretion. The trend of smaller hole radii reached around higher X-ray luminosity stars is due to the ability of thermal sweeping to begin at higher surface densities with higher X-ray luminosities. At low X-ray luminosities (<10^{29} \text{ erg s}^{-1}), the hole radius is independent of X-ray luminosity: this is because both \( \Sigma_{\text{TS}} \) and the surface density of the disc at a given radius at the point of photoevaporative hole opening scale linearly with X-ray luminosity, and so an equality holds.

![Figure 9](https://example.com/figure9.png)

**Figure 9.** Comparison between the thermal sweeping model presented in equation (15): plotted as \( \Sigma_{\text{TS}}/L_X \) versus radius, and the simulations presented in Section 4. The open points and dashed line are for the 0.7 M_\odot star and the filled points and solid line are for the 0.1 M_\odot star. The square, circular and triangular points are for X-ray luminosities of \( 2 \times 10^{29} \), \( 2 \times 10^{30} \) and \( 2 \times 10^{31} \text{ erg s}^{-1} \), respectively.

![Figure 10](https://example.com/figure10.png)

**Figure 10.** Maximum inner hole radii reached by clearing discs before thermal sweeping takes over. Each point represents an individual disc model taken from the population synthesis calculation of Owen et al. (2011), which follows the evolution of a group of discs evolving under photoevaporation around a 0.7 M_\odot star with an X-ray luminosity drawn from the observed X-ray luminosity function (Güdel et al. 2007).
is always achieved at a fixed radius. At larger X-ray luminosities, the process of ‘photoevaporation starved accretion’ (Drake et al. 2009; Owen et al. 2011) causes a slightly sublinear dependence on X-ray luminosity for the disc surface density at hole opening. Consequently, the thermal sweeping radius declines somewhat with X-ray luminosity. However, at the highest X-ray luminosities this is countered by the fact that gap formation occurs at earlier times and hence higher surface densities for stars with higher X-ray luminosities.

Furthermore, the rapid dispersal of transition discs created by photoevaporation at radii ~45 au means that the ratio between the time spent in an accreting transition disc phase and that in a non-accreting transition disc phase is greatly enhanced, compared to the photoevaporation model analysed in Owen et al. (2011). For example, our new thermal sweeping models predict that the median model \( L_X = 1.1 \times 10^{38} \text{ erg s}^{-1} \) (fig. 9 of Owen et al. 2011) initially opens the hole at ~2.5 au and the disc is destroyed by thermal sweeping when \( R_{\text{hole}} \sim 30 \text{ au} \), with a remaining disc mass of a few Jupiter masses. Furthermore, it spends approximately \( 2 \times 10^7 \text{ years} \) as an accreting transition disc and \( 6.5 \times 10^2 \text{ years} \) as a non-accreting transition disc. Thus, the addition of thermal sweeping suggests that transition discs created by photoevaporation spend the majority of their life as accreting transition discs. Due to the variation of thermal sweeping with stellar mass, thermal sweeping can begin earlier around lower mass stars and may become inefficient around intermediate-mass stars, where the X-ray luminosity no longer increases with the stellar mass above 1–2 \( M_\odot \) (e.g. Güdel et al. 2007; Albacete Colombo et al. 2007).

5.3 Possible limits for the time-scale of thermal sweeping

One possible limit on the time-scale at which thermal sweeping proceeds through the disc is energetic. The multidimensional simulations that were initially used to investigate thermal sweeping enforced local radiative equilibrium. This condition meant that energetic considerations were neglected [see discussion in Owen et al. 2010, who demonstrated that for standard X-ray photoevaporation, the mechanical luminosity of the wind is a comfortably small fraction (\( \lesssim 8 \) percent) of the X-ray luminosity of the source]. So, \( P_{\text{Dv}} \) work done escaping from the potential was not accounted for. The dynamical time-scale on which thermal sweeping begins is not going to be limited by this consideration (because the mass at the inner edge is small). However, as most of the disc material is at large radius, it may take longer than the dynamical time-scale at the inner edge for the disc material to heat up to the equilibrium temperature. We can estimate an ‘energy-limited’ time-scale for clearing to occur by comparing the gravitational potential energy stored in the disc to the rate at which the X-rays inject energy into the system.

We can calculate the gravitational binding energy of a power-law disc of the form \( \Sigma = \Sigma_{\text{hole}}(R/R_{\text{hole}})^{-d} \) as

\[
U_{\text{disc}} = \int_{R_{\text{hole}}}^{R_{\text{out}}} 2\pi R dR \frac{GM_\star}{R} \Sigma_{\text{hole}} \left(\frac{R}{R_{\text{hole}}}\right)^{-1} \\
= 2\pi G M_\star \Sigma_{\text{hole}} R_{\text{hole}} \log \left(\frac{R_{\text{out}}}{R_{\text{hole}}}\right) \\
\sim 10^{40} \text{ erg} \left(\frac{M_\star}{0.7 M_\odot}\right) \left(\frac{\Sigma_{\text{hole}}}{0.3 \text{ g cm}^{-2}}\right) \left(\frac{R_{\text{hole}}}{10 \text{ au}}\right).
\]  

Comparing this binding energy to the received X-ray luminosity, we find an energy-limited clearing time-scale of

\[
\tau_{\text{clear}} = \frac{U_{\text{disc}}}{\epsilon L_X} = 2 \times 10^3 \text{ years} \left(\frac{L_X}{10^{38} \text{ erg s}^{-1}}\right)^{-1} \left(\frac{\epsilon}{0.25}\right)^{-1} \times \left(\frac{R_{\text{hole}}}{10 \text{ au}}\right) \left(\frac{M_\star}{0.7 M_\odot}\right) \left(\frac{\Sigma_{\text{hole}}}{0.14 \text{ g cm}^{-2}}\right),
\]  

where \( \epsilon \) represents the fraction of X-rays intercepted by the disc (the X-ray photosphere to the star occurs at height \( >H_\star \)).

6 SUMMARY

In this work, we have investigated the rapid disc dispersal mechanism for holed disc ‘thermal sweeping’, introduced by Owen et al. (2012), which takes over from photoevaporation and destroys the remaining disc material on a short (\( \lesssim 10^5 \) year) time-scale once the surface density of the disc has dropped to sufficiently low values. We use the criterion suggested by Owen et al. (2012) (involving equality of the disc’s vertical scaleheight and the radial thickness of the X-ray heated bound gas at the rim of the holed disc) in order to define the point at which a disc will undergo rapid dispersal, using a large suite of 1D radiation-hydrodynamical simulations. In this way, we identify the surface density at which thermal sweeping proceeds as a function of X-ray luminosity, stellar mass and inner hole radius. Our results however do not replicate the analytic estimate presented by Owen et al. (2012) since this was based on the assumption that the extent of the X-ray heated region at the disc’s inner rim is set by a fixed absorption column. Instead, we find (apart from the largest X-ray luminosities and smallest radii) that the limits of X-ray heating are set by the ionization parameter falling to the value \( \xi_{\text{min}} \), below which the temperature–ionization parameter relation steepens strongly (see Fig. 8), at which point the X-rays do not heat the gas above the local disc temperature. We use this result to derive a new criterion (equation 15) for how the critical surface density for thermal sweeping varies with the physical properties of disc systems and find that this provides an excellent match to the simulation results (see Fig. 9).

Our main findings are summarized below.

(i) Thermal sweeping can rapidly clear the disc once the disc mass drops below a few Jupiter masses and the inner hole is at sufficiently large radii \( \gtrsim 10 \text{ au} \).

(ii) The critical surface density for thermal sweeping to proceed scales linearly with X-ray luminosity and increases with inner hole radius.

(iii) Thermal sweeping is considerably more efficient at clearing discs around lower mass stars and could possibly become ineffective around intermediate-mass stars which have comparatively weaker X-ray emission.
Using the derived condition for thermal sweeping, we show that transition discs created through photoevaporation are completely destroyed once their inner hole radii reach sizes $>20–40$ au.

Future multidimensional radiation-hydrodynamic simulations are required to understand the details of thermal sweeping. In particular, energy considerations raise questions about how rapid thermal sweeping can be in heating the outer parts of the disc once runaway penetration is in progress. Additionally, penetration and heating of the disc by the far-ultraviolet (FUV) radiation field (e.g. Gorti & Hollenbach 2009; Gorti, Dullemond & Hollenbach 2009) may aid in thermal sweeping. The inclusion of FUV heating in hydrodynamic simulations still remains challenging, and we cannot speculate on the role played by FUV heating in thermal sweeping at this stage.

ACKNOWLEDGEMENTS

We are grateful to the referee for comments that improved the paper. We thank Giovanni Rosotti, Barbara Ercolano, Geoff Vasil, Emmanuel Jacquet and Tom Haworth for insightful discussions. JEO is grateful to hospitality from the IoA, Cambridge, during the initial stages of the work. MHB acknowledges support of an NSERC summer research grant held at CITA during 2012. The calculations were performed on the Sunnyvale cluster at CITA which is funded by the Canada Foundation for Innovation. We would like to acknowledge the Nordita programme on Photo-Evaporation in Astrophysical Systems (June 2013) where part of the work for this paper was carried out.

REFERENCES

Albacete Colombo J. F., Flaccomio E., Micela G., Sciortino S., Damiani F., 2007, A&A, 464, 211

Alexander R. D., Armitage P. J., 2009, ApJ, 704, 989

Alexander R. D., Clarke C. J., Pringle J. E., 2006, MNRAS, 369, 229

Andrews S. M., Williams J. P., 2005, ApJ, 613, 1134

Andrews S. M., Williams J. P., 2007, ApJ, 671, 1800

Andrews S. M., Wilner D. J., Espaillat C., Hughes A. M., Dullemond C. P., McClure M. K., Qi C., Brown J. M., 2011, ApJ, 732, 42

Armitage P. J., Hansen B. M. S., 1999, Nat, 402, 633

Birnstiel T., Andrews S. M., Ercolano B., 2012, A&A, 544, A79

Calvet N., D’Alessio P., Hartmann L., Wilner D., Walsh A.,Sitko M., 2002, ApJ, 568, 1008

Calvet N. et al., 2005, ApJ, 630, L185

Chiang E. I., Goldreich P., 1997, ApJ, 490, 368

Cieza L. A. et al., 2013, ApJ, 762, 100

Clarke C. J., Owen J. E., 2013, MNRAS, 433, L69

Clarke C. J., Gendrin A., Sotomayor M., 2001, MNRAS, 328, 485

D’Alessio P., Calvet N., Hartmann L., 2001, ApJ, 553, 321

Drake J. K., Ercolano B., Flaccomio E., Micela G., 2009, ApJ, 699, L35

Dullemond C. P., Dominik C., 2005, A&A, 434, 971

Duvert G., Guilloteau S., Ménard F., Simon M., Dutrey A., 2000, A&A, 355, 165

Ercolano B., Drake J. J., Raymond J. C., Clarke C. C., 2008, ApJ, 688, 398

Ercolano B., Clarke C. J., Hall A. C., 2011, MNRAS, 410, 671

Espaillat C. et al., 2010, ApJ, 717, 441

Gorti U., Hollenbach D., 2009, ApJ, 690, 1539

Gorti U., Dullemond C. P., Hollenbach D., 2009, ApJ, 705, 1237

Güdel M. et al., 2007, A&A, 468, 353

Haisch K. E., Jr Lada E. A., Lada C. J., 2001, ApJ, 553, L153

Hayes J. C., Norman M. L., Fiedler R. A., Bordner J. O., Li P. S., Clark S. E., ud-Doula A., Mac Low M.-M., 2006, ApJS, 165, 188

Hernández J. et al., 2007, ApJ, 671, 1784

Kennedy G. M., Kenyon S. J., 2009, ApJ, 695, 1210

Kenyon S. J., Hartmann L., 1995, ApJS, 101, 117

Koeperl C. M., Ercolano B., Dale J., Teixeira P. S., Ratzka T., Spezzi L., 2013, MNRAS, 428, 3327

Luhman K. L., Allen P. R., Espaillat C., Hartmann L., Calvet N., 2010, ApJS, 186, 111

Mamajek E. E., 2009, in Usuda T., Tamura M., Ishii M., eds, AIP Conf. Ser. Vol. 1158, Exoplanets and Disks: Their Formation and Diversity. Am. Inst. Phys., New York, p. 3

Mathews G. S., Williams J. P., Menard F., 2012, ApJ, 753, 59

Nayakshin S., 2013, MNRAS, 431, 1432

Owen J. E., Clarke C. J., 2012, MNRAS, 426, L96

Owen J. E., Jackson A. P., 2012, MNRAS, 425, 2931

Owen J. E., Ercolano B., Clarke C. J., Alexander R. D., 2010, MNRAS, 401, 1415

Owen J. E., Ercolano B., Clarke C. J., 2011, MNRAS, 412, 13

Owen J. E., Clarke C. J., Ercolano B., 2012, MNRAS, 422, 1880

Rice W. K. M., Armitage P. J., Wood K., Lodato G., 2006, MNRAS, 373, 1619

Rosotti G. P., Ercolano B., Owen J. E., Armitage P. J., 2013, MNRAS, 430, 1392

Skrutskie M. F., Dutkevitch D., Strom S. E., Edwards S., Strom K. M., Shure M. A., 1990, AJ, 99, 1187

Stone J. M., Norman M. L., 1992, ApJS, 80, 753

Strom K. M., Strom S. E., Edwards S., Cabrit S., Skrutskie M. F., 1989, AJ, 97, 1451

Wyatt M. C., 2008, ARA&A, 46, 339

Xu Y., Mccray R., 1991, ApJ, 375, 190

Zhu Z., Nelson R. P., Dong R., Espaillat C., Hartmann L., 2012, ApJ, 755, 6

This paper has been typeset from a $\TeX$ file prepared by the author.