Hot Super Earths: disrupted young jupiters?

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ABSTRACT

Recent Kepler observations revealed an abundance of ‘hot’ Earth-size to Neptune-size planets in the inner 0.02–0.2 au from their parent stars. We propose that at least some of these smaller planets are the remnants of massive giant planets that migrated inwards quicker than they could contract. We show that such disruptions occur if the young giant planet embryo is initially extended. The characteristic planet–star separation at which such ‘hot disruptions’ occur is $R \approx 0.03–0.2$ au. The disruption is most likely at high accretion rates, $\dot{M} \gtrsim 10^{-6} \, M_\odot \, \text{yr}^{-1}$, when the migration is rapid and the embryo is unable to contract quickly enough. At late times, when the accretion rate drops to $\dot{M} \lesssim 10^{-7} \, M_\odot \, \text{yr}^{-1}$, the embryos migrate sufficiently slow to be not disrupted. These ‘late arrivals’ may explain the well-known population of hot jupiters, although they could also be accounted for by the more compact giant planets in the framework of the core accretion scenario for planet formation.

Key words: planets and satellites: formation – planet–star interactions – protoplanetary discs.

1 INTRODUCTION

Standard protoplanetary disc models (Chiang & Goldreich 1997) show that the inner $\sim 0.1$ au region is too hot to allow for the existence of small solid particles there. Thus, planets should not be able to grow there. Yet observations of nearby solar-type stars show that many of them do host planets in that inhospitable-to-planet-formation region. The very first exoplanet to be convincingly detected had the separation of $R \sim 0.05$ au from its parent star and had a mass of about that of Jupiter (Mayor & Queloz 1995). Such gas giant planets circling their parent stars in a close proximity ($R \lesssim 0.1$ au) are now called ‘hot jupiters’. It is believed that they are explained by the inward radial drift (migration) of the planets born farther out (Lin, Bodenheimer & Richardson 1996).

The Kepler mission has recently produced a number of surprising results (Borucki et al. 2011), one of which is that there is an even greater number of smaller planets, for example, Earth-size to Neptune-size planets, in that region. It is similarly obvious that these smaller ‘hot’ planets also had to be brought there from farther out by an inward radial migration. One practical difficulty in testing this idea, though, is that the migration of smaller planets is expected to occur in the poorly understood ‘type I’ regime (Paardekooper & Papaloizou 2008). This theoretical difficulty leads to a large range of uncertainties in the predictions of the detailed core accretion model calculations (e.g. see fig. 5 in Ida & Lin 2008) for smaller planets.

Recently, the key importance of the radial migration of the earliest gas condensations formed in the massive protoplanetary discs by the gravitational instability was realized (Vorobyov & Basu 2006; Boley et al. 2010). Analytical estimates (Nayakshin 2010a) and numerical simulations (Vorobyov & Basu 2006, 2010; Boley et al. 2010; Cha & Nayakshin 2011) show that these condensations can migrate all the way from their birthplace in the outer $R \sim 100$ au disc to the inner $\sim$ few au disc and be disrupted there during the earliest massive disc phase ($t \lesssim$ few $\times 10^6$ yr, typically). Recently, a detailed numerical investigation into the radial migration of massive planets in self-gravitating discs (Baruteau, Meru & Paardekooper 2011) confirmed the rapid migration seen by the authors quoted earlier. The underlying physical reason for the rapid migration was found to be that the planets do not open a gap and migrate in the faster type I regime.

It was pointed out by Boley et al. (2010) and Nayakshin (2010a) that this migration-and-disruption sequence yields an unexplored way of forming terrestrial-like planets. If dust grows and sediments in the centre of the clump and forms a solid density core there, then the tidal disruption of the gas clump may leave a solid core – an Earth-like protoplanet (note the connection to earlier ideas of McCrea & Williams 1965; Boss 1998; Boss, Wetherill & Haghipour 2002). Nayakshin (2010b, 2011b) used a simple spherically symmetric radiation hydrodynamic code with the dust grains as a second fluid to delineate the conditions when such a mechanism for solid core growth can work. Based on the potential promise of these ideas, Nayakshin (2010a) formulated the ‘Tidal Downsizing’ (TD; Nayakshin 2010a) hypothesis for planet formation. In this picture, a partial disruption of a $\sim 10 M_J$ gas clump [which we also call a giant embryo (GE)] leaves a giant planet, whereas a complete disruption yields a terrestrial-like planet.

In this paper, we continue to assess the potential utility of the TD hypothesis to planet formation. We point out that another...
ingredient, noted but not explicitly included in the calculations by Boley et al. (2010) and Nayakshin (2010a,b, 2011b), must be added to the scheme. To explain it, consider isolated GEs first. As they contract, their internal temperature increases. At early times, the rate of this contraction is controlled by the cooling rate of the embryo, that is, by the rate at which the embryo can get rid of the excess thermal energy. However, when the temperature $T_{2nd} \approx 2500$ K is reached, molecular hydrogen ($H_2$) dissociates. This process is an efficient energy sink, which allows the embryo to contract rapidly – in fact collapse hydrodynamically – without the need to radiate the energy away. The embryo collapse stops only at much higher densities, and temperatures as high as a few $10^5$ K, at which point hydrogen atoms are partially ionized. The embryo must then continue a slower contraction, again regulated by the rate at which its energy is radiated away. The collapse is known as the ‘second collapse’ in the star formation literature (Larson 1969), when the ‘first cores’ of masses $\sim 50 M_1$ collapse (Masunaga & Inutsuka 2000) to become the ‘second cores’, which are the proper protostars.

The second collapse may be the last step to making a gas giant planet in the gravitational disc instability (GI) model for giant planet formation (e.g. Boss 1997). However, in the TD hypothesis for planet formation, as we show below, this final step is not automatically successful – planets continuing to migrate rapidly towards their parent stars may still be disrupted at $R \sim 0.1$ au. We suggest this process as a way of forming the hot Super Earths observed by the Kepler mission (Borucki et al. 2011).

In analogy to the star formation literature, we refer to the GEs that are mainly molecular, and the embryo’s temperature $T_e < T_{2nd}$, as the ‘first GEs’; those where $H_2$ is disassociated are termed ‘second GEs’ instead.

2 CONTRACTION AND COLLAPSE OF THE FIRST EMBRYO

To illustrate the main point of this paper, we calculate the contraction of a giant embryo with ‘typical’ parameter values (e.g. those that appear quite reasonable to us for a solar metallicity disc around an $\sim$ solar mass star; see Nayakshin 2010a). In particular, the embryo mass $M_e = 100 M_1$ (this mass is chosen as the most likely mass of the gas clumps formed by GI in the disc; see below and also Boley et al. 2010), the normalized dust opacity $\kappa_0 = 0.5$ and the grain mass fraction $f_g = 0.01$. The embryo is initialized as a first core of the same mass (see Nayakshin 2010b).

The calculation is carried out with an updated version of the 1D gas-dust grains radiative hydrodynamics code of Nayakshin (2010b, 2011b). Instead of using an ideal gas equation of state with $\gamma = 5/3$, the code now uses the equation of state appropriate for a hydrogen and helium mix with the relative weight abundances $X = 0.7$ and $Y = 1 - X$ (the contribution of metals to the internal energy is neglected). The calculation includes dissociation of $H_2$, its rotational and vibrational degrees of freedom, with the ortho-hydrogen-to-para-hydrogen ratio fixed at 3:1 (cf. Boley et al. 2007), but neglects hydrogen and helium atom ionization. This is justified as the first embryos are too cold for the ionization effects to be important. Operationally, having advanced the internal specific energy, $u$, for every grid through a time-step, we then iterate to solve for the gas temperature and the abundances of molecular and atomic hydrogen corresponding to that value of $u$ and density $\rho$. The equations solved are identical to equations (30)–(41) and (46) and (47) of Nayakshin (2010b). Instead of using an ideal gas equation of state with $\gamma = 5/3$, the code now uses the equation of state appropriate for a hydrogen and helium mix with the relative weight abundances $X = 0.7$ and $Y = 1 - X$ (the contribution of metals to the internal energy is neglected). The calculation includes dissociation of $H_2$, its rotational and vibrational degrees of freedom, with the ortho-hydrogen-to-para-hydrogen ratio fixed at 3:1 (cf. Boley et al. 2007), but neglects hydrogen and helium atom ionization. This is justified as the first embryos are too cold for the ionization effects to be important. Operationally, having advanced the internal specific energy, $u$, for every grid through a time-step, we then iterate to solve for the gas temperature and the abundances of molecular and atomic hydrogen corresponding to that value of $u$ and density $\rho$. The equations solved are identical to equations (30)–(41) and (46) and (47) of Nayakshin (2010b).

Figure 1. The embryo’s temperature (in units of $10^3$ K), density (in $10^{-8}$ g cm$^{-3}$) and radial size, $r_e$ (in au), as a function of time, as labelled in the figure. Note the abrupt change near the end of the calculation, marking the second collapse, when the central temperature reaches $T \sim 2300$ K. The bump in the central temperature near $t = 1000$ yr is caused by the formation of an $\approx 5 M_\oplus$ solid core inside the embryo.

Viau, Bastien & Cha (2006) with the metal abundance $Z$ set to zero.

Fig. 1 shows the time-evolution of the embryo’s central temperature (solid line, in units of $10^3$ K), the gas density (dotted line, in units of $10^{-8}$ g cm$^{-3}$) and the outer radius of the embryo, $r_e$, (dashed line, in units of 1 au).

Despite the updated equation of state, the evolution of the first embryo is quite similar to that of the cases studied in Nayakshin (2011b). This may not be entirely surprising, given the similar insensitivity of the first (gas) cores to the equation of state as found by Masunaga, Miyama & Inutsuka (1998) and Masunaga & Inutsuka (2000). The embryo contracts and heats up, whereas dust grains grow. By the time $t \approx 1000$ yr, the grains increase in size to about 20 cm. Their density exceeds that of the gas in the centre of the embryo; they become self-gravitating and form a solid core of mass $M_c \approx 5 M_\oplus$. Fig. 1, the solid core formation is notable by the bump in the central temperature. After the core formation, the central region becomes hotter than the grain vaporization temperature of $\approx 1400$ K, evaporating the grains, and thus terminating further core growth (see Nayakshin 2011b, for details on this negative feedback loop). The central region also expands slightly. Most of the GEs are, however, unaffected by the solid core in this case, and the curves resume their otherwise monotonic behaviour a few hundred years later.

At $t \approx 6.5 \times 10^3$ yr, the GE goes through the second collapse when the central temperature exceeds about $2300$ K. The embryo radius drops rapidly, while the density and the temperature increase strongly. The first embryo becomes the second in our terminology. We note that the time until the collapse that we find here is rather close to $5.7 \times 10^3$ yr found for the same mass model by Helled & Bodenheimer (2010). It is possible that the formation of the solid core and the slight bounce of the inner gas cloud’s regions (cf. Nayakshin 2010a) associated with the energy release by the core may be responsible for the small difference, although the different opacities used may be a more important factor.
3 THE SECOND EMBRYO

3.1 Uncertainties in the structure of young planets

At a distance \( R \) from the star, the tidal density is

\[
\rho_i = \frac{M_e}{2\pi R^3} \approx 4 \times 10^{-3} \frac{M_e}{\text{M}_\odot} \left( \frac{0.03 \text{ au}}{R} \right)^3,
\]

(1)

where \( M_e \) is the stellar mass. Planets with mean densities below the tidal density are tidally disrupted, partially or wholly (in detail, this depends on the internal structure of the planet down to the core and the strength of the disc torques). Aged planets \( (t \gtrsim 10^9 \text{ yr}) \) have typical densities \( \rho \gtrsim 1 \text{ g cm}^{-3} \), and they are thus stable to tidal disruption except for ultrashort-period orbits or if irradiation and tidal heating effects are important in sweling the planet (e.g. Gu, Lin & Bodenheimer 2003).

However, the density of a young second embryo may be comparable to that given by equation (1) and they may therefore be unstable to the Roche lobe overflow, losing some or all of their gaseous mass before they mature into the present-day hot jupiters.

Unfortunately, our radiation hydrodynamics code is not well suited to follow the evolution of the second embryo, missing key high-density physics and stellar atmosphere modelling of radiative transfer (e.g. see Burrows et al. 2000; Baraffe et al. 2003). Compounding this, there is no consensus in the literature on the evolution of the giant planets younger than approximately a few Myr because of unknown initial conditions, computational difficulties, and also due to the dependence on the microphysics of opacity and grain sedimentation inside the planet (Boley & Durisen 2010; Boley, Helled & Payne 2011; Helled & Bodenheimer 2011).

The initial radius of the planet is very important as the initial density varies as the inverse third power of it. Starting from the planet of \( \sim 20R_J \) radius, Graboske et al. (1975) find that the embryo of mass \( M_e = 1M_J \) spends the initial contraction stages on the Hayashi track. During this phase, the planet’s effective temperature, \( T_{\text{eff}} \), is almost constant at \( \log T_{\text{eff}} \approx 3.1 - 3.3 \), and the planet’s radius evolves from \( \sim 20R_J \) to \( \sim 2.5R_J \) in \( 10^3 \text{ yr} \) (see their figs 1 and 2).

Bodenheimer (1974) and Decampli & Cameron (1979) estimate the \( 1M_J \) planet radius right after the hydrogen dissociation collapse to be \( \sim 5R_J \). Bodenheimer et al. (1980), on the other hand, find that hydrogen molecule dissociation results in extremely small planets, \( R \approx 1.3R_J \), without the Hayashi phase. At the end of their calculation, the central density of the planets is \( \gtrsim 0.1 \text{ g cm}^{-3} \).

We are mainly interested here in more massive giant embryos, \( 5 \lesssim M_e \lesssim 20M_J \). The lower mass limit is set by the fragmentation criteria of self-gravitating gas discs: the minimum mass of the gas fragment must exceed the opacity limit for gravitational fragmentation, \( M_{\text{min}} \approx (5-10)M_J \) (Low & Lynden-Bell 1976; Rees 1976; Masunaga & Inutsuka 1999, see also the first embryo mass estimates by Boley et al. 2010). The upper mass limit is set by the gas clump becoming hotter before grain sedimentation could occur (Nayakshin 2010b). These more massive clumps cool rapidly, go through the second collapse and become very massive coreless planets or brown dwarfs.

In principle, such objects, if contracting slowly enough in the post-second-collapse stage, could also be tidally disrupted in the inner disc region. The result of such a disruption is currently not clear as more detailed modelling is needed, but obviously a complete disruption would leave no trace of the planet/brown dwarf, whereas a partial disruption would leave a less massive gas object. These more massive embryos are therefore unlikely to have anything to do with the subgiant planets in the inner \( \sim 0.1 \text{ au} \) of the parent stars that we discuss in this paper.

However, even for the more massive embryos that we consider here, the uncertainty in the young age remains. The radii of giant isolated planets with mass between 1 and 10\( M_J \), as found in the literature, are \( \approx 2R_J \) at 1 Myr according to Nelson et al. (1986, fig. 1); the same but at a few Myr according to Burrows et al. (2000) and Bodenheimer et al. (2003); and 2–3\( R_J \) at 1 Myr according to Burrows et al. (2001, fig. 3) and Baraffe et al. (2003, fig. 6).

Addition of stellar irradiation could significantly puff up the early planets (e.g. fig. 1 in Burrows et al. 2000). In a stark contrast to these results, working within the framework of the core accretion model for planet formation, Marley et al. (2007) showed that if one assumes that accretion shock radiates away the accretion energy of the gas infalling on to the planet, then the resulting early radius and luminosity of the planet can be much smaller than the values quoted above, obtained under what the authors called ‘hot-start initial conditions’. We refer to such models as ‘low-density-start’ ones as these reflect more directly the importance of the planet initial density in determining whether it will or will not be disrupted. The radii of the planets at age \( t = 10^6 \text{ yr} \) are found by Marley et al. (2007) to be between 1.2 and 1.4\( R_J \), consistent with the value found by Bodenheimer et al. (1980) (which could be called a ‘high-density-start’ model).

We believe that low-density initial conditions are more relevant to the problem at hand because (i) our embryos are not isolated, thus irradiation and tidal heating effects are present; (ii) the embryo may be rapidly rotating before the second collapse (Boley et al. 2010; Nayakshin 2011a) and thus specific angular momentum conservation could slow down the contraction of the planets after \( H_2 \) dissociation; and (iii) Nayakshin (2011b) finds that the accretion luminosity of the solid core inside the giant embryo may be significant and may even unbind the gas envelope if the opacity is large enough. This extra heating source inside the embryo is likely to keep it hotter.

Nevertheless, keeping in mind these present uncertainties in the properties of the young massive gas giants, we consider below a simple model with a range of initial second embryo radii that we believe roughly brackets the possible outcomes of a planet migrating very close to its parent star.

3.2 A simple model of the second collapse

We estimate the radius of the embryo right after the \( H_2 \) dissociation using the energy conservation arguments (neglecting the helium contribution to the planet’s energy). The total energy of the embryo before dissociation is \( E_0 \approx -GM_e^2/2r_0 \), where \( r_0 \) is the radius of the embryo. The same quantity immediately after the collapse is

\[
E_0 = -\frac{GM_e^2}{2r_0} + \frac{M_e}{2m_H}D + \frac{X_e}{m_H} \frac{M_e}{m_H} \frac{1}{D} + \frac{X_e}{m_H} \frac{1}{D} \frac{1}{D},
\]

(2)

where \( D = 4.5eV \) and \( \epsilon = 13.6 \) are the dissociation energy of hydrogen molecules and the ionization energy of a hydrogen atom, respectively, and \( 0 \lesssim X_e \lesssim 1 \) is the number fraction of fully ionized hydrogen. Since \( |E_0| \) is much less than the absolute magnitudes of the first two terms on the right-hand side of this equation, we can approximate \( E_0 \approx 0 \). With this, we estimate the planet’s initial radius in the second configuration as

\[
r_c(t_2) \approx \frac{GM_e m_H}{D + 2X_e \epsilon} \approx 5.4 R_J \frac{M_e}{M_1} \left( \frac{1}{1 + 2X_e \epsilon/D} \right)^{0.4},
\]

(3)

where \( t_2 \) is the time of the second collapse. For \( X_e = 0 \), this equation reproduces the larger initial radii of a Jupiter-mass planet obtained by Bodenheimer (1974) and Decampli & Cameron (1979);
for $X_i \approx 0.4$, the smallest value of $1.3R_i$ (obtained by Bodenheimer et al. 1980) is recovered. For more massive embryos, the second collapse may in principle ionize even more hydrogen, so we shall also consider larger values of $X_i$ below.

Having estimated the initial radius of the embryo after H$_2$ dissociation, we need a cooling model to calculate the rate of planet’s contraction. The ‘hot-start’ (low-density-start, in our terminology) models typically predict effective temperatures for young (t $\sim$ 10$^6$ yr) planets in the range of 1000–2000 K, decreasing with time to smaller values. The high-density-start models of Marley et al. (2007), on the other hand, give $T_{\text{eff}}$ $\approx$ 500–800 K for the same time period. We shall make the simplest assumption that the effective temperature of the planet remains constant during the time of interest. [This is also consistent with the ‘Hayashi tracks’ for young and massive $M_e \gtrsim 10M_J$ planets found by Graboske et al. (1975).]

To sample the uncertain range of cooling rates, we consider three values of $T_{\text{eff}}$ = 500, 1000 and 2000 K. This translates into a difference in the cooling rates of a factor of $\sim$250 from the smallest to the largest $T_{\text{eff}}$ models, thus covering a broad region in the parameter space.

With this simple cooling model, we solve for the evolution of the second GE radius, $r_2(t)$, as a function of time
\[
\frac{\text{d}E_e}{\text{d}t} = -4\pi r_e^2 \sigma k T_{\text{eff}}^4, \tag{4}
\]
where the energy of the embryo is $E_e \approx GM_e^2/2r_e(t)$. A trivial integration yields
\[
r_e(t) = \frac{r_e^2(t)}{1 + Ar_e^2(t - t_2)}. \tag{5}
\]
Here $r_2$ (equation 3) is the planet’s radius at the time of the second collapse, $t_2$, and $A = 24\pi\sigma T_{\text{eff}}^4/(GM_e^2)$. This model does not take into account the electron degeneracy pressure in the GE, which becomes important at high densities, keeping the planet from further contraction when the planet cools and exhausts its thermal pressure support. Inclusion of the degeneracy pressure would thus strengthen our conclusions.

The second embryo’s mean density, $\rho_e = 3M_e/(4\pi r_e^3)$, as a function of time in our model, is
\[
\rho_e(t) = \rho_0(t_2) \frac{18\sigma T_{\text{eff}}^4}{GM_e}(t - t_2). \tag{6}
\]
In terms of absolute values, the initial density of the second embryo is at least $10^{-3}$ g cm$^{-3}$ and rises with time according to
\[
\rho_e(t) = \rho_0(t_2) + 5 \times 10^{-4} \text{ g cm}^{-3} \frac{(t - t_2) 10M_J}{10^4 \text{ yr}} \frac{T_{\text{eff}}^4}{M_e}, \tag{7}
\]
where $T_3 = T_{\text{eff}}/1000$ K. Comparing this with the tidal disruption density $\rho_1$ given by equation (1), we conclude that, to be disrupted, the giant planet needs to (i) have $\rho(t_2) < \rho_1$ and (ii) be also young, for example, $t \lesssim 10^6$ yr.

### 3.3 Tidal vulnerability of young ‘low-density-start’ planets

Before we consider the results of our exploratory model, we point out a limiting case which is useful in the interpretation of these results. We estimate the planet’s migration rate, assuming that there is only one giant planet in the inner $R \gtrsim$ few au at any one time. Due to its significant mass, the embryo migrates in the type II regime. As shown by Ivanov, Papaloizou & Polnarev (1999), the migration time, $t_{\text{mig}}$, is shorter than but comparable to the ‘accretion time’, $t_a = M_e/M$, where $M$ is the accretion rate in the disc. Thus,
\[
t_{\text{mig}} \lesssim \frac{M_e}{M} = 10^3 \text{ yr} \frac{M_e}{10M_J} \frac{10^{-6} M_\odot \text{ yr}^{-1}}{M} \tag{8}
\]
Now, for a low-density-start planet, the initial density $\rho(t_2)$ is much smaller than the second term in equation (7), and we can neglect it. In this limit, the location of the second (hot) disruption of the embryo depends only on the accretion rate through the disc (since this controls the planet’s migration rate). We can estimate the embryo–star separation at which the disruption occurs, $R_{\text{hot}}$, defined as the radius, where $\rho_e = \rho_0$. For this, we note that $(t - t_2)$ in equation (6) should be of the order of $t_{\text{mig}}$ as this is the time it takes for the embryo to migrate inwards.

Replacing $(t - t_2)$ with $t_{\text{mig}}$ in equation (6) and requiring the embryo density to equal the tidal disruption density, we arrive at
\[
R_{\text{hot}} = \left[ \frac{GM_eM}{36\pi\sigma T_{\text{eff}}^4} \right]^{1/3} = 0.041 \text{ au} \left( \frac{M}{10^{-6} M_\odot \text{ yr}^{-1}} \right)^{1/3} \tag{9}
\]

Rephrasing the statement made after equation (7), to be disrupted, the giant planet needs to start in the low-density configuration and also reside in a high-accretion-rate disc which would transport it into the hot $R < 0.1$ au region rapidly. Independently of the initial configuration, ‘late arrivals’, when $M \lesssim 10^{-8} M_\odot \text{ yr}^{-1}$, may survive and contract into hot jupiters as the disc runs out of mass, presumably due to photoevaporation (Alexander, Clarke & Pringle 2006).

### 4 ILLUSTRATIVE PLANET PLUS DISC MIGRATION CALCULATIONS

Nayakshin (2010a) used a toy analytical model to describe planets migrating in a protoplanetary accretion disc. In this simple model, the inner disc radial structure is calculated assuming the standard accretion disc model of Shakura & Sunyaev (1973) for a given mass of the star, $M_*$, the doubling time-scale for the star, $t_0$, and the disc viscosity parameter, $\alpha$. Where the self-gravity parameter $Q$ drops below unity (the outer $R \gtrsim$ tens of au), the disc structure is modified using the additional constraint, $Q = 1$, for the marginal self-gravitational stability of the disc (cf. Levin 2007). At $t = t_0$, the assembly of the star is assumed complete (its mass reaches 1 $M_\odot$), and the disc torques are abruptly removed. Here we use this toy disc migration model to exemplify the possible outcomes of young massive embryos migrating in the disc. We use the model embryos calculated as described in Sections 2 and 3.

For all of the representative cases presented below, the embryo’s mass is $M_e = 10M_J$, born at $R = 100$ au and initialized as described in Nayakshin (2010b). For simplicity, we assume that the mass of the giant embryo is constant until it is disrupted dynamically. 3D simulations of embryos migrating in gas discs show the destruction of the embryos in a matter of two to three orbits (Boley et al. 2010; Cha & Nayakshin 2011) once they fill their Roche lobes.

#### 4.1 An example of a disrupted hot jupiter

We first present one particular low-density-start calculation, setting $X_i = 0$ in equation (3) and $T_{\text{eff}} = 1000$ K. The parameters of the calculation are $M_e = 0.5M_\odot$, $t_0 = 10^5$ yr and $\alpha = 0.05$. This yields an accretion rate through the disc of $\dot{M} = M_e/t_0 = 5 \times 10^{-6} M_\odot \text{ yr}^{-1}$.

The upper panel of Fig. 2 shows the radial location of the embryo as a function of time. The embryo migrates into the inner $\sim 0.1$ au region of the disc in less than $10^4$ yr due to the high disc accretion rate. The middle panel shows the evolution of the Hills radius, $r_H$, and the embryo’s radius, $r_e$. $r_e$ is initially one order of magnitude smaller than the Hills radius. In fact, this difference increases to
a factor of around 50 when the embryo goes through the second collapse at $t \sim 7 \times 10^3$ yr.

The lower panel shows the embryo’s temperature evolution. The solid core formation occurs at $t \sim 10^3$ yr (cf. Fig. 1), marked by the vertical dashed line. The ‘hot’ disruption of the embryo occurs at time $t \approx 3 \times 10^4$ yr, at the embryo–star separation of $R = 0.07$ au. Note that the disruption of the gaseous envelope should leave a solid core behind, since the solid core formation occurred much earlier. In our calculation, the solid core migrates slightly to $R \sim 0.055$ au by $t = t_{db}$. This last bit of radial migration is via type I and may in reality be far less efficient (Ida & Lin 2008; Paardekooper & Papaloizou 2008).

### 4.2 The parameter space for the disruption of young planets

We now vary several parameters of our toy model to investigate how the outcome varies. In particular, we pick three values of $X_i = 0, 0.3$ and 0.85, which translates into the difference in the initial $r_e(t_e)$ by a factor of 6, and also three values of $T_{\text{eff}} = 500, 1000$ and 2000 K. While not being self-consistent, these simple model giant planet $r_e(t)$ functions cover the parameter space from the low-density-start to the high-density-start models. In addition, we vary the doubling time-scale for the star, $t_{db}$, which is equivalent to varying the accretion rate through the disc, $\dot{M}$.

We arbitrarily assume that the protoplanetary disc inner boundary is at $R = 0.02$ au, which may well be reasonable for the high accretion rates we are interested in. Planets that reach this inner disc radius intact are expected to either remain there or be swallowed by the protostar if it extends to this radius.

Fig. 3 shows the models with $T_{\text{eff}} = 1000$ K. The upper panel shows the location of the tidal disruption of the planet, if it occurs. The solid curve corresponds to the low-density-start model with $X_i = 0$, whereas the red dotted and the blue dashed curves show the cases for $X_i = 0.3$ and 0.85, respectively (the dot–dashed curve is explained below). The lower panel of the figure shows the disruption time, if the disruption occurs, or infinity, if it does not.

For the highest accretion rates, $\dot{M} \gtrsim 10^{-5} M_\odot$ yr$^{-1}$, the embryo migrates very quickly and therefore undergoes ‘cold’ disruption at $R \sim 3–4$ au, as in the base model of Nayakshin (2010a). The embryo at that point is still dominated by molecular hydrogen. Since the solid core formation occurs at $t \sim 10^3$ yr, the disruption should leave the core behind. This outcome is common to all the three curves as the disruption occurs before the embryo gets to the second stage.

At smaller accretion rates, $\dot{M} \lesssim 10^{-5} M_\odot$ yr$^{-1}$, the result depends crucially on how compact the embryo is right after the $\mathrm{H}_2$ dissociation collapse. For the most extended planets, for example, the solid curve in the figure, the disruption occurs for accretion rates between $\sim 10^{-3}$ and $\sim 10^{-2} M_\odot$ yr$^{-1}$. The disruption radius drops as $\dot{M}$ drops because the planet migrates slower and is able to contract further by the time it is disrupted. At the lowest accretion rates considered, the planet cools sufficiently to avoid tidal disruption altogether. As explained above, it either stalls there or is swallowed by the protostar.

The black dot–dashed curve shows the analytical approximation to the ‘hot’ disruption radius, $R_{\text{hot}}$, given by equation (9). It is seen to be a very good approximation to the solid curve in the appropriate parameter range, which implies that for the lowest initial density model, the particular initial density of the planet, $\rho_e(t_e)$, is quickly forgotten (cf. equation 6).

The mid-range model, $X_i = 0.3$, shown with the red dotted curve, shows a similar pattern, but here the disruption location varies less. This is readily explicable by the larger initial density $\rho_e(t_e)$, so that the cooling of the planet in the initial $\sim 10^3–10^4$ yr after the second collapse does not actually compact the embryo significantly. Finally, the blue dashed curves show that the most compact models, $X_i = 0.85$, do not get disrupted at all after they have undergone the second hydrodynamical collapse.

We repeated these simple calculations for the two other values of the planet effective temperature, $T_{\text{eff}} = 500$ and 2000 K, and plot the disruption radius in Figs 4 (a) and (b), respectively. The slowest
Disrupted hot jupiters

The radial location of the tidal disruption of the giant embryo as a function of accretion rate in the disc for $T_{\text{eff}} = 1000 \text{ K}$ and three values of the hydrogen ionization fraction $X_i$ (see equation 3), $X_i = 0$, 0.3 and 0.85, shown with the solid black, dotted red and dashed blue curves, respectively. The dot-dashed curve shows the analytical approximation to the ‘hot disruption’ radius (equation 9). The lower panel shows the disruption time for the same three models. At the highest accretion rates, the embryos migrate so rapidly that they are disrupted while their hydrogen is still molecular (‘the cold disruption’ at $3–4 \text{ au}$). At lower accretion rates, the hot disruption occurs for the hot-star model but not for the high-density-start one. At the lowest accretion rates, $\dot{M} \lesssim 10^{-7} \text{ M}_\odot \text{ yr}^{-1}$, tidal disruption does not occur for any model as the planet contracts sufficiently rapidly and never fills its Roche lobe.

Figure 3. Upper panel: the radial location of the tidal disruption of the giant embryo as a function of accretion rate in the disc for $T_{\text{eff}} = 500 \text{ K}$ and two other values of the effective temperature, as labelled in the panels. As expected, slowly cooling models (the upper panel) are much more prone to tidal disruption in the inner disc than the faster cooling ones (the lower panel).

dependencies connected with the microphysics of these planets, such as dust sedimentation and opacity, the mass and the luminosity of the solid core (if present). The global environment of the forming planet is also important since this may result in a rapid rotation of the planet and/or strong stellar irradiation.

Therefore, we chose to consider a simple model with a range of values for the planet’s radius after the second collapse and a range of effective temperatures. Varying the parameters of this simple model, we cover the parameter space from the ‘cold-start’ to the ‘hot-start’ planets. With the help of the analytical toy model for planet migration from Nayakshin (2010a), we find that (i) cold-start planets are unlikely to be tidally disrupted at any value for the disc accretion rate; (ii) hot-start models are disrupted in the inner $\sim 0.02–0.1 \text{ au}$ region, provided that their Kelvin–Helmholtz contraction time-scale is longer than $\sim 10^3–10^5 \text{ yr}$, depending on the disc accretion rate; and (iii) at the lowest accretion rates, for example, below $\dot{M} \sim 10^{-7} \text{ M}_\odot \text{ yr}^{-1}$, the disruption becomes increasingly unlikely.

We should emphasize that in the context of the TD hypothesis, all planets except the most massive ones, $M \gtrsim 5 \text{ M}_J$, should have undergone at least some envelope removal. The initial mass of gas clumps formed by the gravitational fragmentation of the outer disc is expected to be at least $\sim 5–10 \text{ M}_J$ (Boley et al. 2010; Nayakshin 2010b). Hence, even the hot jupiters of mass $\sim \text{ M}_J$ should have lost most of their gas envelopes. This could have occurred at large $\sim \text{ au}$ to tens of au distances if the embryo filled its Roche lobe while hydrogen was still molecular (as discussed by Boley et al. 2010;
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