REGULATED STAR FORMATION IN FORMING DISK GALAXIES EXPOSED TO THE ULTRAVIOLET RADIATION BACKGROUND

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ABSTRACT

We perform radiation hydrodynamics simulations on the evolution of galactic gas disks irradiated by the ultraviolet radiation background. We find that gas disks with $N_{HI} \geq 10^{21} \text{cm}^{-2}$ that are exposed to ultraviolet radiation at a level of $I_\lambda = 1$ can be self-shielded from photoheating, whereas disks with $N_{HI} \leq 10^{21} \text{cm}^{-2}$ cannot. We also find that the unshielded disks keep a smooth density distribution without any sign of fragmentation, while the self-shielded disks easily fragment into small pieces through self-gravity, possibly followed by star formation. The suppression of star formation in unshielded disks is different from the photoevaporation effect, since the assumed dark halo potential is deep enough to retain the photoheated gas. The presence of such a critical threshold column density would be one of the reasons for the so-called downsizing feature of present-day galaxies.

Subject headings: galaxies: formation — hydrodynamics — radiative transfer

1. INTRODUCTION

Recent advances of available computational resources have enabled us to simulate the entire disk galaxy in order to investigate parsec-scale structures in the disk (Wada & Norman 2007; Robertson & Kravtsov 2008; Tasker & Bryan 2008; Saio et al. 2008). In addition, numerical techniques for radiation hydrodynamics (RHD) simulations are now being developed, especially in the field related to cosmic reionization (Gnedin & Abel 2001; Melema et al. 2006; Rijkhorst et al. 2006; Susa 2006; Iliev et al. 2006; Yoshida et al. 2007; Qiu et al. 2007; Whalen & Norman 2008; Altay et al. 2008). Using these schemes, we are now ready to tackle the numerical simulations of the detailed structures in galactic disks coupled with radiation transfer.

The importance of ultraviolet radiation transfer effects on the galaxy formation/star formation in galaxies has been pointed out by several authors. First of all, the observed star formation threshold in galactic disks (Kennicutt 1989; Martin & Kennicutt 2001) could be explained by the self-shielding effects of galactic disks from the external ultraviolet radiation field (Schaye 2001, 2004). Schaye (2004) obtained the physical state of the disks exposed to the external ultraviolet field, using one-dimensional calculations with Cloudy (Ferland 2000). On the basis of the linear stability arguments (Toomre 1964), he has demonstrated that the inner parts of the galactic disks shielded against an external radiation field by dust absorption could be gravitationally unstable, whereas the unshielded outer parts of the disks are stable. These results nicely explain the observed features of galactic disks; however, they have not been tested by multidimensional RHD simulations.

Secondly, in the context of galaxy formation, radiative feedback by the ultraviolet background is expected to be very important, especially for low-mass galaxies. Photoheating can heat the gas up to $10^4 \text{K}$, which prevents the gas from collapsing, in case the gravitational potential of the dark matter halo is not deep enough to retain the photoheated gas. Such feedback effects are quoted as photoevaporation, which is well studied at various levels (Umemura & Ikeuchi 1984; Efstathiou 1992; Babul & Rees 1992; Thoul & Weinberg 1996; Barkana & Loeb 1999; Ferrara & Tolstoy 2000; Gnedin 2000; Kitayama et al. 2000, 2001; Susa & Umemura 2004a, 2004b). In addition, there is some evidence in dwarf galaxies that the star formation is suppressed in the “hot” ($\sim 10^4 \text{K}$) phase (Young & Lo 1997a, 1997b), from which we infer that photoheating can suppress star formation in galaxies, even if the potential of the dark halo is deep enough to prevent the gas from evaporating. This issue could be relevant to the so-called downsizing problem in nearby galaxies (Cowie et al. 1996; Kauffmann et al. 2003). These authors found that old galaxies are massive, whereas the young ones are less massive. This trend continues to higher redshifts (e.g., Kodama et al. 2004; Erb et al. 2006; Reddy et al. 2006; Papovich et al. 2006). One of the possible interpretations of this feature is that star formation proceeds very rapidly in massive galaxies, whereas it is a slow process in less massive galaxies for some reason. Kauffmann et al. (2003) also pointed out that these two groups are well defined; i.e., they have a clear boundary in stellar mass at $\sim 10^{10} M_\odot$. This critical mass scale is too large to be related to the photoevaporation mechanism; however, photoheating could still be the candidate with which to explain the downsizing mass if we can show the mechanism by which the star formation is suppressed even in halos in which the photoionized gas can be inherent.

In this paper, we examine the fragmentation of gas disks embedded in a dark halo potential under the ultraviolet radiation field, using the recently developed RHD code RSPH (Susa 2006). This paper is organized as follows. In the next section, the numerical scheme is briefly summarized. In §3, the setup of our numerical simulations is described. We show the results of our numerical simulations, as well as the analytic estimate, in §4. Sections 5 and 6 are devoted to discussion and conclusions.

2. METHODOLOGY

We perform numerical simulations with the radiation SPH (RSPH) code (Susa 2006). The code can compute the fractions of the primordial chemical species $e^-$, $H^+$, $H^-$, $H_2$, and $H_2^+$ by fully implicit time integration. It also can deal with multiple sources of ionizing radiation, as well as the radiation at the Lyman-Werner band.

The hydrodynamics is calculated with the smoothed particle hydrodynamics (SPH) method. We use the version of SPH from Umemura (1993), with the modification to the SPH kernel function.
and the symmetrization of the equations of motion according to Steinmetz & Müller (1993). We also adopt the particle resizing formalism by Thacker et al. (2000), in which the number of neighbor SPH particles is kept almost constant, without sudden changes.

The nonequilibrium chemistry and radiative cooling for primordial gases are calculated by the code developed by Susa & Kitayama (2000), in which the H$_2$ cooling and reaction rates are mostly taken from Galli & Palla (1998). As for the photoionization process, we employ the so-called on-the-spot approximation (Spitzer 1978). In this paper, we also added the radiative cooling due to metals, employing the formula given in Dalgarno & McCray (1972) and assuming that $Z = 10^{-2} Z_{\odot}$. The metallicity of nearby disk galaxies is normally larger than $Z = 10^{-1} Z_{\odot}$; however, the observed metallicity of the intergalactic medium, from which the galaxies are formed, is $Z \sim 10^{-2}$ to $10^{-3} Z_{\odot}$ (e.g., Songaila 2001; Schaye et al. 2003). Since we are interested in the forming disk galaxies, we employ a value of $Z = 10^{-2} Z_{\odot}$ in the present set of simulations. Note that the radiative transitions of heavy elements are not the dominant process of radiative cooling for $Z \lesssim 1000$ K (Bohringer & Hensler 1989; Susa & Umemura 2004a) in the case in which $Z = 10^{-2} Z_{\odot}$ is assumed. H and H$_2$ molecules are the main coolants of such gas.

The optical depth is integrated by using the neighbor lists of SPH particles. In our old scheme (Susa & Umemura 2004a), we created many grid points on the light ray between the radiation source and an SPH particle. In the present scheme, we do not create so many grid points. Instead, we create one grid point on the light ray for each SPH particle on its upwind. We find the neighboring “upstream” particle for each SPH particle on its line of sight to the source, which corresponds to the grid point. Then the optical depth from the source to the SPH particle is obtained by summing up the optical depth at the “upstream” particle and the differential optical depth between the two particles. This process is described in more detail in Susa (2006). We assess the optical depth for the ionizing photons, as well as the Lyman-Werner band photons, by the method described above. In the present version of the code, dust opacity is not included in the calculation.

The code has been tested for various standard RHD problems (Susa 2006) and has already been applied to the issues of the radiative feedback effects of first-generation stars (Susa & Umemura 2006; Susa 2007). We also took part in a code comparison project with other radiation hydrodynamics codes (Iliev et al. 2006), in which we found reasonable agreements between the codes.

### TABLE 1

| Model | $I_{21}$ | $\rho_i$ (M$_\odot$ pc$^{-3}$) | Simulated Time (Myr) | Number of SPH Particles | $n_{H_2}$ (cm$^{-3}$) | $\epsilon$ (pc) |
|-------|---------|-----------------|-----------------|-----------------|-----------------|----------------|
| A……… | 0 | 0.05 | 300 | $1.28 \times 10^6$ | 235 | 3.05 |
| B……… | 0 | 0.1 | 120 | $2.56 \times 10^6$ | 235 | 3.05 |
| C……… | 0 | 0.3 | 40 | $7.68 \times 10^6$ | 235 | 3.05 |
| Ar………. | 1 | 0.05 | 350 | $1.28 \times 10^6$ | 235 | 3.05 |
| Br/2……. | 0 | 0.1 | 120 | $2.56 \times 10^6$ | 235 | 3.05 |
| Cr/2……… | 1 | 0.3 | 40 | $7.68 \times 10^6$ | 235 | 3.05 |
| B/8……… | 0 | 0.1 | 120 | $2.56 \times 10^6$ | 58.8 | 6.09 |
| Br/2……… | 0 | 0.1 | 120 | $2.56 \times 10^6$ | 58.8 | 6.09 |

3. SETUP OF NUMERICAL SIMULATIONS

We perform numerical simulations of a galactic gas disk embedded in a dark halo potential. The dark halo potential is fixed as

$$\Phi_{DH}(r) = \frac{1}{2} \left[ \frac{a v_1^2}{r^2 + a_1^2} + \frac{a v_2^2}{r^2 + a_2^2} \right]^{1/2},$$

where $r$ denotes the distance from the center, $a = 1$ kpc, $a_1 = 0.3$ kpc, $a_2 = 5$ kpc, and $v_1 = v_2 = 100$ km s$^{-1}$. Thus, the assumed potential is similar to the one employed by Wada & Norman (2007), except that the rotation velocity at a given radius is smaller by a factor of 2, since we are interested in forming galaxies that are relatively less massive than our Galaxy.

In the present set of simulations, we assume an initially uniform disk (Fig. 1) with slight perturbations (the displacements of the SPH particles are $\pm 10\%$). The initial thickness and radius of the disks, $H_i$ and $R_{disk}$, are 100 pc and 3 kpc for all of the runs, respectively. The initial velocity of the gas particles in the disk is assumed to be the Kepler rotation velocity around the dark halo potential center. Using equation (1), we find that the rotation velocity at the edge of the disk, $r = R_{disk}$, can be as large as $\sim 100$ km s$^{-1}$. Thus, the dark halo potential is deep enough to retain the gas even if it is photoheated to $\sim 10^4$ K. The initial densities of the disks, $\rho_i$, are changed depending on the models, which are listed in Table 1.

The light rays of the external ultraviolet radiation field are assumed to be perpendicular to the gas disk. We consider two directions (upward and downward), as shown in Figure 1. This particular choice of the ray direction is an approximation if we consider the background radiation field exactly, since background photons are coming from all angles. If we account for such effects, the boundary between the shielded and unshielded regions will be softened, which might change the self-shielding effects slightly. However, as long as we do not change the number of photons coming into the disk, the self-shielding condition is also almost unchanged, since it is determined by the balance between the number of recombining hydrogen atoms and the number of photons. The flux of the field at the Lyman limit is given such that the mean intensity at the midplane of the disk is equal to $10^{-21} I_{21}$ ergs s$^{-1}$ cm$^{-2}$ Hz$^{-1}$ sr$^{-1}$ in the case in which the opacity of the gas disk is ignored, where $I_{21}$ is the intensity normalized by $10^{-21}$ in cgs units. We consider an AGN-type spectrum, which is proportional to $\nu^{-1}$.

For all of the runs, we employ the common physical/numerical parameters of a gas disk that are listed in Table 2, whereas the model-dependent parameters are listed in Table 1. As shown in
Table 1, we perform 10 runs, changing the density of the disk and the intensity of the radiation field from run to run.

Since the resolution of the simulation is constrained by the mass of an SPH particle, there is a critical density above which the gravitational fragmentation is not captured properly. According to Bate & Burkert (1997), the density $n_{\text{H, res}}$ is given as

$$n_{\text{H, res}} = \frac{3}{4\pi m_p} \left( \frac{5k_B T_{\text{min}}}{2Gm_p} \right)^{\frac{3}{2}} \left( \frac{1}{2m_{\text{SPH}} n_{\text{H, res}}} \right)^{\frac{1}{2}},$$

(2)

where $m_{\text{SPH}}$, $n_{\text{H, res}}$, and $T_{\text{min}}$ denote the mass of each SPH particle, the number of neighbor particles, and the minimal temperature set in the simulations, respectively. The symbols $k_B$, $G$, and $m_p$ have their ordinary meanings; i.e., they represent the Boltzmann constant, the gravitational constant, and the proton mass, respectively. In the present set of simulations, the value of $n_{\text{H, res}}$ is equal to 235 cm$^{-3}$, except in the runs B/2, Br/2, B/8, and Br/8, in which the masses of the SPH particles are larger than those in the other regular runs by a factor of 2 or 8 (see Table 1). The softening length of the gravitational interaction between SPH particles, $\epsilon$, is set so as to satisfy

$$\epsilon = \frac{1}{2} \left( \frac{3N_{\text{nei}} m_{\text{SPH}}}{4\pi m_p n_{\text{H, res}}} \right)^{1/3}.$$  

(3)

This expression guarantees that the self-gravity of a dense clump above $n_{\text{H, res}}$ will be softened.

Given these configurations, we can assess the Toomre $Q$-value in order to understand the stability of the initially uniform disk. Toomre's $Q$ is defined as

$$Q = \frac{c_s \kappa}{\pi G \Sigma}.$$  

(4)

Here $c_s$, $\kappa$, and $\Sigma$ denote the sound speed of the gas, the epicyclic frequency, and the surface density of the disk, respectively. In the case in which $Q > 1$ holds, the disk is gravitationally stable because of the Coriolis force/thermal pressure, whereas it is unstable if $Q < 1$ is satisfied (Toomre 1964).

The actual $Q$ parameters for models A, B, and C (see Table 1) are plotted in Figure 2. Here $c_s$ is assumed to be the sound speed corresponding to $T_{\text{min}} = 300$ K. Since we assume uniform disks for all models, the disks are more unstable at the outer radii. In the present simulations, the radius of the disks is 3 kpc for all models. Thus, there exist critical radii above which the disks are unstable if they are cooled down to $T_{\text{min}}$, although the unstable region is narrow for model A (it is unstable only at the edge of the uniform disk). It is worth noting that if we employ values of $T_{\text{min}}$ that are lower than 300 K, the unstable region should expand, since the gas temperature will go down to $T_{\text{min}}$, which results in smaller values of $Q$.

4. RESULTS

4.1. Stability of the Disks without Radiation

First we present the results for $I_{z1} = 0$; i.e., no external radiation field. The face-on/edge-on view of the snapshots for the disks in models A, B, and C are shown in Figure 3. The snapshots are taken at $t = 300$ Myr for model A, $t = 120$ Myr for model B, and $t = 40$ Myr for model C. The particular choices of these times when the snapshots are taken basically corresponds to the times when the time steps at dense gas clumps collapsed...
below $n_{H,\text{res}}$ become so short ($\leq 10^3$ yr) that physical time in the numerical computation evolves very slowly. Since the growth timescale of gravitational instability is shorter in denser disks, we take an earlier snapshot in model C than in the others.

It is clear that the disks in all models fragment into small filaments/knots. At the same time, the inner parts of the disks are stable. As is shown in Figure 2, the $Q$-values are larger than unity at smaller radii, whereas they are smaller than unity at larger radii. Therefore, the inner parts of the disks are stable, which is consistent with the present results. In addition, the critical radii outside of which the disks are unstable as predicted by the Toomre criterion are also consistent with our numerical results. Indeed, the critical radii read from Figure 2 are 2.8, 1.6, and 0.8 kpc for models A, B, and C, which are almost consistent with the boundary radii of the smooth parts of the disks.

In Figure 4, the phase diagrams (in the $n_H$-$T$ plane) of the three runs at the same times as in Figure 3 are shown. The gas temperature in all of the models is cooled efficiently, which is almost close to $T_{\text{min}} = 300$ K in the high-density realm ($n_H \gtrsim 10^3$ cm$^{-3}$). Such efficient cooling justifies the stability argument based on the $Q$-values plotted in Figure 2, in which $T = T_{\text{min}}$ is assumed.

These phase diagrams also show that very dense regions above $n_{H,\text{res}}$ are formed in all of the models. These dense regions correspond to the fragments presented in Figure 3.

4.2. Effects of Ultraviolet Radiation on the Disk Stability

Now we show the snapshots of the runs with $I_{21} = 1$ (models Ar, Br, and Cr) in Figure 5; i.e., $t = 300$ Myr for model Ar, $t = 120$ Myr for model Br, and $t = 40$ Myr for model Cr. The
snapshots are synchronized to the corresponding ones in Figure 3. At a first glance, we find clear differences from the runs without ultraviolet radiation. In model Cr, we find small filaments and knots, as were found in model C, whereas we cannot find any sign of fragmentation in models Ar and Br. In addition, the density of the disks in models Ar and Br becomes lower than that in models A and B. We also observe that the disks are geometrically thicker than the previous ones. In Figure 6, the phase diagrams are shown for models Ar, Br, and Cr. In all of the runs, a significant amount of material is photoheated up to $\geq 10^4$ K in the low-density region, which contributes to the disk thickening. In addition, no very dense gas component above $n_{\text{H}_2}\text{cm}^{-3}$ is found in models Ar and Br, while we find it again in model Cr, as was found in model C. Thus, these results indicate that the presence of the external ultraviolet radiation field suppresses the fragmentation of the disks for our low-density models (Ar and Br), although it does not for our high-density model (Cr).

4.3. Self-shielding from Ultraviolet Radiation

The difference between our low-density models (Ar and Br) and our high-density model (Cr) comes from the self-shielding effects. Figure 7 shows the density (top), temperature (middle), and H $i$ fraction $\chi_{\text{H}_i}$ (bottom) distributions along the direction perpendicular to the disks. In model Cr (right), the gas near the midplane (at $z = 0$) is neutral, cold, and dense, because of the self-shielding. On the contrary, in models Ar and Br, the cold and dense regions near the midplane are relatively smaller than that in model Cr. To be more quantitative, Figure 8 shows the temperature probability distribution functions (PDFs) for models Ar, Br, and Cr (top) and A, B, and C (bottom). The temperature

Fig. 5.—Same as Fig. 3, but for models Ar, Br, and Cr.

Fig. 6.—Same as Fig. 4, but for models Ar, Br, and Cr.
PDFs show the mass fraction of the disk found within a logarithmic temperature bin of width $\Delta$ (log $T$) = 0.06. In the runs without the radiation field (models A, B, and C), most of the mass is condensed in the coldest phase at $T_{\text{min}}$. On the other hand, the temperature PDFs are broader in the runs with radiation. In models Ar and Br, most of the mass in the disks is in the hot phase ($>10^4$ K), whereas the cold phase ($<10^3$ K) is dominant in model Cr. Such a large difference in the temperature distribution results in a different stability of the disks. According to Figure 2, the $Q$-values for models Ar and Br are $\geq 0.5$ if $r \leq 3$ kpc and $T = T_{\text{min}}$. In the actual simulations, we find that $T \sim 10^4$ K, which means that the $Q$-values are larger than 0.5 by a factor of $(10^4/300)^{1/2} = 5.77$; i.e., $Q > 1$ is achieved everywhere in the disks for these models. Thus, these disks are stable, whereas the disk in model Cr is unstable because of its coldness.

A rough estimate of the self-shielding criterion for primordial gas has been derived by Susa & Umemura (2000a), who evaluate the shielded photoheating rate in order to compare it with the peak $H_2$ cooling rate below $10^4$ K. We remark that heavy elements are not the dominant coolant for $10^3$ K < $T$ < $10^4$ K in the case in which $Z \leq 10^{-2} Z_\odot$ (Susa & Umemura 2004a).

A similar argument is also made in Corbelli et al. (1997), which describes the thermal instability of the primordial gas. According to Susa & Umemura (2000a), the photoheating rate per unit volume at the midplane of the disk for the optically thick limit is given as $n_{H_\alpha} y_{H_\alpha} \mathcal{H}$, where

$$\mathcal{H} = \frac{4 \pi I_\alpha n_{H_\alpha} \sigma_{H_\alpha} \Gamma(\beta)}{3(1 + \beta) \tau_{\lambda^{-\beta}}}.$$  

Here $n_{H_\alpha}$ and $y_{H_\alpha}$ are the number density of the hydrogen nuclei and the $H_\alpha$ fraction at the midplane of the disk, $\sigma_{H_\alpha}$ is the photoionization cross section at the Lyman limit, $I_\alpha$ is the Lyman limit frequency, $I_\alpha$ denotes the incident ultraviolet intensity at the Lyman limit, and $\Gamma(\beta)$ is the gamma function. The quantity $\beta$ used in this equation is defined as $\beta \equiv 1 + (\alpha - 1)/3$, where $\alpha$ denotes the spectral index; i.e., $I_\nu \propto \nu^{-\alpha}$ is assumed. The quantity $\tau_{\lambda^{-\beta}}$ denotes the optical depth at the Lyman limit, which is written as

$$\tau_{\lambda^{-\beta}} = \frac{\langle y_{H_\alpha} \rangle N_{HI} \sigma_{H_\alpha}}{2}.$$  

Here $N_{HI}$ denotes the column density of the disk and $\langle y_{H_\alpha} \rangle$ is the $H_\alpha$ fraction averaged along the direction perpendicular to the disk.

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2 The expression $\mathcal{H}$ is 2 times larger than eq. (A12) in Susa & Umemura (2000a), since the radiation field irradiates both the right side of the disk and the reverse side in the present simulations.
disk. The peak H$_2$ cooling rate per unit volume below 10$^4$ K is described as $n^2_{H_2}C_1$, where $C_1$ $\approx$ 10$^{-26}$ erg s$^{-1}$ cm$^{-3}$ (Shapiro & Kang 1987; Susa & Umemura 2004a). In addition, we assume hydrostatic equilibrium in the vertical direction. As a result, the number density of hydrogen nuclei at the midplane of the disk is related to the column density as follows:

$$n_{H_2} = \frac{\pi G n^2_{H_2} m_p}{2 c^2_s}.$$  

(7)

Here the gas disk is assumed to be isothermal, with a sound velocity of $c_s$. Letting the photoheating rate be equal to the H$_2$ peak cooling rate, and combining that with equations (5), (6), and (7), we obtain the threshold column density $N_{H_{\text{sh}}}$ above which the photoheating is shielded enough for the midplane to cool down to $\leq$1000 K:

$$N_{H_{\text{sh}}} = \left[ \frac{2^{b+1} c^2_s I_{\text{sh}} \nu_{\alpha} \sigma_{\alpha}^{-\beta} \Gamma(\beta)}{3G \mu m_p C_1 \left< \gamma H_2 \right> (\gamma + 1)} \right]^{1/(b+1)}.$$  

(8)

Using the present assumptions ($\alpha = 1$, $I_{\text{sh}} = 1$), we have

$$N_{H_{\text{sh}}} \approx \left(2.5 \times 10^{21} \text{ cm}^{-2} \right) \left( \frac{I_{\text{sh}}}{1} \right)^{1/3} \left( \frac{c_s}{10 \text{ km s}^{-1}} \right)^{2/3} \left( \frac{\left< \gamma H_2 \right>}{0.5} \right)^{-1/3}.$$  

(9)

Here we also assume that $\gamma_{H_{\text{sh}}} = 1$ in the above expression, since it is true for all three of the models Ar, Br, and Cr.

On the other hand, the initial column density of the disks in models Ar, Br, and Cr is 6.3 $\times$ 10$^{20}$, 1.3 $\times$ 10$^{21}$, and 3.8 $\times$ 10$^{21}$ cm$^{-2}$, respectively. Thus, the numerical simulations indicate that 1.3 $\times$ 10$^{21}$ cm$^{-2} \leq N_{H_{\text{sh}}} \leq$ 3.8 $\times$ 10$^{21}$ cm$^{-2}$, which almost agrees with the rough analytic estimate shown above.

4.4. Probability Distribution Functions of the Density Field

The differences among the six models (A, Ar, Br, B, C, and Cr) are clarified in Figure 9 in terms of their density probability distribution functions. Six panels represent the snapshots of the density PDFs for these six models. Here the density PDFs are defined as the mass fraction found within a logarithmic density bin of width $\Delta(\log n)$ = 0.18. The histograms drawn with thin lines represent the initial PDFs, whereas those drawn with thick lines show the evolved ones.

It is clear that in models A, B, C, and Cr, the PDFs extend beyond the resolution limit $n_{H_{\text{res}}}$, although the gravitational force is softened for $n_{H_{\text{res}}} \lesssim n_{H_{\text{res}}}$ (see eq. [3]). In fact, even if gravitational force could be neglected at $n_{H_{\text{res}}} \approx n_{H_{\text{res}}}$, matter accretes from the outer envelope. As a result, the density (or pressure) of the clump becomes higher than $n_{H_{\text{res}}}$, in order to compensate for the increasing ram/thermal pressure of the accreting material.

On the other hand, in models Ar and Br, the PDFs have a very sharp cutoff below $n_{H_{\text{res}}}$, which could be interpreted to say that the star formation activity is suppressed beyond the resolution limit $n_{H_{\text{res}}}$. Thus, dense self-gravitating fragments formed in these disks, which would lead to star formation activity.

4.5. Convergence of the PDF

The differences among the six models Ar, Br, B, C, and Cr are clarified in Figure 9 in terms of their density probability distribution functions. Six panels represent the snapshots of the density PDFs for these six models. Here the density PDFs are defined as the mass fraction found within a logarithmic density bin of width $\Delta(\log n)$ = 0.18. The histograms drawn with thin lines represent the initial PDFs, whereas those drawn with thick lines show the evolved ones.

It is clear that in models A, B, C, and Cr, the PDFs extend beyond the resolution limit $n_{H_{\text{res}}}$. Thus, dense self-gravitating fragments formed in these disks, which would lead to star formation activity.

It looks strange at first that very dense clumps form beyond $n_{H_{\text{res}}}$, although the gravitational force is softened for $n_{H_{\text{res}}} \approx n_{H_{\text{res}}}$ (see eq. [3]). In fact, even if gravitational force could be neglected at $n_{H_{\text{res}}} \approx n_{H_{\text{res}}}$, matter accretes from the outer envelope. As a result, the density (or pressure) of the clump becomes higher than $n_{H_{\text{res}}}$, in order to compensate for the increasing ram/thermal pressure of the accreting material.

On the other hand, in models Ar and Br, the PDFs have a very sharp cutoff below $n_{H_{\text{res}}}$ for the snapshots at $t$ = 300 Myr (model Ar) and $t$ = 120 Myr (model Br), although models A and B have a clear fragmentation signature at the same times. Thus, the impression from the montage of the disk is correct; i.e., the disk in models Ar and Br does not fragment into dense clouds. This result could be interpreted to say that the star formation in these disks is heavily suppressed by the ultraviolet radiation field.
Fig. 9.—Density PDFs for models Ar, Br, Cr, A, B, and C. Horizontal axes show the density, whereas the vertical axes show the differential mass fraction $\Delta M/M_{\text{total}}$ found within a logarithmic density bin of width $\Delta (\log n_H) = 0.18$. The histograms drawn with thin solid lines indicate the initial PDFs in the models. Those drawn with thick solid lines indicate the evolved PDFs at $t = 300$ Myr for models A and Ar, $t = 120$ Myr for models B and Br, and $t = 40$ Myr for models C and Cr. The thin dotted vertical lines show the resolution limits of the runs.

Fig. 10.—Comparison of the density PDFs in models B and Br to the results from the low-resolution runs (B/2, Br/2, B/8, and Br/8). The axes are the same as in Fig. 9. Top: PDFs for models Br (black histograms), Br/2 (red histograms), and Br/8 (blue histograms) at $t = 40$ Myr (left), 100 Myr (middle), and 120 Myr (right). Density resolution limits, $n_{H,\text{res}}$, corresponding to the runs Br, Br/2, and Br/8 are shown as the vertical dotted lines with colors that are the same as the histograms. Bottom: PDFs for models B, B/2, and B/8 at the same three epochs. The colors are the same as in the top panels.
as B/2, Br/2, B/8, and Br/8 (Table 1). Figure 10 shows the density PDFs for models B, B/2, and B/8 (bottom) and models Br, Br/2, and Br/8 (top) at t = 40, 100, and 120 Myr. In the half-resolution runs (B/2 and Br/2), the density resolution limits are 4 times smaller than those in models B and Br, because $n_{\text{H, res}}$ is inversely proportional to the square of the mass resolution (see eq. [2]). Similarly, the density resolution limits in runs B/8 and Br/8 are 64 times smaller than those in runs B and Br. The PDF histograms basically agree very well below $n_{\text{H, res}}$, as we expected.

In runs B/8 and Br/8, the density resolution limits are so low that the high-density regions ($\gtrsim 100$ cm$^{-2}$) are not captured properly. As a result, it seems to be difficult to distinguish the PDFs for models B/8 and Br/8 from each other. On the other hand, the results from runs B/2 and Br/2 indicate the same conclusion as we found in runs B and Br, since the density resolution is sufficient to capture the self-shielded dense regions. Thus, the disk in run B/2 fragments because of gravitational instability, whereas that in run Br/2 is stable. Therefore, we conclude that the resolutions of the present regular simulations are sufficient to capture the fragmentation of the disks under the assumptions we employed, whereas in lower resolution runs such as B/8 and Br/8, we are not able to describe the present physical situation.

5. DISCUSSION

It should be emphasized that the present results directly prove the suppression of star formation activities by ultraviolet background radiation in halos with $T_{\text{dr}} > 10^4$ K. It has already been established that star formation is suppressed in less massive halos with $T_{\text{dr}} \lesssim 10^4$ K, since the gas in the dark halo potential evaporates because of the high thermal pressure of the photoheated gas (Umemura & Ikeuchi 1984; Efstathiou 1992; Babul & Rees 1992; Thoul & Weinberg 1996; Barkana & Loeb 1999; Ferrara & Tolstoy 2000; Gneden 2000) if the self-shielding effect is not important (Kitayama et al. 2000, 2001; Susa & Umemura 2004a, 2004b). On the other hand, in the present simulations, the gas in the dark halo does not evaporate because of the deep gravitational potential. We find that even in such halos, star formation could be heavily suppressed in the case in which the gas is configured to form disks with low column density.

Another important result found in the present calculation is the presence of a clear boundary in the column density of the disk, below which star formation is heavily suppressed. According to the numerical results, the critical column density is $N_{\text{H}} \sim 1-4 \times 10^{21}$ cm$^{-2}$. This threshold is very interesting, mainly for two reasons. First, although the present simulations are performed with a fixed dark halo mass, we can try to convert the critical column density into the dark halo mass. The mass of the uniform disk with given values of $N_{\text{H}}$ and disk radius $R_{\text{disk}}$ is

$$M_{\text{disk}} = \left(6 \times 10^9 \, M_\odot\right) \left(\frac{R_{\text{disk}}}{10 \, \text{kpc}}\right)^2 \frac{N_{\text{H}}}{2.5 \times 10^{21} \, \text{cm}^{-2}}.$$  \hspace{1cm} (10)

Therefore, if the mass of the host dark halo is 7 ($\gtrsim \Omega_b/\Lambda$) times the disk mass, the critical dark halo mass is $\sim 4 \times 10^{10} \, M_\odot$. This value is still smaller than the critical scale found by Kauffmann et al. (2003) by a factor of a few, but if we take the internal feedback effects such as UV radiation from internal sources (AGNs, massive stars) or supernovae into account, the threshold might account for the observed critical downsizing mass. We also point out that if we start the simulation from a cosmological setup, the threshold mass could be raised more, since photoheating might be able to penetrate deeper into the disk because of the initially less dense configurations. Thus, including such physics, as well as starting simulations from a cosmological setup, will be necessary in order to obtain a more precise understanding of star formation in forming disk galaxies.

Secondly, the critical column density is as large as the star formation threshold found in local disk galaxies (Kennicutt 1989; Martin & Kennicutt 2001). In fact, the observed star formation threshold of galactic disks has been investigated by Schaye (2001, 2004) from the theoretical side. He suggested that the galactic disk could be gravitationally stable if it is not shielded from the external UV radiation, since the $Q$-value of the photoheated disk easily exceeds unity. As a result, gravitational fragmentation of the disk is suppressed, so the star formation activities are also suppressed. Schaye & Dalla Vecchia (2008) performed simple hydrodynamics simulations that take the star formation threshold column density into account, using the effective equation of state in a multiphase medium, although they do not solve the radiation transfer equations explicitly. They found that the threshold column density can be as large as $4 \, M_\odot$ pc$^{-2}$ (see Figs. 4 and 5 in Schaye & Dalla Vecchia 2008), which is $5 \times 10^{20}$ cm$^{-2}$ in ergs units. This value is smaller than that obtained in this paper by a factor of a few ($\sim 2.5 \times 10^{20}$ cm$^{-2}$). The basic mechanism that they propose to suppress the star formation of the disk is same as the one found in the present simulations, except that we do not include the effects of dust extinction, which is especially important for present-day disk galaxies. The gas disk is more easily self-shielded by dust absorption, since the dust opacity at the Lyman limit frequency is as large as the H$\beta$ continuum for solar metallicity. Thus, it is reasonable that we have a larger critical column density than that obtained in Schaye & Dalla Vecchia (2008). On the other hand, in the present simulations, we assume that $Z = 10^{-2} \, Z_\odot$. Therefore, the dust opacity at the Lyman limit is smaller than the H$\beta$ continuum opacity by 2 orders of magnitude. As a result, dust opacity has a much smaller impact on the self-shielding effects at such a low metallicity. In any case, we dare to mention that our present results almost succeeded in probing the presence of the star formation threshold column density, utilizing the full three-dimensional radiation hydrodynamics simulations. More realistic calculations for present-day galaxies that include the effects of dust extinction are left for future works.

It is also worth noting that the star formation rate (SFR) of the disk above the threshold column density is consistent with the observed value, although it is difficult to compare the present results directly with the observed SFRs, since we do not take the local stellar feedback, the effects of dust particles, or the metal enrichment followed by the radiative cooling by abundant metals into account. Despite such issues, the SFR can be obtained with the density PDF, if we assume that the stars are formed in dense clumps within a local free-fall time. This is given as

$$\dot{\Sigma}_* = \epsilon_* \sqrt{G \rho_c \Sigma_{\text{disk}} f_c},$$  \hspace{1cm} (11)

where $\epsilon_*$ is the star formation efficiency, $\rho_c$ is the threshold density above which the gas is converted to stars, $\Sigma_{\text{disk}}$ denotes the surface density of the disk, and $f_c$ denotes the mass fraction of the gas in the disk that has condensed into dense clumps with $\rho > \rho_c$, which can be obtained by integrating the density PDF above $\rho_c$. In the present set of simulations with radiative feedback, model Cr is the only one in which star formation is expected, since self-shielded cold fragments emerge in the run. If we use $\epsilon = 0.1$ and $n_{\text{H, res}}$ to assess the threshold density $\rho_c$, we obtain a value of $f_c \approx 0.066$. Using the surface density of the disk in model Cr, we evaluate the SFR as $\dot{\Sigma}_* = 0.03 \, M_\odot \, \text{yr}^{-1} \, \text{kpc}^{-2}$. Although we cannot discuss the dependence of the SFR on $\Sigma_{\text{disk}}$.
the value is consistent with observations (Kennicutt 1998). This reasonable agreement infers the validity of the present numerical models.

We also point out another perspective on this issue from Susa & Umemura (2000a), in which they suggested that self-shielding from the ultraviolet background could be a key mechanism with which to determine the morphology of galaxies (Susa & Umemura 2000b), although their arguments are based on one-dimensional radiation hydrodynamics calculations. Unfortunately, it is impossible to relate the present results to morphology bifurcation of galaxies, since we assume a disk by hand in our simulations performed so far. From this point of view, again we need to perform simulations from cosmological initial conditions.

6. CONCLUSION

In this paper, we perform radiation hydrodynamics simulations on the fragmentation of galactic disks under the ultraviolet radiation background. We find that the ultraviolet radiation field strongly suppresses star formation in the disks in the case in which the photoheating is not shielded enough. We emphasize that this suppression is different from the photoevaporation effect, because the rotation velocities at the outer boundary of the disks in the present set of simulations are ~100 km s^{-1}, which is fast enough to retain the photoheated gas. In our simulations, we find a threshold column density of the disk (~10^{21} cm^{-2}) above which the fragmentation is not suppressed. This is similar to the star formation threshold column density observed in nearby galaxies. The presence of such a critical threshold would be one of the reasons for the so-called downsizing problem in nearby galaxies.

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