

Heating of the intracluster medium by quasar outflows

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ABSTRACT

We study the possibility of quasar outflows in clusters and groups of galaxies heating the intracluster gas in order to explain the recent observation of excess entropy in this gas. We use the extended Press–Schechter formalism to estimate the number of quasars that become members of a group or cluster of a given mass and formation epoch. We also estimate the fraction of mechanical energy in the outflows that is imparted to the surrounding medium as a function of the density and temperature of this gas. We finally calculate the total amount of non-gravitational heating from such outflows as a function of the cluster potential and formation epoch. We show that outflows from broad absorption line and radio-loud quasars can provide the required amount of heating of the intracluster gas. We find that in this scenario most of the heating takes place at $z \sim 1-4$, and that this ‘pre-heating’ epoch is at lower redshift for lower-mass clusters.

Key words: intergalactic medium – quasars: general – cosmology: theory – X-rays: galaxies.

1 INTRODUCTION

Clusters and groups of galaxies contain a large amount of hot gas, besides galaxies and the gravitationally dominant dark matter. This hot X-ray emitting gas, known as the intracluster medium (ICM), represents a part of the baryonic matter of the Universe that is not associated with individual galaxies but remains trapped in the deeper gravitational potential of galaxy clusters. Hierarchical models of structure formation have been very successful in explaining many observed properties of galaxies and galaxy clusters. Nevertheless, some puzzling problems remain open and unexplained. Models of cluster formation in which the intergalactic gas simply falls into the dark matter dominated gravitational potential well (so-called infall models) fail to reproduce all the structural properties of the local cluster population (e.g., Evrard & Henry 1991; Navarro, Frenk & White 1995; Mohr & Evrard 1997; Bryan & Norman 1998). There is certainly some additional physics driving ICM evolution.

Recent X-ray observations have provided evidence for some non-gravitational heating of the diffuse, high-density baryons in the potential wells of groups and clusters of galaxies, in addition to the heating during the gravitational collapse. One of the first pieces of evidence was in the shape of the L_x-T relation, which is steeper than the self-similar behaviour $L_x \propto T^2$ predicted in the case of gravitational processes only. As early as the emergence of *ROSAT* and *Einstein* data, several authors proposed that the missing element is the existence of a ‘pre-heated high-entropy’ intergalactic gas prior to a cluster’s collapse (David, Forman & Jones 1991; Evrard & Henry 1991; Kaiser 1991; White 1991). Later,

Ponman, Cannon & Navarro (1999) and Lloyd-Davies, Ponman & Cannon (2000) found direct evidence of an entropy excess with respect to the level expected from gravitational heating in the centre of groups. Ignoring the constant and logarithms, one can define the ‘entropy’ as $S \equiv T/n^{2/3}$. The excess entropy, or equivalently the excess specific energy, flattens the density profile, decreasing the X-ray luminosity, which is proportional to the square of the density. The effect is stronger in poorer clusters, where the excess energy associated with the excess entropy is comparable to the gravitational binding energy, while rich clusters, where gravity is dominant, are mostly unaffected. This produces a steepening of the L_x-T relation.

The most popular scenario to explain these thermal properties of the ICM successfully has been the ‘pre-heating’ scenario. For this scenario, the candidate processes that have been looked into are strong galactic winds driven by supernovae. However, Valageas & Silk (1999) showed that the energy provided by supernovae cannot raise the entropy of the intergalactic medium (IGM) up to the level required by current observations. The observed amount of required energy injection depends on the epoch, and values given have been in the range of 0.4–3 keV per gas particle (Navarro et al. 1995; Cavaliere, Menci & Tozzi 1997; Balogh, Babul & Patton 1999; Wu, Fabian & Nulsen 2000; Lloyd-Davies et al. 2000; Borgani et al. 2001). For example, Wu et al. (2000) showed that galactic winds can impart only ≤ 0.1 keV per particle. Moreover, Kravtsov & Yepes (2000) estimated the energy provided by supernovae from the observed metal abundance of the ICM and found that heating only by supernovae-driven outflows requires unrealistically high efficiency. On the other hand, quasar outflows may be much more powerful and plausible candidates for the heating (Valageas & Silk

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1999). In this paper, we focus on the role of quasar outflows in this regard.

The epoch of the energy input also remains uncertain. Lloyd-Davies et al. (2000) put an upper limit of $z_{\max} \sim 7-10$ on the pre-heating epoch, from their estimate of excess entropy in groups. For active galactic nuclei (AGNs), there have been no additional constraints like the metal abundance in the case of supernova heating (Kaiser & Alexander 1999). Recently, Yamada & Fujita (2001) have looked into the deformation of the cosmic microwave background (the Sunyaev–Zel’dovich effect) by hot electrons produced at the shocks produced by jets from AGNs. They showed that the observed excess entropy of the ICM and the *COBE/FIRAS* upper limit for the Compton y parameter are compatible with each other only when the heating by jets occurred at relatively small redshift ($z \leq 3$). Thus they questioned the ‘pre-heating’ scenario as their result suggests that the heating occurred simultaneously with or after cluster formation.

In this paper, we calculate the heat input from quasar outflows inside clusters. We calculate the mechanical work done by various kinds of quasar outflows and the excess energy imparted by them on to the intracluster medium via $p dV$ work. For the statistics of quasars inside clusters, we use the extended Press–Schechter formalism. Finally we calculate the excess energy per particle and tally the results with available observations.

In the next section, we discuss the abundance of quasars inside clusters of a given mass. We then discuss the evolution of quasar outflows and the mechanical work done by them in Section 3. We use these concepts to calculate the heating of the ICM in Section 4. We then discuss the implications of our results in Section 5. Throughout the paper we assume a flat universe with a cosmological constant, with $\Omega_0 = 0.35$, $\Omega_{\Lambda 0} = 0.65$ and $h = 0.65$.

2 QUASARS INSIDE CLUSTERS

For a proper evaluation of the heat input from quasars (QSOs) inside clusters, one first needs to calculate their abundance and its dependence on the quasar mass, cluster mass and the cluster formation redshift. Observationally, it is still difficult to obtain good statistics of quasars inside clusters. Estimations of the galaxy–QSO correlation function have shown that at low redshifts ($z \lesssim 0.4$) quasars typically reside in small to moderate groups of galaxies and not in rich galaxies (e.g., Bahcall & Chokshi 1991; Fisher et al. 1996). It is still uncertain whether or not their optical or radio luminosities depend on their environments. Some studies (e.g., Ellingson, Green & Yee 1991) have shown that radio-loud quasars preferentially reside in richer clusters at higher redshift, although more recent studies (e.g., Wold et al. 2001) do not find any such correlation.

There has been, however, considerable work in relating the observed quasar luminosity function or the radio luminosity function (for radio-loud quasars) with the mass function of galaxies as prescribed by the Press–Schechter (PS) formalism (e.g., Haehnelt & Rees 1993; Haiman & Loeb 1998; Yamada, Sugiyama & Silk 1999).

At $z = 0$, Yamada et al. (1999) found that one can reproduce the abundance of the radio sources powered by AGNs (i.e., leaving aside the radio sources powered by starbursts) by assuming that a fraction f_r of the haloes from the PS formalism become radio-loud quasars, where $f_r \sim 0.01$ for $M_h \gtrsim 10^{12} M_\odot$, and $f_r = 0$ for $M < 10^{12} M_\odot$. They assume an upper limit on M_h of $10^{14} M_\odot$. Since it is known that radio-loud quasars constitute a fraction 0.1 of the quasar population (Stern et al. 2000), this means that a fraction

$f_q = 10f_r$ of the haloes from the PS formalism become quasars. In other words, $f_q \sim 0.1$ for $10^{12} \leq M_h \leq 10^{14} M_\odot$, and $f_q = 0$ otherwise. Yamada et al. (1999) assumed this fraction to be a constant for all redshifts. In this model, the rate of formation of quasars is given by the derivative of the PS mass function at the relevant mass-scale.

This is similar to the model adopted by Haiman & Loeb (1998) and Furlanetto & Loeb (2001, hereafter FL01). FL01 showed that at high redshift ($z \gtrsim 4$) the rate of formation of quasars is a fraction $f_q \sim 0.1$ of the rate of formation of haloes from the PS formalism, if a lifetime of order 10^7 yr is assumed for the quasar. There is, however, a difference, in that they had only a lower limit to M_h and no upper limits [see equations (2) and (3) of Haiman & Loeb (1998)]. They mention that at low redshifts their formalism does not predict any decline as observed in reality, and for which reason they consider their model only at high redshift. It is possible that this mismatch is due to the lack of upper limits, since the differentiation of PS mass function for objects with an upper limit in mass decreases at low redshift (Haiman, private communications). In any case, Haiman & Loeb (1998) found that this prescription yields a matching quasar luminosity function that is observed at redshift $z \gtrsim 2.5$. We find later that most of the heating of the ICM gas (even for our least massive cluster) occurs at $z \gtrsim 2$ (Fig. 7). We will therefore assume for simplicity that this fraction $f_q \sim 0.1$ at all redshifts (as in Yamada et al. 1999).

In this paper, we would like to have a conservative estimate of the quasar abundance. Also, we would like to calculate the abundance of quasars in clusters including low-mass groups of galaxies. For this reason, we assume the value of f_q as above, but use an upper limit of $10^{13} M_\odot$. Since the mass function decreases steeply at the higher-mass end, this should not change the value of f_q substantially. In brief, we assume that,

$$f_q \sim \begin{cases} 0.1 & \text{if } 10^{13} \gtrsim M_h \gtrsim 10^{12} M_\odot, \\ 0 & \text{if } M_h > 10^{13} M_\odot, \quad M_h < 10^{12} M_\odot, \end{cases} \quad (1)$$

motivated by the model of Yamada et al. (1999), and by the fact that a similar prescription by Haiman & Loeb (1998) recovers the quasar population at high redshift, and relate the rate of formation of quasars with that of haloes in the PS formalism.

We are, however, concerned with the statistics of quasars *inside* clusters. For this one needs to have an extension of the PS mass function which can predict the probability of a given halo becoming a part of a bigger object later, or the probability of an object having had a progenitor of a given mass at an earlier epoch. Such extensions of the PS theory have been studied in detail by Bower (1991) and Lacey & Cole (1993), for example.

In the standard PS theory, the mass function, i.e., the fraction of regions with mass in the range $M, M + dM$ that have overdensity δ in excess of δ_c (which is the threshold for perturbations becoming non-linear), is given by

$$f_{\text{PS}} = f(M, \delta_c) dM = \frac{-1}{\sqrt{2\pi} (\sigma_M^2)^{3/2}} \exp\left(-\frac{\delta_c^2}{2\sigma_M^2}\right) \frac{d\sigma_M^2}{dM} dM. \quad (2)$$

Here, σ_M is the mass variance of the perturbation at the mass-scale M . The relation between the number density of objects in the mass range $M, M + dM$ with $f(M) dM$ is

$$n(M) dM = \frac{\rho_0}{M} f(M) dM, \quad (3)$$

where ρ_0 denotes the background mass density.

In the extended PS theory, the fraction of regions of mass M ,

contained within a larger scale region of mass M' and overdensity δ' , which are more overdense than δ_c , is given by

$$f(M, \delta_c | M', \delta') dM = \frac{-1}{\sqrt{2\pi} (\sigma_M^2 - \sigma_{M'}^2)^{3/2}} \times \exp \left[-\frac{(\delta_c - \delta')^2}{2(\sigma_M^2 - \sigma_{M'}^2)} \right] \frac{d\sigma_M^2}{dM} dM. \quad (4)$$

This expression recovers the simple PS mass function in the limit $M' \rightarrow \infty$ and $\delta' \rightarrow 0$, relevant for the whole universe.

If we then identify M' with M_{cl} , the mass of a cluster, and $\delta' = \delta_c(z_f)$, the threshold overdensity of the cluster at its formation epoch z_f , we can then obtain the mass fraction of M_{cl} that have been parts of progenitors of a given mass range (M to $M + dM$) at a given (earlier) redshift. If we also identify this mass range with that of the quasars as in the standard PS theory (10^{12} – $10^{13} M_\odot$), and use the fraction f_q of these haloes that become quasars (equation 1), we will obtain the mass fraction of the final cluster that have been quasars at some given earlier epoch.

To obtain the rate of formation of these quasars (inside a future cluster of mass M_{cl}), we should differentiate the above expression. A simple differentiation will, however, not give the correct result, since there will be a negative contribution from the merging of haloes out of this mass range. One would get a negative rate of formation of such quasars at some point if the rate at which they disappear beyond this mass limit is not taken into account in a proper manner.

Consider the abundance of objects in a given mass range M , $M + dM$ at two successive epochs z_1 and z_2 ($z_2 < z_1$). The abundance $f(M) dM$ at z_2 will be given by $f_2 = f_1 + F - D$, where f_2 and f_1 are the abundances at epochs z_2 and z_1 , the term F denotes the abundance of newly formed objects in this mass range during the epoch z_1 and z_2 (from merger of smaller objects), and D signifies the abundance of objects that moved out of this mass range as a result of merger (into bigger objects). A simple differentiation of the PS function, involving the difference ($f_2 - f_1$) will therefore depend on both F and D . For a given range of mass, D is very small at a very early epoch, but it increases with time (see, e.g., fig. 5 of Haiman & Menou 2000), and at a later epoch can become larger than F . Therefore, at lower redshifts, a simple differentiation can imply a negative rate of change of abundance. If the contribution of D is neglected, one would then incorrectly get a negative value of F .

This problem has been encountered in the case of the ordinary PS function by many authors. While studying the rate of mergers in the context of background radiation from starbursts and AGNs, Blain & Longair (1993) noted that a simple differentiation of the PS function leads to a negative rate of formation of objects in a given mass range. In other words, the actual rate of formation of objects is given by

$$\dot{f}_{\text{form}} = \dot{f}_{\text{PS}} + \dot{f}_{\text{merger}}. \quad (5)$$

They performed a simulation assuming a simple power spectrum and obtained a fit for the rate of these objects merging to form bigger objects. They found that \dot{f}_{merger} can be approximated well (in the Einstein–de Sitter universe) by

$$\dot{f}_{\text{merger}} = \phi \frac{f_{\text{PS}}}{t} \exp \left[(1 - \alpha) \frac{\delta_c^2}{2\sigma^2} \right], \quad (6)$$

where the value of $\alpha \sim 1.35$ and $\phi \sim 1.3$ – 1.7 . This problem has also been investigated by Sasaki (1994) and Percival & Miller

(1999). With the extension of the PS formalism, one can now calculate this merging rate (see also, Chiu & Ostriker 2000).

We note here that in the case of the merger of a lower-mass quasar into a more massive quasar, our implicit assumption is that the central black holes also merge and form a bigger black hole appropriate for the bigger quasar (see below; equation 13). In other words, we assume that the central black hole mass always traces the halo mass.

The rate at which an object of mass M at epoch z merges to form a bigger object of mass M' is given by (Bower 1991; Lacey & Cole 1993)

$$\frac{d^2 p}{dM' dz} (M \rightarrow M' | z) dM' = \frac{1}{\sqrt{2\pi} \sigma_{M'}^2 (\sigma_M^2 - \sigma_{M'}^2)^{3/2}} \times \exp \left[-\frac{\delta_c(z)^2 (\sigma_M^2 - \sigma_{M'}^2)}{2\sigma_M^2 \sigma_{M'}^2} \right] \times \left| \frac{d\sigma_M^2}{dM} \right| \left| \frac{d\delta_c(z)}{dz} \right| dM. \quad (7)$$

Here, $p(M \rightarrow M' | z) dM'$ is the probability of an object of mass M merging to become an object of mass within the range $M', M' + dM'$ at redshift z . The rate of disappearance of objects of a given mass, \dot{n}_{merger} , should be essentially

$$\frac{df_{\text{merger}}}{dz} (M, z) = f_{\text{PS}}(M, z) \int_{2M}^{\infty} \frac{d^2 p}{dM' dz} (M \rightarrow M' | z) dM', \quad (8)$$

where the lower limit of the integration is chosen to be such that the merged object is at least twice as massive as the merging object. We show the result of this for an sCDM universe (with a *COBE* normalized power spectrum) in Fig. 1, Fig. 1, and show the fit of Blain & Longair (1993) with $\phi = 0.9$ and $\alpha = 1.35$. We find that

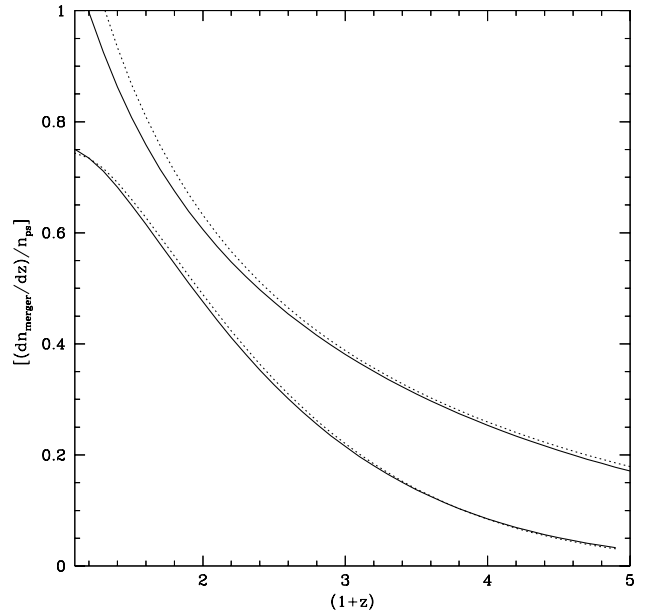


Figure 1. The rate of disappearance of objects compared in different cosmologies. The thin solid line shows the result of integration in equation (8) for an sCDM universe for $M = 10^{14} M_\odot$ and the dotted line shows the Blain & Longair (1993) fit with $\phi = 0.9$ and $\alpha = 1.35$. The thick solid line shows the result of the integration in the Λ CDM universe and the dashed line refers to a fit described in the text. Both curves use four-year *COBE* normalized spectra.

the merging rate is fitted by a lower value of ϕ than they assumed, although the difference is a factor of order unity. It is possible that this difference is due to the specific assumption in the simulation done by Blain & Longair (1993), e.g., the power spectrum being a simple power law (Blain, private communication), or it could be the result of the lower limit ($2M$) chosen by us. At any rate, if we chose the above integral to represent the merger rate then it would be a conservative estimate, since decreasing the lower limit would simply increase the value of the merger rate and, in turn, the formation rate of objects.

We have found that the result in the case of the Λ CDM universe can be fitted by a similar function, with ϕ being replaced by $0.9 d\delta_c(z)/dz$, for $\alpha = 1.35$ with an accuracy of order $\lesssim 5$ per cent. However, we do not use these fits in our calculation, and we evaluate the integral numerically for our purpose.

To be precise, this rate of disappearance is valid for the objects following the PS mass function, i.e., for objects that are not already parts of bigger objects. Motivated by the extension of the PS formalism, we here posit that the rate of disappearance of objects inside a bigger object also has the same form, with $f_{\text{PS}}(M, \delta_c)$ in equation (8) being replaced by $f(M, \delta_c | M', \delta')$. There is admittedly no way of verifying the truth of this Ansatz at present, since this would involve more extensions of the PS theory than presently available. It will also involve comparing the merger rates inside and outside of clusters. However, it leads to a conservative

estimate for the formation rate of quasars in a cluster. As the work of Bower (1991) has shown, growth of perturbations *inside* a cluster is enhanced compared to in the field. This means that the merging rate of objects of a given mass inside a cluster should be larger than that in the field. Here, by assuming a comparable merging probability (the factor that multiplies the abundance of objects, given by f_{PS}), we are in a way underestimating the rate of disappearance (\dot{f}_{merger}) and, in turn, the rate of formation (\dot{f}_{form}) of quasars in a cluster. The final result of total heat input from our formalism should, therefore, be a conservative estimate.

We show the results of adding the rate of disappearance in Fig. 2. The dashed lines, which show the term $-(df_{\text{PS}}/dz)$ (or equivalently df_{PS}/dt), become negative at lower redshift, suggesting the need for the addition of the rate of disappearance of objects. The dotted lines show the result of adding this rate, using the Blain & Longair (1993) fit (with $\phi = 0.9$), and the solid lines show the result using the integral in equation (8). The upper panels of the figure show the case of objects (of mass $10^{13} M_{\odot}$) in the field, and the bottom panels show the case for these objects inside a cluster of $M_{\text{cl}} = 10^{15} M_{\odot}$ ($z_f = 0$). For these curves, we have used the integration in equation (8), and the extension of the PS mass function for the abundance of objects inside the cluster, as explained above. The left and right panels show the cases for sCDM and Λ CDM universes. As the bottom panels show, the addition of the disappearance rate does not suffice to make $-(df_{\text{PS}}/dz)$ (inside

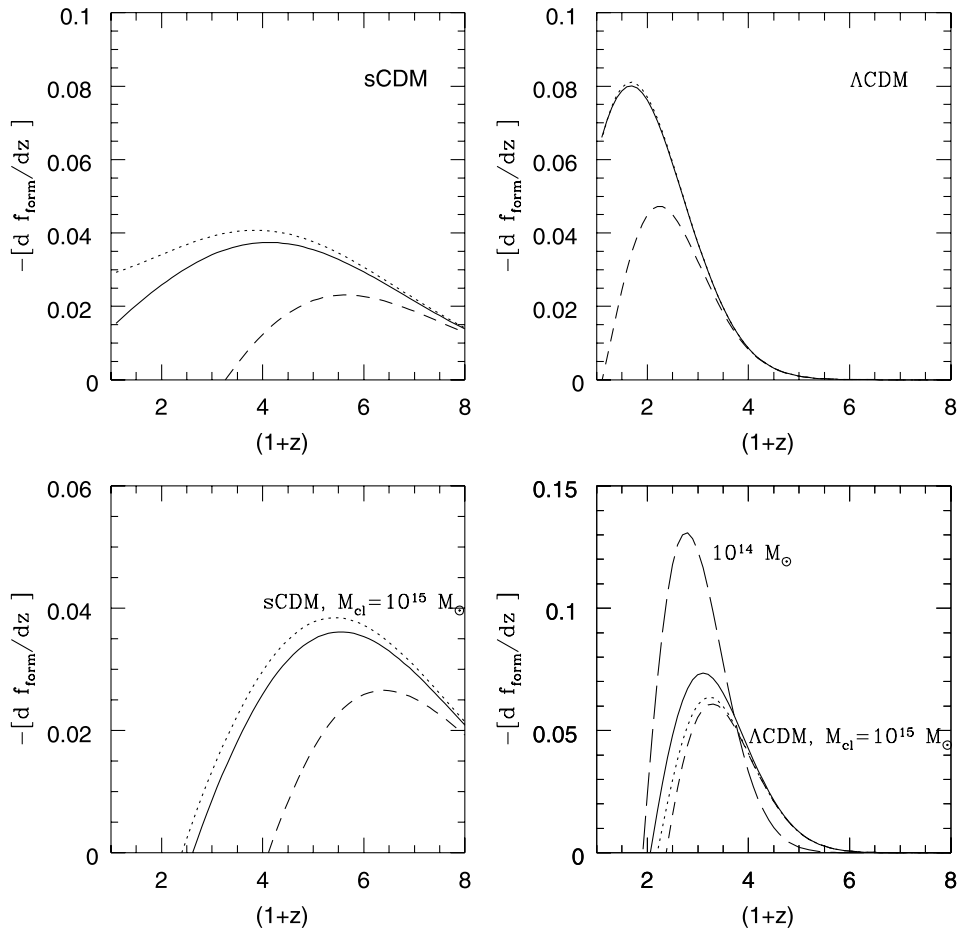


Figure 2. The net formation rate of objects (df_{form}/dz) plotted for objects of mass $10^{13} M_{\odot}$ (for $dM = M$), in clusters (bottom panels, with $M_{\text{cl}} = 10^{15} M_{\odot}$) and in general (top panels), and for sCDM models (left panels) and Λ CDM models (right panels). Dashed lines show the term df_{PS}/dz , dotted lines show df_{form}/dz using the Blain & Longair (1993) fit (with $\phi = 1.0$), and solid lines show df_{form}/dz using the integral in equation (8). The long-dashed line in the lower right panel shows the case for $M_{\text{cl}} = 10^{14} M_{\odot}$.

clusters) positive at low redshift. This is suggestive of the enhanced growth of perturbations, and the need for a *larger* rate of merger *inside* clusters, as previous authors have noted.

We also show the case for $M_{\text{cl}} = 10^{14} M_{\odot}$ by the long dashed line in the lower right panel of Fig. 2. Comparison with the solid line (for $M_{\text{cl}} = 10^{15} M_{\odot}$) shows that the formation rate of these galaxies inside a lower-mass cluster is larger. This follows simply from the extension of the PS formalism [from the dependence on the term $(\sigma_M^2 - \sigma_{M'}^2)$]. It is interesting to note that this is consistent with the suggestion from observation that quasars are preferentially located in groups of galaxies instead of rich clusters (e.g., Bahcall & Chokshi 1991; Fisher et al. 1996).

If our formalism is used without any correction, this will lead to *subtraction* of energy input in the final result. We circumvent this problem by putting $-(df_{\text{PS}}/dz) = 0$ when this term turns negative. This will, therefore, provide a lower limit to the total energy input from quasar outflows in a cluster.

We can finally write down the rate of formation of quasars in a given mass range $M, M + dM$ inside a cluster of a given mass, M_{cl} , in the form of the rate of increase of the fraction of mass of M_{cl} that is in quasars at an epoch z , as (remembering equation 3)

$$\begin{aligned} & \frac{df_{\text{q,cl}}(M, z | M_{\text{cl}}, z_f)}{dz} dM \\ &= f_{\text{q}} \frac{df(M, \delta_{\text{c}}(z) | M_{\text{cl}}, \delta_{\text{c}}(z_f))}{dz} dM \\ &+ f_{\text{q}} f(M, \delta_{\text{c}}(z) | M_{\text{cl}}, \delta_{\text{c}}(z_f)) dM \\ &\times \int_{2M}^{M_{\text{cl}}} \frac{d^2 p}{dM' dz} (M \rightarrow M' | z) dM', \end{aligned} \quad (9)$$

with the condition that $dn_{\text{q}}/dz = 0$, for $z < z_n$, where $(dn_{\text{q}}/dz)|_{z < z_n} > 0$. The integral on the right-hand side is evaluated using equation (7). Here we have also changed the upper limit of the integration to M_{cl} . The integrand is a rapidly decreasing function of M' and the value of the integral depends mostly on the lower limit.

Here, the threshold density contrast in a cosmological constant dominated universe is given by a fit given by Kitayama & Suto (1996),

$$\delta_{\text{c}}(z) = 1.68[g(z=0)/g(z)][1 + 0.0123 \log \Omega_{\text{m}}(z)]. \quad (10)$$

In our calculations, we have used a fit for $g(z)$ from Carroll, Press & Turner (1992),

$$\begin{aligned} & g(\Omega_{\text{m}}(z), \Omega_{\Lambda}(z)) \\ & \sim \frac{5\Omega_{\text{m}}(z)}{2\{\Omega_{\text{m}}(z)^{4/7} - \Omega_{\Lambda}(z) + [1 + \frac{1}{2}\Omega_{\text{m}}(z)][1 + \frac{1}{70}\Omega_{\Lambda}(z)]\}}, \end{aligned} \quad (11)$$

where (Lahav et al. 1991)

$$\begin{aligned} \Omega_{\text{m}}(z) &= \Omega_{\text{m}0}(1+z)^3 / [\Omega_{\text{m}0}(1+z)^3 + \Omega_{\Lambda 0}], \\ \Omega_{\Lambda}(z) &= \Omega_{\Lambda 0} / [\Omega_{\text{m}0}(1+z)^3 + \Omega_{\Lambda 0}]. \end{aligned} \quad (12)$$

We emphasize here that the above formalism leads to a *conservative* estimate of the abundance of quasars in a cluster (and, therefore, the final heat input), because (a) we ignore the increased pace of growth of perturbation and the merging rate inside a cluster, and (b) the lower limit of the integration could in reality be smaller than $2M$, which is probably the reason the rate of

formation still turns negative at low redshifts even after the addition of a merger term.

To summarize the work in this section, we have used existing ideas for relating the quasar formation rate to the Press–Schechter mass function, to estimate the rate of formation of quasars *inside clusters* (as a function of cluster mass and formation redshift), utilizing the extensions of the PS formalism. First, we use equation (1) to relate the PS mass function to quasar abundance. The standard PS mass function (equation 2) is then replaced by its extension (equation 4). Furthermore, we add the contribution due to merger (into larger objects) (equations 5 and 8), again using the extensions of the PS mass function. We finally have the rate of formation of quasars inside clusters as given by equation (9).

We use this formalism to calculate the total energy input from outflows from quasars in a cluster, or a lower limit to it. We next discuss the energy input from individual outflows, which we will combine with our calculation of formation rate of quasars for the final result.

3 WORK DONE BY QUASAR OUTFLOWS

In this section we calculate the energy input from quasar outflows into the ambient medium. We model the outflows as they evolve in the ambient medium and calculate the $p dV$ work done by the outflows. To begin with, we discuss the different kinds of outflows that we consider and the characteristics of the hosts of quasars.

3.1 Quasar outflows

We consider two major types of quasar outflows. For radio-loud quasars (RLQ), the outflow is in the form of a tightly collimated jet, which deposits energetic particles into a cocoon that expands against the surrounding medium. These outflows are characterized by the kinetic luminosity of the jet, L_{k} . According to Willot et al. (1999), this is correlated with the bolometric luminosity L_{bol} of the quasar, and $0.05 \lesssim L_{\text{k}}/L_{\text{bol}} \lesssim 1.0$. We follow FL01 in arguing that, since $L_{\text{bol}} \sim 10L_{\text{B}}$ (Elvis et al. 1994), the rest-frame B -band luminosity, $L_{\text{k}} \sim L_{\text{B}}$.

Radio-loud quasars, however, constitute only about 10 per cent of the total population of quasars (Stern et al. 2000). We therefore define a factor f_{o} for the fraction of quasars with outflows, and define $f_{\text{o}} \sim 0.1$ for our RLQ model. The fraction f_{o} here denotes the number of radio relics/lobes per halo, since a radio-loud quasar may have several outbursts of radio activity. This fraction is therefore a very conservative estimate since it is obtained from observed radio luminosity function and does not take into account the existence of radio relics in clusters.

Another important kind of outflow is encountered in broad absorption line (BAL) quasars. The absorption troughs are thought to be due to absorbing clouds flowing out of the quasars with velocities up to $0.1c$. Although they are encountered in about 10 per cent of quasars, it is believed that all quasars have such outflows (all the time) and the covering fraction of the BAL outflows is about 10 per cent (Weymann et al. 1991; Weymann 1997). Some authors also believe that BAL outflows have a limited lifetime (especially the low ionization BALs) and that they have a large covering fraction in the early phase of a quasar (Voit, Weymann & Korista 1993). For our calculation, we use a fraction $f_{\text{o}} \sim 1$ for the BAL outflows. We discuss the effect of the uncertainty in these factors on the final result in Section 5.

We model the BAL outflows as having a kinetic luminosity L_{k} . Following FL01, if N_{H} is the column density of the absorbing gas,

f_c the covering fraction and R_{BAL} is the size of the absorption system, then L_k is related to the outflow velocity v_{BAL} as

$$L_k \sim 2\pi f_c N_{\text{H}} m_p R_{\text{BAL}} v_{\text{BAL}}^3.$$

The observed ranges of these parameters are as follows: $v_{\text{BAL}} \lesssim 0.1c$, $f_c \sim 0.1$ (Weymann 1997; but see above), $R_{\text{BAL}} \sim 1-500$ pc and $N_{\text{H}} \sim 10^{22}-10^{23}$ cm $^{-2}$ (Krolik 1999; Gallagher et al. 1999). For these values, the magnitude of L_k is close to that of L_B . FL01 also argued that for BAL winds $L_k \sim 0.1-100L_B$, and finally assumed $L_k \sim L_B$. Since this estimate depends crucially on a number of uncertain parameters (for example, the fact that the absorption column density in optical measurement is much smaller than the above-mentioned X-ray column density), it may not really be a conservative estimate, but it does provide a simple scaling which we hope is not too unreasonable. The estimate $L_k \sim L_B \sim 0.1L_{\text{Edd}}$ is probably not a conservative estimate, but an upper limit, in that for a covering fraction of 10 per cent the mechanical luminosity of the wind could not be larger than 0.1 of the Eddington rate. Keeping all these uncertainties in mind, we assume that $L_k \sim L_B$ for BAL outflows.

We then need to connect L_B of a given quasar with the properties of its halo. First, as Haiman & Loeb (1998) have shown, the mass of the black hole at the centre is related to L_B as

$$M_{\text{BH}} = \frac{1}{0.093} \left(\frac{L_B}{1.4 \times 10^{38} \text{ erg s}^{-1}} \right) M_{\odot}, \quad (13)$$

assuming that the quasar radiates at the Eddington luminosity. The factor of 0.093 reflects the fraction of the Eddington luminosity radiated in the B band, taken from the median quasar spectrum of Elvis et al. (1994). Statistically speaking, we therefore assume that a fraction f_q (equation 1) of all black holes radiate at the Eddington rate for a lifetime of $\sim 10^7$ yr, while the rest does not radiate at all.

Secondly, the correlation between the central black hole mass and the total baryonic mass of the galaxy (Magorrian et al. 1998; Gebhardt et al. 2000) gives $M_{\text{BH}} \sim 4 \times 10^{-4} M_{\text{h}}$, where M_{h} is the total mass of the galaxy, using a value of $M_{\text{BH}}/M_{\text{baryonic}} \sim 2-3 \times 10^{-3}$ and $M_{\text{baryonic}}/M_{\text{h}} \sim \Omega_{\text{b}}/\Omega_0 \sim 0.2$.

As far as the collimation is concerned, the geometry of BAL outflows is still uncertain, whereas the radio jets are well collimated. Since some models do suggest a modest collimation even in BAL outflows, with a covering fraction of $f_c \sim 0.1$ (Weymann 1997), we use the idea of collimated outflows for outflows from both radio-loud and BAL quasars.

We next discuss the evolution of the outflow from radio-loud quasars and calculate the fraction of its total kinetic luminosity that it deposits into the surrounding medium in the form of p dV work. For concreteness, we will assume that the corresponding fraction for BAL outflows has similar values, and keep in mind the uncertainty in the geometry and energetics of BAL outflows.

3.2 Evolution of outflows

The standard scenario for outflows from radio-loud quasars involves a ‘cocoon’ surrounding the core and the jet, consisting of a shocked ambient medium and shocked jet material (Scheuer 1974; Blandford and Rees 1974). Begelman & Cioffi (1989) constructed a simple model of the evolution of a cocoon in which the cocoon is overpressured against the ICM. In their model, the expansion along the jet axis is determined by the balance of the thrust of the jet and the ram pressure, whereas the thermal pressure of the cocoon drives along the direction perpendicular to the jet axis. Results of

numerical simulations agree with this scenario (Loken et al. 1992; Cioffi & Blondin 1992).

Here we adopt the model of the evolution of cocoons following the approach of Bicknell, Dopita & O’Dea (1997), which is based on the Begelman & Cioffi (1989) model but includes the p dV work done by the cocoon, in order to find the fraction of total energy lost by the quasar to the ICM through mechanical work (p dV work). Bicknell et al. (1997) derived this fraction ($f_{p,dV}$) to be $f_{p,dV} = 3/8$, for a homogeneous ambient medium (their equation 2.13).

We will, however, calculate this fraction from numerical solution of the equations governing the evolution of the cocoon, for the following reasons. First, the derivation mentioned above implicitly assumes that the mean pressure averaged over the hotspot region is equal to the mean lobe pressure [in the language of Bicknell et al. (1997), this means $\zeta \sim 1$]. In fact, their equation (2.13) shows that for constant ζ one has in general

$$P_c \text{ d}V_c/\text{d}t = (1 + 2\zeta)L_j/8\zeta,$$

which recovers the fraction 3/8 for $\zeta = 1$. In reality, however, this ratio does not remain a constant in time. Secondly, this derivation is valid only during the period when the jet is active. Even after the jet switches off, the cocoon, however, continues to evolve as a result of its overpressure until it reaches an equilibrium pressure with the ambient medium (see also Nath 1995). The cocoon, therefore, continues to do p dV work even after the jet switches off, and the inclusion of this process will lead to an upward revision of the fraction of total energy that is lost in p dV work. Besides, there seems to be some confusion in the literature regarding the fraction. For example, Inoue & Sasaki (2001) have recently adopted a fraction $f_{p,dV} = 1/4$ in their calculation of energy input into the surrounding gas.

We therefore calculate this fraction by numerically solving the equations of cocoon evolution.

We consider two collimated steady jets advancing into the ambient ICM. The thermalized jet matter and the shock-compressed ICM matter form a cocoon around the jets and the cocoon expands with shocks advancing in directions both parallel and perpendicular to the jet axis. After this stage of evolution, when the jet turns off after a lifetime of $t_{\text{life}} \sim 3 \times 10^7$ yr (Kaiser 2000), cocoons still retain high pressure. They cool radiatively and expand due to overpressure till it reaches a pressure equilibrium with the ambient medium. Thus the relevant equations are:

$$dr_{\text{h}}/\text{d}t = (L_j/A_{\text{h}}\rho_{\text{a}}\beta c)^{1/2}, \quad t < t_{\text{life}} \quad (14)$$

$$= (P_c/\rho_{\text{a}})^{1/2}, \quad t > t_{\text{life}}$$

$$dr_c/\text{d}t = (P_c/\rho_{\text{a}})^{1/2}, \quad (15)$$

$$dE_c/\text{d}t = L_j - P_c \text{ d}V_c/\text{d}t \quad (16)$$

where L_j is the jet luminosity, ρ_{a} is the density of the ambient medium and βc is the velocity of the jet material. As the jet is highly relativistic, $\beta \sim 1$. The averaged hotspot area $A_{\text{h}} \sim 30$ kpc 2 (Bicknell et al. 1997) is assumed to be larger than the radius of the jet, according to the ‘dentist’s drill’ model of the jet (Scheuer 1982). Here r_{h} is the length of the jet or the distance of the hotspot from the centre of the galaxy, r_c is the half-width of the cocoon at the centre and V_c is the volume of the cocoon given by $V_c = \epsilon_v(2\pi r_c^2)r_{\text{h}}$ where ϵ_v is the geometrical factor depending on the shape of the cocoon. For our calculations we have taken the shape of the cocoon as biconical and so $\epsilon_v \sim 1/3$. Finally P_c is the pressure inside the cocoon given by $P_c = (E_c/V_c)(\gamma - 1)$ where

$\gamma = 4/3$ and E_c is the total energy inside the cocoon given by $E_c = L_j t$ till $t < t_{\text{life}}$ and $E_c = L_j t_{\text{life}}$ afterwards, where t_{life} is the lifetime of the jet.

This is admittedly a simplified model of the evolution of the cocoon. In reality, after the jet switches off, one expects Rayleigh–Taylor and Kelvin–Helmholtz instabilities to distort the cocoon, giving rise to ‘buoyant’ plumes. This phase of the evolution of the cocoon, and its effect on the ambient medium, has been addressed by various authors (Gull & Northover 1973; Churazov et al. 2001; Brüggén & Kaiser 2001), mainly with the help of numerical simulations. It is possible that this phase adds substantial heating to the intracluster medium, but short of doing a numerical simulation it is difficult to assess its importance. We have also neglected the loss of energy through radiation, as modelling such losses would involve detailed knowledge of various parameters (e.g., electron energy spectrum and the possibilities of re-acceleration of electrons). With the uncertainties involved in modelling these processes, it seems reasonable to adopt the above simplified picture as a pointer and keep the uncertainties in mind while discussing the final result. In light of this discussion, we will also calculate the final result with a value of $f_{p,dV} = 3/8$ as in Bicknell et al. (1997), which we will adopt as a conservative lower limit.

We numerically calculate the volume of the cocoon as it grows into the ICM and the pressure inside the cocoon at each step, and add up the $p dV$ work to get the final amount of energy lost in this mode. The evolution of the cocoon is continued until the pressure inside the cocoon becomes equal to the ambient pressure, $n_a k_B T_a$ ($n_a = \rho_a / \mu m_p$), where T_a is the temperature of the ICM. The fraction $f_{p,dV}$ is calculated by taking the ratio of the total energy lost through mechanical work to the total energy (that is $L_j t_{\text{life}}$). (We found that, for the relevant values of the ambient medium parameters, the time-scale to reach pressure equilibrium is always larger than t_{life} .) The dependence of $f_{p,dV}$ on the ambient density is shown in Fig. 3 for $T_a = 10^6, 10^7$ and 10^8 K, and for $L_k = 10^{46}$ (solid lines) and 10^{47} erg s $^{-1}$ (dashed lines).

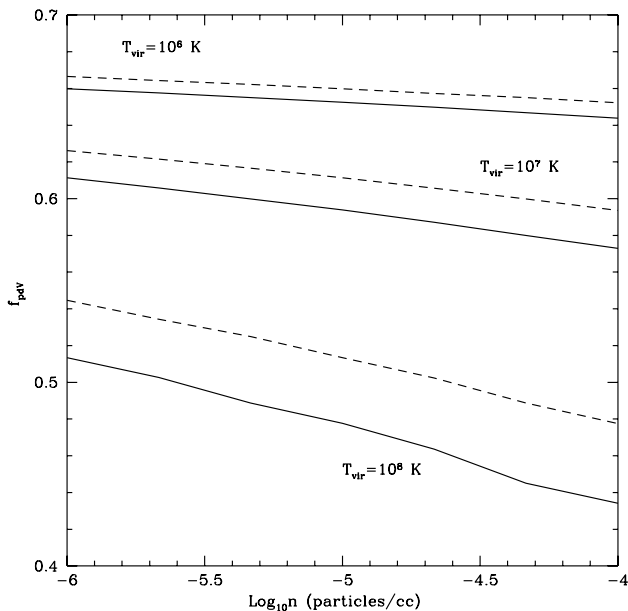


Figure 3. The dependence of $f_{p,dV}$ with the ambient density n for various ambient temperatures T for various L_k (erg s $^{-1}$) with and without radiation loss. Solid and dashed lines are for $L_k = 10^{46}$ erg s $^{-1}$ and 10^{47} erg s $^{-1}$ respectively without radiation loss. Dotted lines are for $L_k = 10^{46}$ with radiation loss.

The plot shows that the fraction $f_{p,dV}$ is a function of the temperature of the cluster and also the density of the ambient medium. The general trend is that the fraction reduces at higher temperatures and higher densities. This is because of the fact that the cocoon reaches pressure equilibrium with the ambient medium sooner for a higher pressure environment (higher T_a and n_a), and the total $p dV$ work ends up being smaller. The plot also shows that the fraction $f_{p,dV}$ depends weakly on the jet luminosity for lower temperatures ($T_a \leq 10^7$ K), whereas there is a bit of a difference for $T_a \sim 10^8$ K.

The fraction $f_{p,dV}$ calculated above is somewhat larger than that used in the literature for ambient medium with low pressure. The difference is mainly the result of our inclusion of cocoon evolution even after the jet has switched off. Incidentally, Inoue & Sasaki (2001), while using a value of $f_{p,dV} = 1/4$, discussed the possibility that this fraction could be larger in reality, *because* of its continued evolution after the switching off of the jet (their section 3.2).

In our calculation for the total $p dV$ work done by BAL and RLQ outflows, we will use the values of $f_{p,dV}$ obtained above. As mentioned earlier, the energetics and geometry of BAL outflows are not clear at present. For concreteness, we have worked out the case of RLQ outflows in detail, and we will use the same values of $f_{p,dV}$ for BAL outflows as well.

4 HEATING OF THE ICM

Equipped with the knowledge of the rate of formation of quasars in clusters (equation 9) and the fraction of total energy that is deposited as $p dV$ work by the outflows from them (Section 3.2), we are now in a position to calculate the total amount of non-gravitational energy provided by quasar outflows in a cluster. If we denote the gas fraction of the total cluster mass by f_{gas} , then the total number of gas particles is $\sim M_{\text{cl}} f_{\text{gas}} / m_p$, where m_p is the proton mass. It is not yet clear if the gas fraction has a universal value for clusters of all masses and at all redshifts. A few authors

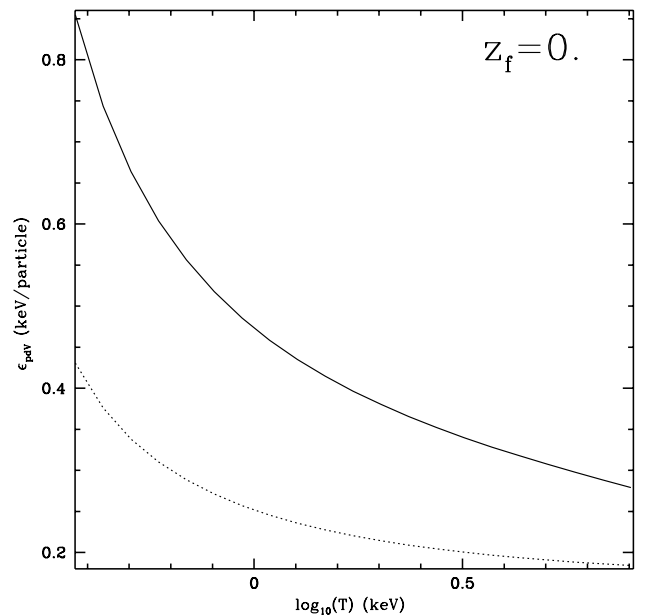


Figure 4. Excess energy (keV) from BAL outflows as a function of the cluster virial temperature (keV) for clusters with $z_f = 0$. The solid line shows the result of our calculation using the density- and temperature-dependent $f_{p,dV}$ and the dotted line shows the results when $f_{p,dV} = 3/8$.

(see, e.g., Schindler 1999) have found no correlation of the gas fraction with the cluster mass, whereas others (e.g., Ettori & Fabian 1999) find some correlation (with low-mass clusters possessing lower gas fraction). This correlation has been attributed to excess energy deposition (since low-mass clusters are more liable to lose gas from heating than massive clusters) (see, e.g., Bialek, Evrard & Mohr 2001). Here, however, we are trying to calculate the magnitude of this very excess energy. It would not be appropriate to include this correlation a priori in our calculation. (We show later that including these correlations only increases our estimate of excess energy input.) We have, therefore, used a value of $f_{\text{gas}} = 0.1$ for our calculation. The total energy per (gas) particle deposited into the ICM of a cluster of mass M_{cl} is then given by

$$\epsilon_{\text{p dV}} = \frac{m_{\text{p}}}{M_{\text{cl}} f_{\text{gas}}} \int_{z_{\text{m}}}^0 \int_{M_1}^{M_{\text{u}}} \frac{df_{\text{q,cl}}(M, z | M_{\text{cl}}, z_{\text{f}})}{dM dz} \times \frac{M_{\text{cl}}}{M} dM dz f_{\text{o}} [L_{\text{k}} t_{\text{life}} f_{\text{p dV}}(n_{\text{a}}, T_{\text{a}})], \quad (17)$$

where

$$\frac{d^2 n_{\text{q}}(M, z | M_{\text{cl}}, z_{\text{f}})}{dM dz}$$

is calculated using equation (9). The factor $f_{\text{o}} \sim 1$ for the BAL outflows, and $f_{\text{o}} \sim 0.1$ for outflows from RLQs (Section 3). The redshift z_{m} is the maximum redshift of heat input. We later show (Fig. 7) that the heat input is negligible for $z \geq 5$. The density and temperature of the ICM of a cluster of a given mass (M_{cl}) and formation redshift (z_{f}) is calculated using (Eke, Navarro & Frenk 1998)

$$T_{\text{a}} = 1.65 \times 10^7 (1+z) \left[\frac{M_{\text{cl}}}{10^{15} h^{-1} M_{\odot}} \right]^{2/3} \left[\frac{\Omega_0 \Delta(\Omega_0, z)}{\Omega(z)} \right]^{1/3} \quad (18)$$

and

$$n_{\text{a}} = \frac{M_{\text{cl}} f_{\text{gas}}}{m_{\text{p}} (4/3) \pi r_{\text{vir}}^3}, \quad (19)$$

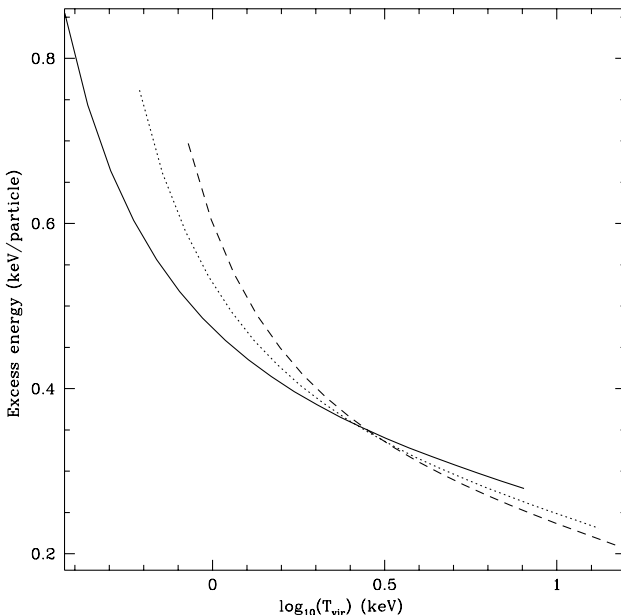


Figure 5. The excess energy for $z_{\text{f}} = 0$ (solid line), 0.5 (dotted line) and 1 (dashed line).

where

$$\Delta(z) = 18\pi^2 + 82x - 39x^2 \quad (20)$$

and $x = \Omega(z) - 1$ (Bryan & Norman 1998), where we use equation (12) to compute $\Omega(z)$. For r_{vir} , we use

$$r_{\text{vir}} = \left[\frac{3M_{\text{vir}}}{4\pi\Delta(z)\rho_{\text{crit}}(z)} \right]^{1/3}, \quad (21)$$

where we have used $M_{\text{vir}} = M_{\text{cl}}$, and ρ_{crit} is the critical density of the universe. The densities used in the following calculations range between 10^{-4} and 10^{-6} cm^{-3} .

The integral in equation (17) is evaluated using $M_1 = 10^{12} M_{\odot}$ and $M_{\text{u}} = 10^{13} M_{\odot}$, for $z_{\text{f}} = 0, 0.5$ and 1 for different values of M_{cl} . We present the results for the total non-gravitational energy input per particle as a function of cluster mass (or, equivalently, gas temperature) in Fig. 4 (for $z_{\text{f}} = 0$). The solid curve shows the heat input calculated using $f_{\text{p dV}}$ from Section 3.2. The dotted line shows the case for a constant $f_{\text{p dV}} = 3/8$ (Bicknell et al. 1997). We show the results for different z_{f} in Fig. 5 (against T) and Fig. 6 (against cluster mass).

We also show in Fig. 7 the rate of deposition of energy ($-d\epsilon_{\text{p dV}}/dz$) as a function of the redshift for three clusters of masses $M = 2 \times 10^{13}, 10^{14}$ and $10^{15} M_{\odot}$, all for $z_{\text{f}} = 0$. All the curves drop to zero at low redshift because of the condition $dn_{\text{q,cl}}/dt = 0$ (in equation 9). In reality the contribution to the heating should be small but non-zero, and will increase the estimate of excess energy.

5 DISCUSSION

It has been estimated that the amount of excess energy required to explain the observation is of the order of 0.5–3 keV per particle (Navarro et al. 1995; Cavaliere et al. 1997; Balogh et al. 1999; Wu et al. 2000). Recently, however, Lloyd-Davies et al. (2000) have shown from observations of groups of clusters that an excess energy of 0.44 ± 0.3 keV per particle suffices to explain the excess

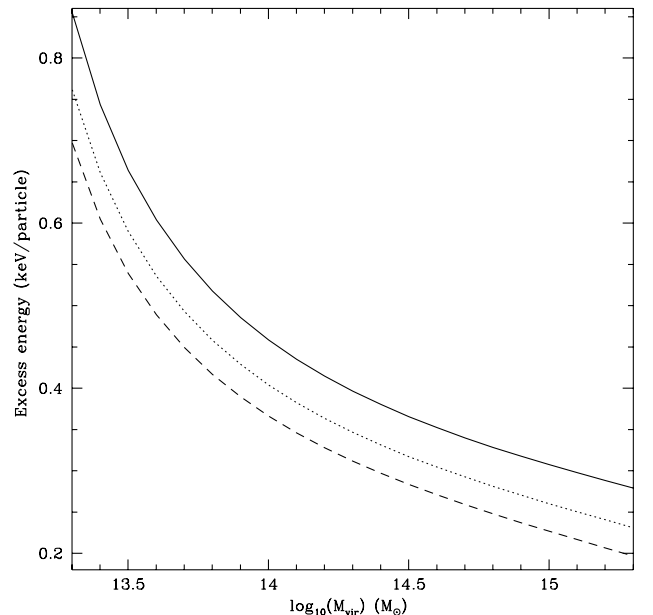


Figure 6. The excess energy plotted against the cluster/group mass for $z_{\text{f}} = 0$ (solid line), 0.5 (dotted line) and 1 (dashed line).

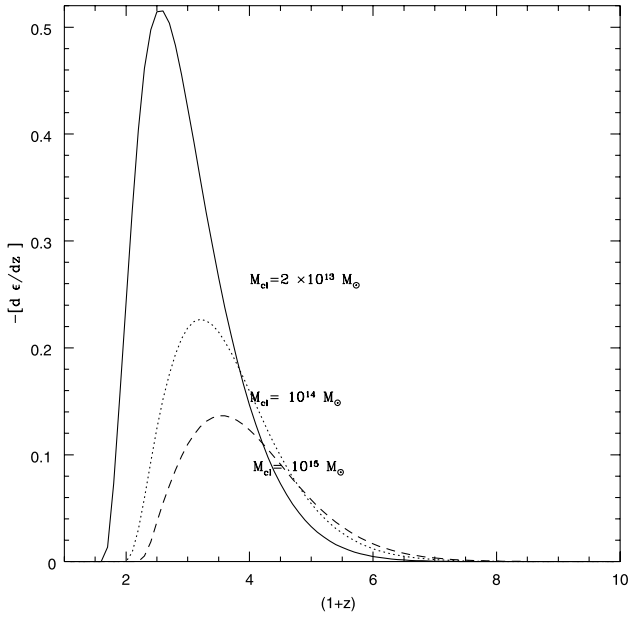


Figure 7. The rate of deposition of excess energy ($-d\epsilon_{p,dV}/dz$) as a function of redshift, for $M_{cl} = 2 \times 10^{13} M_{\odot}$ (solid line), $10^{14} M_{\odot}$ (dotted line) and $10^{15} M_{\odot}$ (dashed line).

entropy in groups. They showed that this can explain the entropy floor for galaxy groups with temperature $T \lesssim 4$ keV. Borgani et al. (2001) have also shown with the aid of numerical simulations that excess energy of order ~ 1 keV per particle reproduces the observations.

The solid and dashed curves in Figs 4 and 5 show that the excess energy from p dV work done by quasar outflows falls in this required range. It is seen that the excess energy per particle is larger for clusters or groups with lower temperature. This is due to two factors: (a) the number of quasars per unit mass is larger for smaller clusters, and (b) the fraction of total energy in outflows that is lost in p dV work is larger for them. We show the results in the case of a constant $f_{p,dV} = 3/8$ (as in Bicknell et al. 1997) with dotted lines. It is interesting to note that even in this case the excess energy is in the required range (0.44 ± 0.3 keV per particle), especially for groups with low temperatures, as advocated by Lloyd-Davies et al. (2000). Incidentally, this is larger than the estimate of excess energy from galactic winds ($\lesssim 0.1$ keV per particle; Wu et al. 2000).

We found that our results for the excess energy (solid line in Fig. 4) can be approximated by a fit of type (in keV per particle)

$$\epsilon_{p,dV} \sim 0.258 \left(\frac{T}{10 \text{ keV}}\right)^{-0.193} + 0.033 \left(\frac{T}{2 \text{ keV}}\right)^{-1.2} \quad (22)$$

for $0.50 \text{ keV} \leq T \leq 8.0 \text{ keV}$.

We also found that the results for the excess energy taking $f_{p,dV} = 3/8$ (Bicknell et al. 1997) (dotted line in Fig. 4) can be approximated by a fit of type (in keV per particle)

$$\epsilon_{p,dV} \sim 0.17 + 0.045 \left(\frac{T}{2 \text{ keV}}\right)^{-0.95} \quad (23)$$

for $0.50 \text{ keV} \leq T \leq 8.0 \text{ keV}$.

We compare our results with the data from Lloyd-Davies et al. (2000) in Fig. 8. We show the predictions of our calculation as the thin solid curve (corresponding to the solid curve in Fig. 4), where

the data points have been taken from Lloyd-Davies et al. (2000) (their fig. 9). We also show the result of the calculations for $f_{p,dV} = 3/8$ by the thick solid line. The data points refer to the binding energy of a (constant) central fraction (0.004) of the virial mass of groups and clusters. The dashed line shows the case for self-similar models, of type $E \propto T$, derived from the data points for rich clusters. The dotted line shows their fit (with a constant excess energy of 0.44 keV per particle) along with a formal 1σ confidence interval shown by the shaded region. The figure shows that our predictions are consistent with the data available at present. The thick line (corresponding to $f_{p,dV} = 3/8$) falls close to the fit provided by Lloyd-Davies et al. (2000), whereas the thin line (using $f_{p,dV}$ from Fig. 3) somewhat overestimates the heat input at the low-mass end. The thick line can be viewed as a conservative estimate of the heat input, since it uses $f_{p,dV} = 3/8$. The thin line, however, provides an estimate of the heat input if $f_{p,dV}$ is much larger than $3/8$. We should remind ourselves here that we have calculated $f_{p,dV}$ for radio galaxies and used the same values for BAL outflows. If a more accurate estimate of $f_{p,dV}$ for BAL outflows is worked out in the future, the resulting heat input into the ICM could then be scaled accordingly using Fig. 8.

We would like to emphasize here again that our calculation provides a conservative estimate of the excess energy, for reasons outlined in Section 2. Moreover, we have used a constant density and temperature in time for the ICM gas (for clusters with a given z_f), which is not very realistic. In reality, the density at higher redshift will be smaller, and the inclusion of a density-dependent $f_{p,dV}$ will only increase the estimate of excess energy (since this fraction increases with decreasing density).

The curves for clusters forming at different epochs show excess energies to decrease somewhat for clusters with higher formation redshift.

The curves of Figs 4 and 5 assume $f_o \sim 1$, which is relevant for BAL outflows. The excess energy from RLQ outflows will be one-tenth of these curves, showing the difficulty of using radio galaxies as the only source of non-gravitational heating, if conservative estimates for their kinetic luminosities are used. Recently, Inoue & Sasaki (2001) have used the radio luminosity functions of Ledlow & Owen (1996) and Willot et al. (2001) to determine the abundance of radio galaxies in clusters, and finally to estimate the total p dV work done by the cocoons of these radio galaxies. They estimated an excess energy of order 1 keV per particle for rich clusters like the Coma cluster, and also for poor groups, assuming that their ratio of radio galaxies per unit cluster mass is universal. From our calculation, we find an excess energy from only radio-loud quasars that is an order of magnitude *smaller* than their estimate. It is possible that the assumptions leading to the estimate of L_k are at the source of this difference (see, e.g., the discussion on the uncertainty in the factor f_j in their section 3.2).

The evolution of the rate of energy input (Fig. 7) shows that the heat input before $z \sim 5$ is almost negligible. Most of the heating occurs in the range $z \sim 1-4$. It is also clearly seen that the ICM of poor groups is heated at lower redshifts, compared to the gas in massive clusters. This follows from the simple consideration that the evolution of objects in a given mass range (here, that of the quasars) typically occurs earlier in more massive clusters. Our results, therefore, suggest that the ICM in clusters were ‘pre-heated’, before the major mergers took place in them, whereas the ICM in groups of galaxies were heated at epochs similar to that of their formation. It is interesting to note that the energy input epoch lies below the upper limit from recent observations of Lloyd-Davies et al. (2000). It is also consistent with the limit on the range

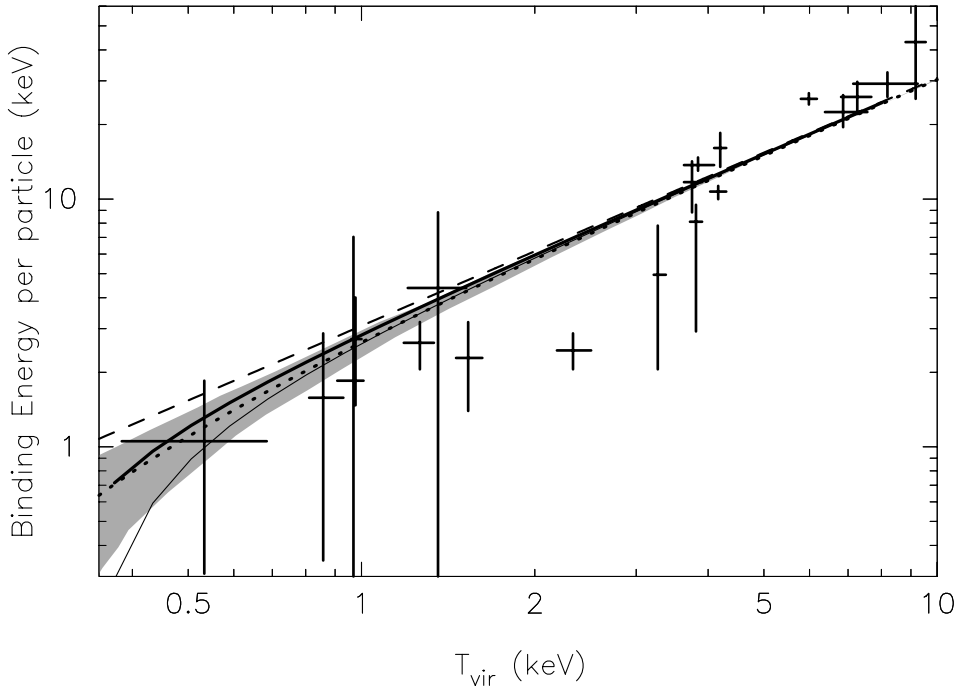


Figure 8. The prediction from our calculations presented in the form of the final binding energy per particle of the central region of groups and clusters plotted against the gas temperature. The data points are from fig. 9 of Lloyd-Davies et al. (2000), and the dashed line refers to their fit $E \propto T$, derived from the data points for clusters with $T \geq 4$ keV. The dotted line refers to their second fit, with a constant excess energy of 0.44 keV per particle (subtracted from the binding energy) along with a formal 1σ confidence interval shown by the shaded region. The thick solid line uses $f_{p,dV} = 3/8$ and the thin solid line uses $f_{p,dV}$ from Fig. 3.

of redshifts as shown by Yamada & Fujita (2001) given the level of uncertainty.

Recently, accelerated particles from shocks as a result of the formation of clusters have been hypothesized to be the source for the diffuse γ -ray background (Loeb & Waxman 2000). This γ -ray production will be suppressed, however, if the gas in the clusters were heated substantially at earlier epochs (Totani & Inoue 2002). Our results show that the ICM in massive clusters were pre-heated (and earlier than the ICM in groups), and the suppression of the γ -ray production will, therefore, be an important effect, if confirmed.

Finally, we discuss the uncertainties involved in our calculation. Apart from the uncertainties in cosmological parameters, the major uncertainties lie in the factors f_q (Section 2) (connecting the abundances of quasars with the PS mass function), f_o (the fraction of quasars with outflows), $f_{p,dV}$ and f_{gas} . Among these, the most uncertain factor is f_o , which we have assumed to be of the order of unity for BAL outflows. The uncertainty in this factor will be reflected in the uncertainty of the final heat input (with a direct proportionality; see equation 17). The uncertainty in $f_{p,dV}$ has already been discussed, and we found that, even if $f_{p,dV}$ is as low as $3/8$ for all cases, the final excess energy is certainly larger than that from supernovae-driven winds, and is still within the required range of excess energy, especially for loose groups. As far as the uncertainty in f_{gas} is concerned, we have also done our calculation with a varying f_{gas} , e.g., of the type,

$$f_{gas} = 0.15 M_{\odot} (1+z)^{-0.5} (M_{cl}/10^{15} h^{-1})^{0.1}, \quad (24)$$

as has been advocated by Ettori & Fabian (1999), and we have found that the excess energy is approximately doubled in this case. However, it is not clear if this correlation is a result of the excess energy, and so it would not be appropriate to attach much

significance to this result. We have also varied the lower limit in our estimate of f_q (equation 1) and found that changing the lower limit from $10^{12} M_{\odot}$ to $10^{11} M_{\odot}$ increases the final heat input by only ~ 10 per cent. This is because of the fact that the increase in the number of quasars is compensated by the decrease in their mechanical luminosity. Lastly, we have already discussed in detail the uncertainty in the net formation rate of quasars in clusters, and, as explained in Section 2, our approach here has been very conservative, and the final results should be regarded as conservative estimates in this regard.

6 SUMMARY

We have calculated the excess energy deposited by quasar outflows in clusters in order to explain the observations of excess entropy in groups and clusters of galaxies. We summarize our findings below:

(1) We have used the extended Press–Sechter formalism to derive a formation rate of quasars inside clusters and groups, as a function of the cluster/group mass and its formation redshift.

(2) We have calculated the fraction of the kinetic luminosity of outflows (RLQ and BAL outflows) that is deposited on to the ambient medium, as a function of the density and temperature of the ambient medium. For outflows from radio-loud quasars, we have included the evolution of the cocoon after the jet turns off.

(3) The final excess energy from the mechanical work done by quasar outflows is found to be of the order of 0.18–0.85 keV per particle, and is consistent with the data available at present. The excess energy in this scenario comes mainly from BAL outflows, with radio galaxies supplying about a tenth of the total. Keeping in mind the uncertainties in the estimate of energetics and abundances of radio and BAL outflows, we conclude that both radio galaxies and BAL outflows are promising candidates for heating the ICM.

We found that this excess energy increases with decreasing mass of the cluster/group. This prediction could be tested with better data in the near future. The excess energy does not depend strongly on the formation redshift.

(4) The epoch of heating is found to be in the range $z \sim 1-4$, where this epoch is at lower redshifts for low-mass clusters.

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