Fast growth of magnetic fields in galaxy clusters: a self-accelerating dynamo

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November 11, 2018; e-print \texttt{astro-ph/0508535}

Abstract. We propose a model of magnetic-field growth in galaxy clusters whereby the field is amplified by a factor of about $10^8$ over a cosmologically short time of $\sim 10^8$ yr. Our model is based on the idea that the viscosity of the intracluster medium during the field-amplification epoch is determined not by particle collisions but by plasma microinstabilities: these give rise to small-scale fluctuations, which scatter particles, increasing their effective collision rate and, therefore, the effective Reynolds number. This gives rise to a bootstrap effect as the growth of the field triggers the instabilities which increase the Reynolds number which, in turn, accelerates the growth of the field. The growth is explosive and the result is that the observed field strength is reached over a fraction of the cluster lifetime independent of the exact strength of the seed field (which only needs to be above $\sim 10^{-15}$ G to trigger the explosive growth).

Key words: galaxies: clusters: general — intergalactic medium — galaxies: magnetic fields

1. Introduction

Rapidly improving measurements of Faraday Rotation in galaxy clusters reveal that intracluster medium is permeated by tangled magnetic fields with strength of a few $\mu$G and typical coherence scales of $1 - 10$ kpc (Kronberg 1994; Clarke, Kronberg & Böhringer 2001; Carilli & Taylor 2002; Govoni & Feretti 2004; Voge\textsuperscript{1} & Enß\textsuperscript{1}lin 2003, 2005). This means that the intracluster medium (ICM) contains huge amounts of magnetic energy and it is natural to ask whether there is a robust and universal mechanism that could be responsible for generating such fields and that would also be insensitive to the particular circumstances in individual clusters (specific age of the cluster, magnitude of the seed field, etc.). We would like to forego the discussion of whether cluster fields could be purely of primordial or external origin (as suggested by several authors; see, e.g., Kronberg et al. 2001; Banerjee & Jedamzik 2003) and instead show that they can be amplified very quickly and with no stringent constraints on the magnitude of the seed field by the small-scale dynamo effect due to the turbulent motions of the ICM. The idea that the fields are dynamo generated has been explored by many authors Jaff\textsuperscript{1} 1980; Roland 1981; Ruzmaikin, Sokoloff & Shukurov 1989; De Young 1992; Goldschmidt & Rephaeli 1993; Kulsrud et al. 1997; Sánchez-Salcedo, Brandenburg & Shukurov 1998; Subramanian, Shukurov & Haugen 2005; Enß\textsuperscript{1}lin & Vogt 2005).

The main issues are whether there is, indeed, turbulence in clusters and whether the rms rate of strain associated with this turbulence is sufficiently large to exponentially grow the field to the observed strength in the cluster lifetime of a few Gyr.

There are many possible energy sources in clusters that almost certainly produce random fluidlike motions of the ICM: merger events, subcluster and galactic wakes, active galactic nuclei. Various indirect observational estimates (e.g., Schuecker et al. 2004; Rebusco et al. 2005; Fujita 2005) appear to converge in expecting random flows with velocity fluctuations $U \sim 10^{-2} - 10^{-3}$ km/s at the outer scale $L \sim 10^2 - 10^3$ kpc (a direct detection may be achieved in the near future; see Inogamov & Sunyaev 2003). Similar numbers are obtained in numerical simulations of cluster formation (Norman & Bryan 1999; Sunyaev, Norman & Bryan 2003; Ricker & Sarazin 2001; Takizawa 2005). There is, however, no consensus on whether turbulence, at least in the usual hydrodynamic sense of a Kolmogorov energy cascade across a broad inertial range of scales, is a generic feature
of clusters (Fabian et al. 2002). The main difficulty is the very large values of the ICM viscosity obtained via the standard estimate $\sim \nu_{thi}\lambda_{mfp}$, where $\nu_{thi} \sim 10^3 \text{ km/s}$ is the ion thermal speed and $\lambda_{mfp} \sim 1-10 \text{ kpc}$ is the mean free path (assuming fully ionised hydrogen ICM with temperature $T \sim 10^8 \text{ K}$ and density $\sim 10^{-3} \text{ cm}^{-3}$). This gives $Re \sim 10^2$ if not less, which makes the existence of a well-developed inertial range doubtful.

In fact, the small-scale dynamo does not require a turbulent velocity field in the sense of a broad inertial range. As the dynamo effect is simply the random stretching of field lines, it is always controlled by the rms rate of strain — and this is associated with the viscous-scale motions, which are random but spatially smooth. Whether the viscous scale $\sim 10^{-3} \text{ km/s} / \nu_{thi}$ is much smaller than the outer scale $L$ is inessential as long as random motions are maintained — not a problem in clusters, given the variety of available stimuli mechanisms. Thus, the field will be amplified and the only issue is whether this happens sufficiently fast, given the relatively small growth rate obtainable in a viscous ICM. We shall see that this turns out not to be a problem because the effective ICM viscosity in the weak-field regime is much smaller than $\nu_{thi}\lambda_{mfp}$.

The widely used viscosity estimate $\sim \nu_{thi}\lambda_{mfp}$ only applies when the magnetic field is so weak that that the ion gyroradius is larger than the mean free path: $\rho_i > \lambda_{mfp}$. This gives $B \lesssim 10^{-18} \text{ G}$ — much smaller than the observed field in clusters. For stronger fields, the plasma becomes magnetised and its viscosity anisotropic: the standard value still works parallel to the magnetic field, while the motions perpendicular to the field are virtually undamped (Braginskii, 1965). As was pointed out recently by Schekochihin et al. (2005), the anisotropy of the viscous stress in clusters means anisotropy of the plasma pressure, which, in turn, always gives rise to well-known plasma microinstabilities (most importantly, mirror and firehose, see, e.g., Gary, 1993) at scales all the way down to the ion gyroscale. The emergence of the pressure anisotropies is intimately related to the changes in the strength of the magnetic field: in a magnetised plasma, the conservation of the first adiabatic invariant for each particle, $\mu = v_i^2 / B$, implies that the perpendicular plasma pressure changes according to

$$\frac{1}{p_\perp} \frac{dp_\perp}{dt} = \frac{\nu_{thi}}{\nu_{ii}} \frac{p_\perp - p_\parallel}{p_\perp},$$

where the first term is due to the conservation of $\mu$ and the second term represents collisional isotropisation ($\nu_{ii} \sim \nu_{thi} / \lambda_{mfp}$ is the ion collision frequency). Balancing the two terms, we get

$$\frac{1}{B} \frac{dB}{dt} = \nu_{ii} \Delta,$$

where $\Delta = (p_\perp - p_\parallel) / p_\perp$. On the other hand, since the resistivity of the ICM is tiny ($\eta_{ii} \sim 10^{20}$), the field lines are frozen into the fluid flow and the change of the magnetic-field strength can be expressed in terms of the fluid velocity $u$:

$$\frac{1}{B} \frac{dB}{dt} = \hat{b} : \nabla u - \nabla \cdot u = \frac{U}{L} Re^{1/2},$$

where $\hat{b} = B / B$. The last equality in Eq. 5 is based on estimating the turbulent rate of strain at the viscous scale. It is the parallel viscosity that matters here because the motions that change $B$ are precisely the ones — and the only ones — that are damped by it (Braginskii, 1965). Since

$$Re = \frac{UL}{\nu_{thi}} = \frac{UL}{\nu_{ii}},$$

we have, from Eqs. 4 and 5,

$$\Delta = \frac{U}{\nu_{thi}} \left( \frac{U}{L} \right)^{1/2} = \left( \frac{U}{\nu_{thi}} \right)^2 Re^{-1/2}.$$

The peak growth rate of these instabilities (at the ion gyroscale) is

$$\gamma_{\max} = \left( \frac{\Delta}{\beta} \right)^{\alpha_1} \Omega_i,$$

where $\beta = 8\pi n T / B^2$, $\Omega_i = eB / cm_i$ is the ion cyclotron frequency, $\alpha_1 = 1$ for the mirror and $\alpha_1 = 1/2$ for the firehose instability. To sum up briefly, external energy sources drive random motions, which change the field [Eq. 5], which gives rise to pressure anisotropies [Eq. 2], which trigger the instabilities. The latter are stabilised if the magnetic field is sufficiently strong: $\beta < 2 / \Delta$.

It is not currently clear exactly how the instabilities affect the structure of the turbulence and magnetic fields in clusters. One plausible hypothesis is that the saturation is quasilinear with small fluctuations $\delta B$ at the gyroscale:

$$\frac{\delta B^2}{B^2} = \left( \frac{\Delta}{\beta} \right)^{\alpha_2}.$$

where $\alpha_2$ is some positive power. These fluctuations will scatter particles, giving an effective collision frequency:

$$\nu_{scatter} = \frac{\delta B^2}{B^2} \gamma_{\max} = \left( \frac{\Delta}{\beta} \right)^{\alpha} \Omega_i,$$

where $\alpha = \alpha_1 + \alpha_2 > 0$. The result is an effective Reynolds number [Eq. 4 with $\nu_{ii} \rightarrow \nu_{scatter}$] that is much larger than the one based on the actual ion-ion collisions and that increases with the magnetic-field strength (because $\nu_{scatter} \propto \Omega_i \sim B$). Since this, in turn, speeds up the small-scale dynamo by increasing the turbulent rate of strain [Eq. 2], we expect explosive growth of magnetic field. This possibility was briefly referred to in Schekochihin et al. (2005). In §4 we shall construct a simple model of such a self-accelerating dynamo and evaluate its effectiveness in generating observed cluster fields for a set of fiducial cluster parameters (introduced in §2). We will conclude with a summary and some words of caution in §4.
2. Cluster parameters

In order to make specific estimates as we proceed, let us assume a fiducial set of parameters for the ICM:

\[ n = 10^{-3} \text{ cm}^{-3}, \]  
\[ T = 10^8 \text{ K} \Rightarrow \nu_{\text{thi}} = \sqrt{\frac{T}{m_i}} \approx 1000 \text{ km/s}, \]  
\[ U = 300 \text{ km/s} \text{ (rms turbulent velocity)}, \]  
\[ L = 200 \text{ kpc} \text{ (outer scale)}, \]  
\[ \lambda_{\text{mfp}} = 1 \text{ kpc}. \]

We shall measure magnetic field in units of

\[ B_{\text{eq}} = (8\pi nT)^{1/2} \approx 20 \mu G, \]  

for which \( \beta(B_{\text{eq}}) = 1 \). For this field, the ion cyclotron frequency and the ion gyroradius are

\[ \Omega_{i,\text{eq}} = \frac{eB_{\text{eq}}}{cm_i} \approx 0.2 \text{ s}^{-1}, \]  
\[ \rho_{i,\text{eq}} = \frac{v_{\text{thi}}}{\Omega_{i,\text{eq}}} \approx 5000 \text{ km}. \]

The magnetic field strength above which plasma is magnetised (\( \rho_i < \lambda_{\text{mfp}} \)) is

\[ B_0 = B_{\text{eq}} \rho_{i,\text{eq}} / \lambda_{\text{mfp}} \approx 10^{-18} \text{ G}. \]

The turbulence is characterised by \( \text{Re} \approx 60 \) and

\[ \frac{L}{U} \approx 10^9 \text{ yr} \text{ (outer timescale),} \]  
\[ \frac{L}{U} \text{Re}^{-1/2} \approx 10^8 \text{ yr} \text{ (viscous timescale),} \]  
\[ l_\nu \approx L \text{Re}^{-3/4} \approx 10 \text{ kpc} \text{ (viscous cutoff).} \]

3. A model of self-accelerating cluster dynamo

We assume that the net effect of the instabilities is to change the effective collision frequency of ions:

\[ \nu_{\text{eff}} = \nu_{ii} + \nu_{\text{scatter}}(B), \]  

where \( \nu_{\text{scatter}} \) is given by Eq. 8. If we use Eqs. 14-15 with \( \nu_{ii} \) replaced by \( \nu_{\text{eff}} \), we find the following equation for the anisotropy ratio

\[ \frac{1}{\Delta^2} = \frac{1}{\Delta_0^2} + \frac{1}{\epsilon} (\Delta - 2B^2)^\alpha B, \]

where magnetic field is in units of \( B_{\text{eq}} \) [Eq. 12] and

\[ \Delta_0 = \left[ \frac{L}{(v_{\text{thi}})^3 \lambda_{\text{mfp}} / L} \right]^{1/2} \approx 0.01, \]  
\[ \epsilon = \frac{L}{v_{\text{thi}}} \left( \frac{B_{\text{eq}}}{\rho_{i,\text{eq}} L} \right) \lambda_{\text{mfp}} \Delta_0^2 \approx 10^{-17}. \]

Our results will not be very sensitive to the value of \( \alpha \). In what follows, we shall use \( \alpha = 3/2 \) for specific estimates.

As the instabilities grow much faster than the field changes, we can assume that, as far as the field is concerned, the adjustment of the collision frequency occurs instantaneously. The growth of the magnetic field satisfies [from Eqs. 4 and 5 with \( \nu_{ii} \rightarrow \nu_{\text{eff}} \)]

\[ \frac{dB}{dt} = \frac{1}{\Delta(B)} \left( \frac{v_{\text{thi}}}{U} \right)^2 \frac{L}{\lambda_{\text{mfp}}} \alpha/2. \]

For \( \alpha = 3/2 \), we have \( B_1 \approx 10^{-15} \text{ G}. \)

3.1. Exponential stage

Let us start with an arbitrarily small seed field. As long as the term in Eq. (22) dominates, the effect of the instabilities is negligible and the field grows exponentially with the characteristic e-folding time \( \sim \Delta_0 \), which in dimensional units translates to the inverse rms rate of strain \( \sim 10^8 \text{ yr} \) [Eq. 19]. This initial dynamo stage continues until the second term in Eq. (22) becomes comparable to the first, which occurs at the field strength

\[ B_1 = \frac{\epsilon}{\Delta_0^{2+\alpha}} = B_{\text{eq}} \frac{\rho_{i,\text{eq}}}{\lambda_{\text{mfp}}} \left[ \frac{v_{\text{thi}}}{U} \right]^3 \frac{L}{\lambda_{\text{mfp}}} \alpha/2. \]

For \( \alpha = 3/2 \), we have \( B_1 \approx 10^{-15} \text{ G}. \)

3.2. Explosive stage

For \( B \gg B_1 \), we may approximate the solution of Eq. (22) by

\[ \Delta(B) \approx \left( \frac{\epsilon}{B} \right)^{1/(2+\alpha)}. \]

Upon substitution into Eq. (25), this gives an explosively growing field:

\[ B(t) = \epsilon \left( \frac{2 + \alpha}{c t_c - t} \right)^{2+\alpha}, \]

where, taking \( B(0) = B_1 \), we get \( t_c = (2 + \alpha) \Delta_0 \). During this very fast stage, the anisotropy drops precipitously, while the effective Reynolds number and, therefore, the rms rate of strain, grow:

\[ \text{Re}_{\text{eff}} = \left( \frac{U}{v_{\text{thi}}} \right)^4 \frac{1}{\Delta(B)^2}. \]

Hence the self-accelerating nature of the dynamo. Eventually (in fact, very soon), the field grows to be strong enough so that \( B^2 \) is comparable to \( \Delta(B) \), the second term in Eq. (22) starts dropping again and the \( \Delta \) stops decreasing and starts increasing. The minimum value of \( \Delta \) is reached at \( B \sim B_2 \), where

\[ B_2 = \epsilon^{1/(5+2\alpha)} \text{ Re}_{\text{eff}} \left( \frac{U}{v_{\text{thi}}} \right)^4 \frac{1/(5+2\alpha)}{\lambda_{\text{mfp}}} \lambda_{\text{mfp}} \Delta_0^{2+\alpha} \sim 10^{-17}. \]

For \( \alpha = 3/2 \), we have \( B_2 \approx 10^{-7} \text{ G}. \) The time it takes to achieve this is

\[ t \approx t_c - (2 + \alpha) \epsilon^{2/(5+2\alpha)} \approx t_c \approx 10^8 \text{ yr}. \]

This is to be compared with the conventional dynamo (e.g., Subramanian, Shukurov & Haugen, 2005), which has the e-folding time of \( 10^8 \text{ yr} \) and would, therefore, need around 2 Gyr to achieve an amplification factor of \( 10^8 \).

\[ \text{Subramanian, Shukurov & Haugen, 2005,} \]
3.4 Nonlinear stage

3.3 Algebraic stage

When \( B > B_2 \), \( \Delta \) stays just above \( 2B^2 \). Computing the small correction, we obtain the following solution of Eq. (22):

\[
\Delta(B) \approx 2B^2 + \left[ \frac{\epsilon}{B} \left( \frac{1}{4B^4} + \frac{1}{\Delta^2} \right) \right]^{1/6}. \tag{32}
\]

Solving Eq. (22) with \( \Delta \approx 2B^2 \) and \( B(0) = B_2 \), we get

\[
B(t) = \sqrt{B_2^2 + t}. \tag{33}
\]

This is a rather slow growth but it only lasts until the first term in Eq. (22) is again dominant: the transition is, in fact, sharp because once \( \Delta < 2B^2 \), the instabilities are shut down and we have to set \( \nu_{\text{scatter}} = 0 \) for all greater values of \( B \). The field at which this happens is

\[
B_3 = \left( \frac{\Delta_0}{2} \right)^{1/2} = \frac{B_{\text{eq}}}{\sqrt{2}} \left[ \left( \frac{U}{v_{\text{thi}}} \right)^3 \frac{\lambda_{\text{mfp}}}{L} \right]^{1/4} \approx 1 \mu G. \tag{34}
\]

The time to get there is \( t \approx B_3^2 = \Delta_0/2 \approx 10^8 \) yr again. The amplification factor since the end of the explosive stage is \( \sim 10 \).

Note that if we compare \( \rho_i \) with the effective mean free path \( \lambda_{\text{mfp},\text{eff}} = v_{\text{thi}}/\nu_{\text{eff}} \), we get

\[
\frac{\rho_i(B)}{\lambda_{\text{mfp},\text{eff}}(B)} = \frac{B_{\text{eq}}}{B} \frac{\epsilon}{\Delta(B)^2}. \tag{35}
\]

which is readily seen to decrease throughout all of the dynamo stages discussed above. Therefore, the magnetised-plasma assumption, once true at \( B = B_0 \), is never broken. For example, at the point of the largest effective Reynolds number (\( \text{Re} \sim 10^6 \) at end of the explosive regime, \( B = B_2 \)) we have \( \lambda_{\text{mfp},\text{eff}} \approx 0.03 \) pc and \( \rho_i \approx 10^6 \) km.

3.4 Nonlinear stage

So far, the dynamo we have considered has been kinematic, i.e., the magnetic field has remained too weak to exert a back reaction on the flow. The condition for this to be true is that the field energy remains smaller than the energy of the smallest-scale turbulent motions that can change the field, i.e., the viscous-scale motions (see Schekochihin et al. 2004). In Kolmogorov turbulence, their velocity is \( U \text{Re}^{-1/4} \), whence, using Eqs. (29) and (14), we find that the condition for the nonlinearity to become important is \( B^2 > 10^6 \Delta(B) \). Thus, all through the algebraic regime, the magnetic field hovers around this threshold. After instabilities are stabilised, it can finally exceed it by growing above \( B_3 \). For \( B > B_3 \), we have to model the effect of the back reaction. How to do this is a highly nontrivial problem for which a definitive solution has not so far been found. Here, we shall use the model proposed by Schekochihin et al. (2002), which is constructed in a somewhat similar spirit as the self-accelerating dynamo model above. The main idea is that the growing magnetic field gradually suppresses the ability of turbulent eddies to stretch it. At any given point in time, all eddies whose energy is less than the energy of the field are thus suppressed. Final saturated state is reached when the energy of the field is comparable to the energy of the outer-scale eddies, i.e., to the total energy of the fluid turbulence:

\[
\frac{B_4}{B_{\text{eq}}} = \frac{U}{v_{\text{thi}}} \approx 6 \mu G. \tag{36}
\]

To model this process, we assume that the magnetic field evolves according to

\[
\frac{1}{B} \frac{dB}{dt} = \gamma(B), \tag{37}
\]

with the \( B \)-dependent growth rate (Schekochihin et al. 2002)

\[
\gamma(B) = \frac{U}{L \text{Re}^{1/2}} \left[ 1 - \frac{1}{(1 + B_4^2/B_3^2)^{1/2}} \right]^{-1/2} \times \left[ \frac{1}{(1 + B^2/B_3^2)^2} - \frac{1}{(1 + B^2/B_4^2)^2} \right]^{1/2}. \tag{38}
\]

\[\text{For numerical evidence in support of the model, see Schekochihin et al.} \ (2004).\]
It is not hard to see that, for $B_3 \ll B \ll B_4$, $\gamma(B) \propto 1/B^2$, which again gives the magnetic field growing as $B \sim \sqrt{t}$. Thus, the slow algebraic growth continues as the field exceeds $B_3$ until it saturates at $B = B_4$. Again, it does not take long because the inertial range in clusters is very narrow and the field does have to grow very much to get from $B_3$ to $B_4$.

4. Conclusions

All stages of the evolution of the magnetic field described in the previous section are illustrated in Fig. 1, which presents the numerical solution of Eqs. (25) and (22) [followed by Eqs. (47, 48) in the nonlinear stage]. We see that the field grows by a factor of about $10^{12}$ over a time period of $\sim 2$ Gyr. Most of the growth occurs over a short fraction of this time during the explosive self-accelerating dynamo stage, which takes the field from any value above $\sim 10^{-15}$ G to $\sim 10^{-7}$ G in a cosmologically short time of about $10^{10}$ yr. After that, the field does not have to grow further for a very long time before it reaches the strength of several $\mu$G consistent with the observational data. The importance of the self-accelerating dynamo mechanism we have proposed is that it provides for field amplification to the observed strength of the seed field over a short period of time independent of the precise value of the seed field or, indeed, of exactly how many Gyr a cluster has been around. This removes the need to envision ways of generating seed fields satisfying the often quite large lower bounds that theories (cited in the Introduction) and numerical simulations (e.g., Roettiger, Stone & Burns 1999, Dolag, Bartelmann & Lesch 1999, 2002, Dolag et al. 2005, Brüggen et al. 2008) infer to be necessary for magnetic fields of observed strength to be produced by the conventional turbulent dynamo over the cluster lifetime. Since the amplification time is so short, it is also probably not crucial whether the cluster turbulence responsible for amplification is best modelled as forced or decaying (cf. Subramanian et al. 2005). We conclude that, no matter what the seed field is, the random motions in the ICM will have no difficulty in amplifying it to the observed level in a fraction of the cluster lifetime.

Finally, we must alert the reader to an important issue, which we have ignored above. This concerns the spatial structure of the cluster fields. Theory and simulations of the small-scale dynamo in the MHD description indicate that the dynamo-generated fields have a folded structure with direction reversals at the resistive scale $\sim L \Omega_{\text{turb}}^{-1/2}$ (Scheukhin et al. 2004, Brandenburg & Subramanian 2005). In clusters, this scale turns out to be extremely small ($\sim 10^4$ km), while analysis of the observational data suggests that the typical reversal scale is $\sim 1$ kpc (Voet & Enßlin 2003, 2005). It is likely that the reversal scale is, in fact, set by some form of anomalous resistivity associated with electron scattering by the plasma instabilities discussed above. We relegate further development of this idea to future work.

Acknowledgements. It is a pleasure to acknowledge helpful discussions with T. Enßlin, G. Hammett, R. Kulpsrud, E. Quataert, P. Sharma, A. Waelkens. A.A.S. was supported by a UKAFF Fellowship. This work has benefited from research funding from the EU Sixth Framework Programme under RadioNet R113CT 2003 5058187 and from the US DOE Center for Multiscale Plasma Dynamics, Fusion Science Center Cooperative Agreement ER54785.

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