A possible origin of magnetic fields in galaxies and clusters: strong magnetic fields at $z \sim 10$?

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ABSTRACT

We propose that strong magnetic fields should be generated at shock waves associated with the formation of galaxies or clusters of galaxies by the Weibel instability, an instability in collisionless plasmas. The strength of the magnetic fields generated through this mechanism is close to the order of those observed in galaxies or clusters of galaxies at present. If the generated fields do not decay rapidly, this indicates that strong amplification of magnetic fields after the formation of galaxies or clusters of galaxies is not required. This mechanism could have worked even at a redshift of $\sim 10$, and therefore the generated magnetic fields may have affected the formation of stars in protogalaxies. This model will partially be confirmed by future observations of nearby clusters of galaxies. Mechanisms that preserve the magnetic fields for a long time without considerable decay are discussed.

Key words: instabilities – magnetic fields – galaxies: clusters: general – galaxies: general.

1 INTRODUCTION

The question of the origin of galactic magnetic fields is one of the most challenging problems in modern astrophysics. It is generally assumed that magnetic fields in spiral galaxies are amplified and maintained by a dynamo through rotation of the galaxies (Widrow 2002). The dynamo requires seed fields to be amplified. However, observations of microgauss fields in galaxies at moderate redshifts strongly constrain the lower boundary of the seed fields (Athreya et al. 1998). Moreover, magnetic fields are also observed in elliptical galaxies and galaxy clusters, in which rotation cannot play a central role as the dynamo mechanism (Clarke, Kronberg & Böhringer 2001; Widrow 2002; Vallée 2004).

The Weibel instability is another mechanism to generate strong magnetic fields (Weibel 1959; Fried 1959). This instability is driven in a collisionless plasma, or a tenuous ionized gas, by the anisotropy of the particle velocity distribution function (PDF) of the plasma. When the PDF is anisotropic, currents and then magnetic fields are generated in the plasma, so that the plasma particles are deflected and the PDF becomes isotropic (Medvedev & Loeb 1999). Through the instability, the free energy attributed to the PDF anisotropy is transferred to magnetic field energy. This instability does not need seed magnetic fields. It can be saturated only by non-linear effects, and thus the magnetic fields can be amplified to very high values.

This instability has been observed directly in recent laser experiments (Wei et al. 2002). In astrophysical plasmas, the instability is expected to develop at shocks or at steep temperature gradients, where the PDF is anisotropic. Examples of such sites are pulsar winds, shocks produced by γ-ray bursts, jets from active galactic nuclei (AGNs), cosmological shocks, and cold fronts (contact discontinuities between cold and hot gas) in clusters of galaxies (Kazimura et al. 1998; Medvedev & Loeb 1999; Nishikawa et al. 2003; Schlickeiser & Shukla 2003; Okabe & Hattori 2003). Although the instability was found in 1959, its non-linear nature has prevented us from understanding its long-term evolution. Recently, however, as computer power has increased, detailed particle simulations of plasmas have been initiated, and they have revealed the evolution of magnetic fields even after saturation of the instability (Silva et al. 2003; Medvedev et al. 2005). Based on these results, we consider the generation of magnetic fields at galaxy- and cluster-scale shocks through the Weibel instability at the formation of galaxies (both ellipticals and spirals) and clusters. We use the cosmological parameters $\Omega_0 = 0.3, \lambda = 0.7$, the Hubble constant $H_0 = 70$ km s$^{-1}$ Mpc$^{-1}$, and $\sigma_8 = 0.9$.

2 MODELS

2.1 Generation of magnetic fields at shocks

In the vicinity of the shock front of a collisionless shock, particles from upstream are mixed up with those downstream, and an anisotropy of the PDF will be generated in the plasma. As discussed by Schlickeiser & Shukla (2003), the particles from the upstream region will first be affected by the Langmuir instability. Since the Langmuir instability is a longitudinal electrostatic mode, the velocity component parallel to the shock normal will be thermalized,
while the other components will remain unaffected. Therefore, it is natural to assume that the thermal velocity parallel to the shock normal is of the order of the relative velocity between the upstream and downstream regions or of the order of the shock velocity $V_{\text{sh}}$, and those perpendicular to the shock normal are of the order of the thermal velocity of the upstream plasma. This velocity or temperature anisotropy should develop the Weibel instability.

For shocks in electron–proton plasmas, Schlickeiser & Shukla (2003) indicated that the temperature anisotropy for electrons is too small to drive the Weibel instability if $M \lesssim (m_e/m_p)^{1/2} = 43$, where $M = V_{\text{sh}}/v_{\text{th},p}$ is the shock Mach number, $v_{\text{th},p}$ is the proton thermal velocity in the upstream region, and $m_e$ and $m_p$ are the proton and electron masses, respectively. However, if $M \gtrsim 2$ (note that $M$ provides a measure of the anisotropy in protons), the temperature anisotropy for the protons is large enough to drive the Weibel instability; this is the case even if the PDF of the electrons is completely isotropic (see Appendix A).

The magnetic field strength reaches its maximum when the Weibel instability saturates, and the saturation level would be given as follows. As the Weibel instability develops, magnetic fields are generated around numerous current filaments (Medvedev et al. 2005; Kato 2005). The instability saturates when the generated magnetic fields eventually interrupt the current in each filament, or in other words when the particles’ gyroradii in the excited magnetic fields are comparable to the characteristic wavelength of the excited field (Medvedev & Loeb 1999). For electron–positron plasmas, Kato (2005) showed that the typical current at and after saturation is given by the Alfvén current (Alfvén 1939). In terms of the Alfvén current, the saturated magnetic field strength is given by

$$B_{\text{sat}} \sim \sqrt{2} I_A/(Rc),$$

(1)

where $R$ is the typical radius of a current filament at saturation and $I_A \equiv (v_{\text{th},e}/c)m_e e^2/E$ is the Alfvén current (for non-relativistic cases).

Here, $v_{\text{th},e}$ is the electron thermal velocity in the direction of the higher temperature. For strongly anisotropic cases (i.e. the particle limit cases discussed by Kato (2005)), the radius at saturation is given by $R \sim 2c/\omega_{\text{pe}}$, where $\omega_{\text{pe}} \equiv (4\pi n_0 e^2/m_e)^{1/2}$ is the electron plasma frequency, and the saturated magnetic field strength is given by

$$B_{\text{sat}} \sim 0.5 v_{\text{th},e}(4\pi n_0 m_e)^{1/2},$$

(2)

where $n_0$ is the electron number density.

For the electron–proton plasmas we consider below, currents carried by heavier protons will generate stronger magnetic fields at saturation. It is natural to assume that saturation for the proton currents is determined by the Alfvén current defined for protons, $I_A' \equiv (v_{\text{th},p}/c)m_p e^2/E$, where $v_{\text{th},p}$ is the proton thermal velocity in the direction of the higher temperature. In a shock wave, this leads to

$$B_{\text{sat}} \sim 0.5 V_{\text{sh}}(2\pi n_0 m_p)^{1/2},$$

(3)

where $n_0$ is the proton number density and we have used the assumption that $v_{\text{th},p} \sim V_{\text{sh}}$. This expression is valid when the PDF of the protons has a strong anisotropy or $M \gtrsim 2$ (Kato 2005). Note that the energy density of the saturated magnetic fields attains a sub-equipartition level with the particle kinetic energy. The numbers in parentheses in equations (2) and (3) are different by a factor of 2, because in electron–positron plasmas both electrons and positrons produce currents at the same time, while in electron–proton plasmas only protons contribute to currents when the proton Weibel instability develops. In electron–proton plasmas, if the electron anisotropy is sufficiently large, first electrons generate magnetic fields, which saturate at the early stage of the instability, and then protons generate magnetic fields, which leads to the second saturation. Such an evolution was also observed in numerical simulations (Frederiksen et al. 2004).

The magnetic fields generated at the shock front will be convected downstream at the fluid velocity (in the shock rest frame). At least for a short period after saturation, the magnetic field strength decreases as the current filaments merge together, because the current in each filament is limited to the proton Alfvén current while the size of the current filaments increases (Kato 2005). However, the long-term evolution is still an open question. Medvedev et al. (2005) indicated that the field should reach an asymptotic value beyond which it will not decay, although their model can be applied only when currents are not dissipated. Recent particle simulations have also shown that the final magnetic field strength is given by $B_{\text{sat}} = \eta_{\text{loc}} B_{\text{sat}}$, where $\eta_{\text{loc}} \sim 0.01$ (Silva et al. 2003; Medvedev et al. 2005). In the following, we assume that the magnetic fields $B_{\text{sat}}$ are preserved. The very long-term evolution of the magnetic fields, for which the results of Medvedev et al. (2005) and Silva et al. (2003) cannot be directly applied, will be discussed in Section 4.

We note that, unfortunately, at present most simulations treat electron–positron shocks and there are virtually no simulations that reveal the evolution of collisionless electron–proton shocks satisfactorily (in terms of box sizes, duration times and so on), because of the lack of computational power. Therefore, the model stated above and used in this paper should be confirmed by direct numerical simulations in the future.

2.2 Shock formation

According to the standard hierarchical clustering scenario of the Universe, an initial density fluctuation of dark matter in the Universe gravitationally grows and collapses; its evolution can be approximated by that of a spherical uniform overdense region (Gunn & Gott 1972; Peebles 1980). The collapsed objects are called ‘dark haloes’ and the gas in these objects later forms galaxies or clusters of galaxies. At collapse, the gas is heated by ‘virial shocks’ to the virial temperature of the dark halo, $T_{\text{vir}} = GM/(2\pi c)\rho_{\text{crit}}$, where $G$ is the gravitational constant, and $M$ and $\rho_{\text{vir}}$ are the mass and virial radius of the dark halo, respectively. The relation between the virial radius and the virial mass of an object is given by

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

(4)

where $\rho_{\text{crit}}(z)$ is the critical density of the Universe, and $\Delta_c(z)$ is the ratio of the average density of the object to the critical density at redshift $z$. The critical density depends on redshift because the Hubble constant depends on that, and it is given by

$$\rho_{\text{crit}}(z) = \rho_{\text{crit},0} \frac{\Omega_m(1+z)^3}{\Omega(z)},$$

(5)

where $\rho_{\text{crit},0}$ is the critical density at $z = 0$, and $\Omega(z)$ is the cosmological density parameter given by

$$\Omega(z) = \frac{\Omega_m(1+z)^3}{\Omega_0(1+z)^3 + \lambda},$$

(6)

for a flat Universe with non-zero cosmological constant. The ratio $\Delta_c(z)$ is given by

$$\Delta_c(z) = 18\pi^2 + 82x - 39x^2,$$

(7)

for a flat Universe (Bryan & Norman 1998), where the parameter $x$ is given by $x = \Omega(z) - 1$. The virial shocks form at $r \approx r_{\text{vir}}$, and the velocity is $v_{\text{sh}} \approx \sqrt{GM/r_{\text{vir}}}$. 

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In addition, recent cosmological numerical simulations have shown that ‘large-scale structure (LSS) shocks’ form even before collapse (Cen & Ostriker 1999; Miniati et al. 2000; Davé et al. 2001; Ryu et al. 2003). They form at the turnaround radius \( r_{\text{ta}} \sim 2r_{\text{vir}} \), the point at which the density fluctuation breaks off from the cosmological expansion. For simplicity, we assume that \( r_{\text{ta}} = 2r_{\text{vir}} \), which is close to the self-similar infall solution for a particular mass shell in an Einstein–de Sitter universe \( r_{\text{ta}} = (1/0.56)r_{\text{vir}} \) (Bertschinger 1985). The gas that later forms a galaxy or a cluster passes two types of shocks; the gas passes first the outer LSS shock and then the inner virial shock. The typical velocity of the LSS shocks is \( V_{\text{sh}} \approx H(z)r_{\text{ta}} \) (Furlanetto & Loeb 2004), where \( H(z) \) is the Hubble constant at redshift \( z \), and \( r_{\text{ta}} \) is the physical radius that the region would have had if it had expanded uniformly with the cosmological expansion. The temperature of the post-shock gas is \( T_{s} \approx (3/16)(\mu m_{p}/k_{B})V_{\text{sh}}^{2} \), where \( \mu m_{p} \) is the mean particle mass, and \( k_{B} \) is the Boltzmann constant. Note that, although the model of Furlanetto & Loeb (2004) has been compared with numerical simulations at \( z \approx 0 \), it has not been at high redshifts; it might have some ambiguity there. We do not consider mergers of objects that have already collapsed as the sites of magnetic field generation because the Weibel instability applies only to initially unmagnetized or weakly magnetized plasmas; at the merger, collapsed objects just bring their magnetic fields to the newly born merged object.

Since the Weibel instability develops in ionized gas (plasma), we need to consider the ionization history of the Universe. After the entire Universe is ionized by stars and/or AGNs (\( z \lesssim 8 \)), magnetic fields are first generated at the LSS shocks. In this case, we do not consider the subsequent generation of magnetic fields at the inner virial shocks, because the strength is at most comparable to that of the magnetic fields generated at the LSS shocks. On the other hand, when the Universe is not ionized (\( z \gtrsim 8 \)), the Weibel instability cannot develop at the outer LSS shocks. However, if the LSS shocks heat the gas (mostly hydrogen) to \( T_{s} \gg 10^{4} \text{K} \) and ionize it, the instability can develop at the inner virial shocks.

In this case, the gas ionized at the LSS shock may recombine before it reaches the virial shock. The recombination time-scale is given by

\[
\tau_{\text{rec}} = \frac{1}{\alpha n_{e}} \approx 1.22 \times 10^{3} \text{yr} \frac{1}{y} \left( \frac{T}{10^{4} \text{K}} \right)^{0.7} \left( \frac{n_{\text{H}}}{\text{cm}^{-3}} \right)^{-1},
\]

where \( \alpha \) is the recombination coefficient, \( T \) is the gas temperature, \( y \) is the ionization fraction, and \( n_{\text{H}} \) is the hydrogen density (Shapiro & Kang 1987). If we assume that \( \tau_{\text{rec}} = \tau_{\text{dyn}} \), where \( \tau_{\text{dyn}} \approx (1/2)r_{\text{ta}}/V_{\text{sh}} \) is the time-scale in which the gas moves from the LSS shock to the virial shock, the ionization rate when the gas reaches the virial shock is

\[
y \approx \left( \frac{\tau}{1.22 \times 10^{3} \text{yr}} \right)^{-1} \left( \frac{T}{10^{4} \text{K}} \right)^{0.7} \left( \frac{n_{\text{H}}}{\text{cm}^{-3}} \right)^{-1}
\]

for \( y < 1 \). We found that \( y \approx 0 \) for \( z \gtrsim 9 \), and the minimum value when the generation of magnetic fields is effective (\( z \lesssim 12 \), see Section 3) is \( y \approx 0.3 \). When \( y < 1 \), we simply replace \( n_{\text{H}} \) in equation (3) with \( y n_{\text{H}} \). For the temperature, we assumed that \( T = T_{s} \) in equation (9). Since the magnetic fields do not depend much on the temperature \( (B_{\text{sh}} \propto y^{0.5} \propto T^{0.35}) \), they do not change much even when radiative cooling reduces the temperature; at least \( B_{\text{sh}} \) does not change by many orders of magnitude. It was shown that the ionization rate just behind a shock is \( y \approx 0.1 \), if the shock velocity is relatively small \( (V_{\text{sh}} \sim 40 \text{km s}^{-1}) \); Shapiro & Kang 1987; Susa et al. 1998). If the shock velocity is larger, \( y \sim 1 \). In our calculations, the velocity of the LSS shocks is \( V_{\text{sh}} \gtrsim 40 \text{km s}^{-1} \), when the generation of magnetic fields is effective. Thus, \( y \) just behind the shocks is at least comparable to that obtained through the condition \( \tau_{\text{rec}} \sim \tau_{\text{dyn}} \), and incomplete ionization does not greatly affect the results shown in the next section (see also Section 4).

3 RESULTS AND DISCUSSION

Fig. 1 shows the typical mass of objects, \( M \), as a function of redshift \( z \); the labels \( 1\sigma, 2\sigma \) and \( 3\sigma \) indicate the amplitudes of initial density fluctuations in the Universe from which the objects form, on the assumption of the cold dark matter (CDM) fluctuations spectrum (Barkana & Loeb 2001). Fig. 2 shows the downstream temperature of magnetic fields generated at the LSS shocks. Possible origin of magnetic fields...
at the virial shock \( T_{\text{vir}} \) and that at the LSS shock \( T_s \) for the objects; \( T_{\text{vir}} \) is always larger than \( T_s \). The ratio \( T_{\text{vir}}/T_s \) indicates that \( M \sim 4 \) for the virial shock. For the LSS shock, \( M \gg 1 \), because the gas outside the shock is cold. Thus, equation (3) can be applied to both shocks. In Fig. 3, we present the strength of magnetic fields \((B_r)\) at a scale of \( r_{\text{vir}} \) for the collapsed objects. We assume that the entire Universe is reionized at \( z = 8 \). Thus, for \( z > 8 \), magnetic fields are generated only at the virial shocks if \( T_s > 10^4 \) K. We assume that \( B_r = B_t \) and plot the lines only when \( T_s > 10^4 \) K. The recombination is effective at \( z > 10 \) for the \( 3\sigma \) model. On the other hand, for \( z < 8 \), the magnetic fields are generated at the LSS shocks. We consider the compression of the fields while the size of the gas sphere decreases from \( r = r_{\text{vir}} \) to \( r = r_{\text{vir}} \), and thus we assume that \( B_r = B_t \). Moreover, we plot the lines only for \( T_{\text{vir}} > 2 \times 10^4 \) K, because below this temperature gas infall is suppressed by photoionization heating (Efstathiou 1992; Furlanetto & Loeb 2004).

In Fig. 3, the strength of the magnetic fields generated by protons reaches \( \sim 10^{-8} - 10^{-7} \) G and is very close to the values observed in nearby galaxies and clusters of galaxies \( \sim 10^{-8} \) G; Clarke et al. 2001; Widrow 2002; Vallée 2004). On a galactic scale, the gas sphere may further contract to \( r < r_{\text{vir}} \) because of radiative cooling. If magnetic fields are frozen in the gas, the strength exceeds \( \sim 10^{-6} \) G. However, if this happens, the magnetic energy exceeds the thermal or kinetic energy of the gas. As a result, magnetic reconnection may reduce the strength. In fact, equipartition between the magnetic energy density and the thermal or kinetic energy density appears to hold in the Galaxy (Beck et al. 1996).

The strong magnetic fields shown in Fig. 3 indicate that strong amplification of magnetic fields, such as dynamo amplification, is not required after the formation of galaxies and clusters. This is consistent with the observations of galactic magnetic fields at \( z > 2 \) (Athreya et al. 1998). Future observations of higher-redshift galaxies would discriminate between our model and strong dynamo amplification models; the latter predict much weaker magnetic fields at higher redshifts. Moreover, since the predicted galactic magnetic fields are comparable to those at present, they might have affected the formation of stars in protogalaxies. Fig. 3 also shows that our model naturally explains the observational fact that the magnetic field strengths of galaxies and galaxy clusters fall in a small range (a factor of 10). Since our model predicts that magnetic fields are generated around objects, magnetic fields in intergalactic space are not required as the seed or origin of galactic magnetic fields.

Some recent numerical simulations indicated that gas is never heated to \( \sim T_{\text{vir}} \) for less massive objects because radiative cooling is efficient and the shocks forming at \( z \geq z_{\text{vir}} \) are unstable (Birnboim & Dekel 2003; Kereš et al. 2005). This effect may be important for \( M \lesssim 10^{12} \) M\(_\odot\). If it is correct, the generation of magnetic fields is only effective for \( z \lesssim 5 \) (the \( 3\sigma \) curve in Fig. 1). However, even for \( M \lesssim 10^{12} \) M\(_\odot\), multiple shocks may form at the inner halo when the cold gas reaches there and they collide with each other (Kereš et al. 2005). Therefore, magnetic fields may be created there. The details of these shocks could be studied by high-resolution numerical simulations.

Although it would be difficult to directly observe the generation of magnetic fields through the Weibel instability for distant high-redshift galaxies, it would be easier for nearby clusters of galaxies. Since clusters are now growing, LSS shocks should be developing outside of the virial radii of the clusters (Miniati et al. 2000; Ryu et al. 2003). Since particles are often accelerated at shocks, the synchrotron emission from the accelerated particles could be observed with radio telescopes with high sensitivity at low frequencies (Keshet, Waxman & Loeb 2004). The total non-thermal luminosity \( \text{[synchrotron luminosity plus inverse Compton scattering with cosmic microwave background (CMB) photons]} \) is estimated as

\[
L_u \approx \frac{\epsilon f M}{r_{\text{vir}}/m_p} V_{\text{vir}}^2, \tag{10}
\]

where \( \epsilon \) is the acceleration efficiency and \( f \) is the gas fraction of a cluster. If we assume \( \epsilon = 0.03 \) (Taniguchi et al. 1998; Muraiishi et al. 2000) and \( f = 0.15 \) (Mohr, Mathiesen & Evrard 1999), the maximum value of \( L_u \) is \( \sim 10^{41} \) erg s\(^{-1}\). Since the energy density of magnetic fields \( u_B \) is smaller than that of the CMB \( u_{\text{CMB}} \), most of the non-thermal luminosity \( (L_u) \) is attributed to inverse Compton scattering such as \( u_{\text{CMB}} L_u/(u_B + u_{\text{CMB}}) \), which may have been detected in the hard X-ray band (\( \gtrsim 20 \) keV; Fusco-Femiano et al. 2004). Thus, the synchrotron radio luminosity is \( u_B L_u/(u_B + u_{\text{CMB}}) \lesssim 10^{40} \) erg s\(^{-1}\). Some of the diffuse radio sources observed in the peripheral cluster regions (‘radio relics’) may be this emission (Govoni & Feretti 2004). Since the Weibel instability generates magnetic fields on the plane of the shock front, the synchrotron emission should be polarized perpendicular to the shock front (Medvedev & Loeb 1999), which is actually observed for some radio relics (Govoni & Feretti 2004). The synchrotron emission will tell us the positions of the LSS shocks, if they exist. If magnetic fields are generated there, they should be observed only downstream of the shock. This may be confirmed through Faraday rotation measurements of radio sources behind the cluster for both sides of the shock, if the coherent length of the fields increases sufficiently (see Section 4).

4 ON THE LONG-TERM EVOLUTION OF MAGNETIC FIELDS

The model proposed here has large uncertainties, especially on the very long-term (say gigayear) evolution of magnetic fields. The typical time-scale of the Weibel instability until saturation is given by the inverse of the proton plasma frequency, \( \omega_{\text{pe}} = (4\pi n_p e^2/m_p)^{1/2} \). However, it is evident that this time-scale is much smaller than the cosmological time-scale we have considered in this paper. The typical scalelength of the saturated fields is given by the proton inertial.
length $c/\omega_{pe}$, and it is also much smaller than a kiloparsec scale. Thus, there is a ‘missing link’ between the generated fields and those observed in galaxies and clusters at present.

If the observed magnetic fields have their origin in the Weibel instability, there should be another mechanism that takes over the instability. If the saturation and current evolution model of Kato (2005) is applied to proton currents, the magnetic field strength should decrease after saturation as

$$B = B_{sat}(R/R_{c})^{-1}, \quad (11)$$

where $R$ is the radius of a filament, and the filament radius at saturation, $R_{c}$, is defined for protons ($R_{c} \sim 2c/\omega_{pe}$). This is because the current strength is limited to the proton Alfvén current $I_{A}$ while the radius of current filaments increases through mergers of the filaments. Since $R \sim 10^{10} \text{ cm}$ for $n_{p} \sim 10^{-3} \text{ cm}^{-3}$ (at $r \sim r_{eq}$ for a cluster at $z \sim 0$), the magnetic field strength should be reduced by a factor of $\sim 10^{11}$, if relation (11) holds even for a long-term evolution to a kiloparsec ($\sim 10^{21}$ cm) scale field.

One possibility to overcome this difficulty is that the currents are carried in another form at later times. At the early times of evolution, the state of the plasma is in a kinetic regime in which the characteristic scale of magnetic fields is comparable to or smaller than the Larmor radius of the particles. Probably, at later times, the former would increase via current mergers, and would become larger than the latter. After this situation is realized, the straight currents generated by the Weibel instability, which are limited to the Alfvén current, might be superseded by a kind of drift currents, and the plasma might behave as a magnetohydrodynamic (MHD) fluid. In this case, mergers of cylindrical structures, which were originally the current filaments, would occur through the reconnection process as usually considered for MHD fluids, and therefore the magnetic field strength would not decrease considerably unless the diffusion of magnetic fields due to collisions between charged particles or collisions with neutrals becomes effective.\(^1\) This would allow the magnetic fields to survive for a long time. For example, in an MHD fluid, for a scale of $L \sim 10^{10}$ cm and a temperature of a few keV, the dissipation time-scale of magnetic fields through the collisions between charged particles is $t_{diss} \sim 10^{9}$ yr (Spitzer 1962), which is the dynamical time-scale of a cluster. If current mergers make $L$ larger, then $t_{diss}$ also becomes larger. The numerical simulations performed by Silva et al. (2003) showed that magnetic fields do not decay much after saturation. This might reflect the transition to an MHD fluid. Moreover, an MHD inverse cascade mechanism (e.g. Vishniac & Cho 2001) could provide another process to make larger-scale magnetic fields observed in galaxies and clusters ($\gtrsim$ kiloparsec). Since the ionization rate, $\chi$, is fairly large (Section 2.2), the ambipolar diffusion of magnetic fields (collisions with neutrals) can be ignored (e.g. equations 13–57 in Spitzer 1978). At any rate, studies on the long-term evolution of magnetic fields are strongly encouraged.

5 CONCLUSIONS

In this paper, we have showed that the Weibel instability can generate strong magnetic fields in shocks around galaxies and clusters. The strength is comparable to those observed in galaxies and clusters at present. The mechanism could have worked even at $z \sim 10$. The results are based on the assumption that the magnetic fields generated by the Weibel instability are conserved for a long time. The validity of this assumption must be confirmed in future studies.

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1 We note that ‘fast’ reconnection (Petschek 1964) may occur on time-scales much shorter than those of the particle collisions.

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APPENDIX A: DISPERSION RELATION OF THE WEIBEL INSTABILITY

Here, we derive the dispersion relation of the Weibel instability in a plasma that consists of some species (or populations) of charged particles. In the following, each species is denoted by a label \( s \) (\( s = e \) for electron, \( p \) for proton, and so on), and the mass, charge and number density of the species are denoted by \( m_s \), \( q_s \) and \( n_s \), respectively. We assume here that each species has a bi-Maxwellian distribution

\[
f_\mathbf{0}^{(s)}(v) = \frac{n_s}{(2\pi)^{3/2}\alpha_s\sigma_s^3} \exp\left(-\frac{v_x^2 + v_y^2}{2\sigma_x^2} - \frac{v_z^2}{2\alpha_s^2\sigma_s^2}\right),
\]

where the thermal velocity is given by \( \sigma_s \) for the \( x \)- or \( y \)-directions, and by \( \alpha_s\sigma_s \) for the \( z \)-direction.

If the thermal velocity in the \( z \)-direction is larger than those in the other directions,\(^2\) that is, \( \alpha_s \geq 1 \), the linear dispersion relation of the Weibel mode, which is relevant to the \( z \)-component of the current density, is given as

\[
\omega^2 - (ck)^2 + \sum_s \omega_{ps}^2 \left[ \alpha_s^2 \zeta + \alpha_s^2 - 1 \right] = 0.
\]

\(^2\)Note that in this condition the direction of higher temperature is opposite to that of other authors (e.g. Weibel 1959; Davidson et al. 1972). Nevertheless, this would be more reasonable for anisotropy at shock waves.

where

\[
\omega_{ps} = \sqrt{\frac{4\pi q_s^2}{m_s}}, \quad \zeta_s = \frac{\omega}{k\sqrt{2}\sigma_s},
\]

and \( Z(\zeta) \) is the plasma dispersion function defined as

\[
Z(\zeta) = \frac{1}{\sqrt{\pi}} \int_{-\infty}^{\infty} \frac{1}{z - \zeta} e^{-\frac{z^2}{2}} dz.
\]

It is shown that the dispersion relation (A2) can have a purely positive imaginary solution for \( \omega \), i.e. a purely growing mode. Since the growth rate becomes zero at the maximum wavenumber of the unstable mode, \( k_{\text{max}} \), we can set \( \omega = 0 \) and \( \zeta \), \( Z(\zeta) = 0 \) at \( k = k_{\text{max}} \) in equation (A2) to obtain

\[
k_{\text{max}}^2 = \frac{1}{c^2} \sum_s \omega_{ps}^2 \left( \alpha_s^2 - 1 \right).
\]

It is evident that there is no unstable mode if \( k_{\text{max}}^2 \leq 0 \), while unstable modes exist if \( k_{\text{max}}^2 > 0 \). Explicitly, the condition for instability is given by

\[
\sum_s \omega_{ps}^2 \left( \alpha_s^2 - 1 \right) > 0.
\]

It should be noted that, in an electron–proton plasma for example, even when the electron distribution is completely isotropic (\( \alpha_e = 1 \)) at the initial time, unstable modes exist if the proton distribution is anisotropic (\( \alpha_p > 1 \)).

The above results show that, even if the Mach number of a shock is relatively small (\( \alpha_e \approx 1 \)), the magnetic fields generated by the anisotropy of the proton distribution slowly grow and reach the saturation value given by equation (3) (Kato 2005).

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