Color, Spin, and Flavor Diffusion in Quark-Gluon Plasmas

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In weakly interacting quark-gluon plasmas diffusion of color is found to be much slower than the diffusion of spin and flavor because color is easily exchanged by the gluons in the very singular forward scattering processes. If the infrared divergence is cut off by a magnetic mass, $m_{\rm mag} \sim \alpha_s T$, the color diffusion is $D_{\rm color} \sim [\alpha_s \ln(1/\alpha_s)T]^{-1}$, a factor α_s smaller than spin and flavor diffusion. A similar effect is expected in electroweak plasmas above M_W due to $W^{\pm}$ exchanges. The color conductivity in quark-gluon plasmas and the electrical conductivity in electroweak plasmas are correspondingly small in relativistic heavy ion collisions and the very early Universe.

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Transport of color degrees of freedom in a quark-gluon plasma has recently been found to be infrared sensitive [1] and thus differs from other transport processes as viscous and thermal flow and stopping and electrical conduction [2, 3], as well as energy degradation [4]. Color does not flow easily due to the transfer of color in the exchange of a colored gluon in the very singular forward collisions. As will be discussed here, this suppression of the color flow also applies to the color diffusion and color conductivity, which are infrared sensitive like the quark and gluon quasiparticle relaxation rates [5, 6].

Earlier studies of transport processes in relativistic quark-gluon and electron-photon plasmas found that the effect of Landau damping effectively led to screening of transverse interactions and gave the characteristic relaxation rates in transport processes. Transport coefficients for weakly interacting electron-photon and quark-gluon plasmas for both thermal plasmas [2–4] as well as degenerate ones [7] were calculated to leading logarithmic order. Generally the transport relaxation rates have the following dependence on interaction strength:

$$1/\tau_{\rm tr} \sim \alpha_s^2 \ln(1/\alpha_s) T . \qquad (1)$$

However, the quark and gluon quasiparticle damping rates, $1/\tau_p$, were not sufficiently screened by Landau damping for nonvanishing quasiparticle momentum, $\mathbf{p}$, and depend on an infrared cutoff, $m_{\rm mag} \simeq \alpha_s T$, so that [5, 6]

$$1/\tau_p^{(g)} = 3\alpha_s \ln(1/\alpha_s) T , \qquad (2)$$

to leading logarithmic order. Since the quasiparticle decay rates are not measurable transport coefficients the infrared sensitivity was not considered a serious problem. However, it was recently discovered [1] that diffusion of color in some abstract color space suffered from the same infrared divergence which led to the same color relaxation rates as the quasiparticle damping rates.

We will describe the two kinds of transport processes by calculating the flavor, spin, and color diffusion coefficients in a quark-gluon plasma within the Boltzmann kinetic equation

$$\left(\frac{\partial}{\partial t} + \mathbf{v}_{\mathbf{p}_1} \cdot \nabla_{\mathbf{r}} + \mathbf{F} \cdot \nabla_{\mathbf{p}_1}\right) n_1 = -2\pi\nu_2 \sum_{\mathbf{p}_2\mathbf{p}_3\mathbf{p}_4} [n_1 n_2 (1 \pm n_3)(1 \pm n_4) - n_3 n_4 (1 \pm n_1)(1 \pm n_2)] |M_{12\to34}|^2 \times \delta_{\mathbf{p}_1+\mathbf{p}_2,\mathbf{p}_3+\mathbf{p}_4} \delta(\varepsilon_1 + \varepsilon_2 - \varepsilon_3 - \varepsilon_4). \qquad (3)$$

Here $\mathbf{p}_i$ and ε_i are the quasiparticle momentum and energy, respectively, $n_i(\mathbf{p}_i)$ the quasiparticle distribution function, and $\mathbf{F}$ the force on a quasiparticle. The right-hand side (r.h.s.) is the collision integral for scattering particles from initial states 1 and 2 to final states 3 and 4, respectively, with matrix element squared $|M_{12\to34}|^2$ summed over final states and averaged over initial states. The $(1 \pm n_i)$ factors correspond physically to the Pauli blocking of final states, in the case of fermions, and to (induced or) stimulated emission, in the case of bosons. ν_2 is the statistical factor, 16 for gluons and $12N_f$ for quarks and antiquarks. For scattering of quarks of different flavor

$$|M_{12\to34}^{(qq')}|^2 = \frac{4}{9} g^4 \frac{u^2+s^2}{t^2} \frac{1}{16\varepsilon_1\varepsilon_2\varepsilon_3\varepsilon_4} ; \qquad (4)$$

quark-gluon and gluon-gluon interactions are just 9/4 and $(9/4)^2$ times stronger, respectively, near forward scattering. In a medium this singularity is screened as given by the Dyson equation in which a gluon self-energy $\Pi_{L,T}$ is added to the propagator

$$t^{-1} \to \omega^2 - q^2 - \Pi_{L,T} \qquad (5)$$

(we refer to [8] for details on separating longitudinal and transverse parts of the interaction), where the longitudinal and transverse parts of the self-energy in QED and QCD are for $\omega, q \ll T$ given by

$$\Pi_L(\omega, q) = q_D^2 \left(1 - \frac{x}{2} \ln \frac{x+1}{x-1}\right) , \qquad (6)$$

$$\Pi_T(\omega, q) = q_D^2 \left[\frac{1}{2}x^2 + \frac{1}{4}x(1-x^2) \ln \frac{x+1}{x-1}\right] , \qquad (7)$$

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where $x = \omega/qv_p$ and $v_p = c$ for the relativistic plasmas considered here. The Debye screening wave number in thermal QCD is $q_D^2 = g^2(2N + N_f)T^2/6$ where $N = 3$ is the number of colors, N_f the number of quark flavors, T the plasma temperature, and μ_q the quark chemical potential. The many similarities in QCD and QED plasmas are described in [9]. In the static limit, $\Pi_L(\omega = 0, q) = q_D^2$, and the longitudinal interactions are Debye screened. In contrast, $\Pi_T(\omega = 0, q) = 0$ and so the magnetic interactions are unscreened in the static limit. It has therefore been suggested that the transverse interactions are cut off below the "magnetic mass," $m_{\rm mag} \sim g^2T$, where infrared divergences appear in the plasma [10]. However, as was shown in [2, 3], dynamical screening due to Landau damping effectively screens the transverse interactions off in most transport problems at a length scale of order the Debye screening length $\sim 1/gT$ as in Debye screening. Nevertheless, there are three important length scales in the quark-gluon plasma. For a hot plasma they are, in increasing size, the interparticle spacing $\sim 1/T$, the Debye screening length $\sim 1/gT$, and the scale $1/m_{\rm mag} \sim 1/g^2T$ where QCD effects come into play.

Let us first consider a quark-gluon plasma where the particle flavors have been separated spatially; i.e., the flavor chemical potential depends on position, $\mu_i(\mathbf{r})$. In a steady state scenario the quark flavors will then be flowing with flow velocity, u_i. For simplicity we take the standard ansatz for the distribution functions (see, e.g., [7, 11]):

$$n_i(\mathbf{p}) = \left[\exp\left(\frac{\varepsilon_\mathbf{p} - \mu_i(\mathbf{r}) - \mathbf{u}_i\cdot\mathbf{p}}{T}\right) \pm 1\right]^{-1} \simeq n_i^0 - \frac{\partial n_p^0}{\partial \varepsilon_p}\mathbf{u}_i\cdot\mathbf{p}\,. \quad (8)$$

The expansion is valid near equilibrium where μ_i and therefore also $\mathbf{u}_i$ are small. It gives two terms, the equilibrium distribution function $n_i^0 = (\exp\{[\varepsilon_\mathbf{p} - \mu_i(\mathbf{r})]/T\} \pm 1)^{-1}$ and the deviation from that. In general the deviation from equilibrium has to be found self-consistently by solving the Boltzmann equation. However, as in the case of the viscosity [2], we expect the ansatz (8) to be good within a few percent to leading logarithmic order.

The flavor diffusion coefficient, $D_{\rm flavor}$, defined by

$$\mathbf{j}_i = -D_{\rm flavor}\nabla\rho_i\,, \quad (9)$$

is given in terms of the flavor current $\mathbf{j}_i$ and the gradient of the number density $\rho_i = \sum_p n_i^0(\mathbf{p}) = \nu_i T^3 3\xi(3)/4\pi^2$ of a particular flavor i. From (8) we find

$$\mathbf{j}_i = \sum_\mathbf{p} \mathbf{v}_\mathbf{p} n_i(\mathbf{p}) = \mathbf{u}_i\rho_i\,. \quad (10)$$

The density gradient, $\nabla\rho_i = -\nabla\mu_i \sum_p(\partial n_i^0/\partial\varepsilon_p)$, can be found by solving the Boltzmann equation. Linearizing in $\mathbf{u}_i$ we obtain [7]

$$\frac{\partial n_1}{\partial \varepsilon_1}\mathbf{v}_1\cdot\nabla\mu_1 = 2\pi\nu_2 \sum_{\mathbf{p}_2\mathbf{p}_3\mathbf{p}_4} |M_{12\to34}|^2 \times n_1^0 n_2^0(1 - n_3^0)(1 \pm n_4^0) \times (\mathbf{u}_1\cdot\mathbf{p}_1 + \mathbf{u}_2\cdot\mathbf{p}_2 - \mathbf{u}_3\cdot\mathbf{p}_3 - \mathbf{u}_4\cdot\mathbf{p}_4) \times \delta_{\mathbf{p}_1+\mathbf{p}_2,\mathbf{p}_3+\mathbf{p}_4}\delta(\varepsilon_1 + \varepsilon_2 - \varepsilon_3 - \varepsilon_4)\,. \quad (11)$$

It is most convenient to choose the plasma center-of-mass system (c.m.s.) where one flavor is flowing with velocity $\mathbf{u}_1$ and the others with velocity $\mathbf{u}_2 = -\mathbf{u}_1/(N_f - 1)$. The number of scatterers is then $\nu_2 = 12(N_f - 1)$. Equivalently, one can conveniently include the first flavor so that the number of scatterers is $\nu_2 = 12N_f$ but $\mathbf{u}_2 = 0$. In steady state the gluons will not move in the c.m.s.; i.e., $\mathbf{u}_2 = 0$ for quark-gluon scattering. Since the flavor is unchanged in the collisions $\mathbf{u}_3 = \mathbf{u}_1$ and $\mathbf{u}_4 = \mathbf{u}_2$.

To leading logarithmic order the singular interaction near forward scattering allows us to expand around $\mathbf{q} \sim 0$, where $\mathbf{q} = \mathbf{p}_1 - \mathbf{p}_3 = \mathbf{p}_4 - \mathbf{p}_2$ is the momentum transfer in the collision. Multiplying both sides of (11) by $\mathbf{p}_1$ and summing the Boltzmann equation reduces to

$$\rho_1\nabla\mu_1 = -\mathbf{u}_1\frac{\pi}{3}\nu_2 \sum_{\mathbf{q},\mathbf{p}_1,\mathbf{p}_2} n_1^0 n_2^0(1 - n_3^0)(1 \pm n_4^0) \times |M_{12\to34}|^2 q^2 \delta(\varepsilon_1 + \varepsilon_2 - \varepsilon_3 - \varepsilon_4)\,, \quad (12)$$

where we have used the antisymmetry of the r.h.s. by coordinate change $\mathbf{p}_1 \to \mathbf{p}_3$ so that $\mathbf{p}_1 \to \mathbf{q}/2$. The r.h.s. collision integral of Eq. (12) is now straightforward to evaluate to leading logarithmic order when the screening is properly included (see also Refs. [2, 3, 7]). We find

$$D_{\rm flavor}^{-1} \simeq \frac{\pi^5}{3^3 4\xi(3)^2}(1 + N_f/6)\alpha_s^2\ln(1/\alpha_s)T\,, \quad (13)$$

where ξ is the Rieman zeta function. The term 1 arises from quark-gluon scatterings and the $N_f/6$ from quark-quark scatterings. This result is similar to the viscous, thermal, and momentum relaxation rates because the collision term contains the same factors of momentum transfer: the singular q^{-4} factor from the matrix element squared and the suppressing q^2 factor because the quark flavors lose little momentum in forward scatterings. Including screening, $q^{-4} \to (q^2 + \Pi_{L,T})^{-2}$, where effectively $\Pi_{L,T} \sim q_D^2$, and integrating over momentum transfer, d^2q, gives the leading logarithmic term $\ln(T^2/q_D^2) \simeq \ln(1/\alpha_s)$.

Subsequently, let us consider the case where the particle spins have been polarized spatially by some magnetic field [11]; i.e., the spin chemical potential depends on position, $\mu_\sigma(\mathbf{r})$. With the analogous ansatz to (8) for the distribution functions with μ_σ instead of μ_i, we find the spin current $\mathbf{j}_1 = \mathbf{u}_1\rho_1$ of particle 1, where $\mathbf{u}_1$ and ρ_1 are the corresponding flow velocity and spin density. Linearizing the Boltzmann equation we find

$$\frac{\partial n_1}{\partial \varepsilon_1}\nabla\mu_{1,\sigma}\cdot\mathbf{v}_1 = 2\pi\nu_2 \sum_{\mathbf{p}_2\mathbf{p}_3\mathbf{p}_4} n_1^0 n_2^0(1-n_3^0)(1\pm n_4^0) \times \delta_{\mathbf{p}_1+\mathbf{p}_2,\mathbf{p}_3+\mathbf{p}_4}\delta(\varepsilon_1+\varepsilon_2-\varepsilon_3-\varepsilon_4) \times [|M^{\uparrow\downarrow}_{12\to34}|^2(\mathbf{u}_1-\mathbf{u}_2)\cdot(\mathbf{p}_1-\mathbf{p}_2) + |M^{\uparrow\uparrow}_{12\to34}|^2(\mathbf{u}_1-\mathbf{u}_2)\cdot\mathbf{q}]\,, \tag{14}$$

where $M^{\uparrow\downarrow}$ and $M^{\uparrow\uparrow}$ are the amplitudes for interacting with and without spin flip, respectively. Without spin flip the usual factor $\mathbf{q}$ as in flavor diffusion appears. With spin flip, however, $\mathbf{u}_3 = \mathbf{u}_2$ and $\mathbf{u}_4 = \mathbf{u}_1$ and the factor $(\mathbf{p}_1 - \mathbf{p}_2)$ appears. Because of Galilei invariance both terms are necessarily proportional to the relative flow, $(\mathbf{u}_1 - \mathbf{u}_2)$.

The transition current can be decomposed into interactions via the charge and the magnetic moment by the Gordon decomposition rule,

$$J_\mu = \frac{g}{2m}\bar{u}_f\left[(p_f+p_i)_\mu + i\sigma_{\mu\nu}(p_f-p_i)^\nu\right]u_i\,, \tag{15}$$

where only the latter can lead to spin flip. We notice that the spin-flip amplitude is suppressed by a factor $p_f - p_i$. Consequently, the spin-flip amplitude is suppressed by a factor q^2 and the matrix element squared by a factor q^4. We then find that the spin-flip interactions do not contribute to collisions to leading logarithmic order and the collision integral is similar to that for flavor diffusion evaluated above. The corresponding quark spin diffuseness parameter, defined by $\mathbf{j}_1 = -D_\sigma^{(q)}\nabla\rho_1$, is thus

$$D_\sigma^{(q)} = D_{\text{flavor}}\,. \tag{16}$$

Gluon spin diffusion is slower by a factor 4/9, due to the stronger interactions, and by another factor 4/9, due to differences between Bose and Fermi distribution functions, i.e., $D_\sigma^{(g)} \simeq (4/9)^2 D_\sigma^{(q)}$.

Finally, let us, like for the spin diffusion, assume that color has been polarized spatially given by a color chemical potential, $\mu_c(\mathbf{r})$. The basic difference between flavor and spin diffusion is that *quarks and gluons can easily change color directions* in forward scattering by color exchanges; i.e., one does not pay the extra q^2 penalty factor in the amplitude as in the case of spin flip. Consequently, the color-exchange interactions will dominate the collisions since they effectively reverse the color currents. The Boltzmann equation thus gives us an analogous result to Eq. (14) replacing spin by color where the color-exchange amplitude now dominates. The flow velocity of the scatterers, $\mathbf{u}_i$, $i = 1\cdots4$, depends on the color combination of the scattering quarks, antiquarks, and gluons. However, in c.m.s. the scatterer has vanishing flow velocity, $\mathbf{u}_2 = 0$, on average. Likewise the final velocities will be zero on average. Multiplying both sides with $\mathbf{p}_1$ and summing the Boltzmann equation reduces to [cf. Eq.(14)]

$$\rho_1\nabla\mu_{1,c} = -\mathbf{u}_1\pi\nu_2\sum_{\mathbf{q},\mathbf{p}_1,\mathbf{p}_2} n_1^0 n_2^0(1-n_3^0)(1\pm n_4^0)|M_{12\to34}|^2(\mathbf{p}_1-\mathbf{p}_2)^2\delta(\varepsilon_1+\varepsilon_2-\varepsilon_3-\varepsilon_4), \tag{17}$$

where we have used the antisymmetry by interchange of $\mathbf{p}_1 \to \mathbf{p}_2$. $\mathbf{u}_1$ and ρ_1 are now the color flow velocity and color density of particle 1. The matrix element entering in (17) is now averaged over all color combinations.

The transverse interactions actually diverge for small momentum and energy transfers even when integrating over energy transfers; i.e., dynamical screening is insufficient for obtaining a nonzero color diffusion coefficient like for the quasiparticle decay rates in QCD and QED plasmas (see [6]). Concentrating therefore on the leading contribution from transverse interactions at small $x = \omega/q$, where $\Pi_T \simeq i(\pi/4)q_D^2 x$, we find to leading order (see Ref. [6] for details)

$$\rho_1\nabla\mu_{1,c} = -\mathbf{u}_1\nu_1\frac{11\pi^3}{3^3 5}\alpha_s^2 T^7 \times\int_\lambda^{\sim T} q\,dq\int_{-1}^{1}dx\frac{1}{q^4+(\pi/4)^2 q_D^4 x^2} = -\mathbf{u}_1\nu_1\frac{22\pi^3}{3^3 5}\alpha_s^2\frac{T^7}{q_D^2}\ln\left(\frac{q_D^2}{\lambda^2}\right) \tag{18}$$

for quark scatterings with gluons. Quark-quark scattering adds a factor $(1 + 7N_f/33)$ to Eq. (18). The lower limit is now given by the infrared cutoff, λ. The upper limit on momentum transfers, $\sim T$, actually comes from the distribution functions in Eq. (17) but it does not enter here because only $q \lesssim q_D$ contribute to (18) to leading order. From the color current $\mathbf{j}_1 = \mathbf{u}_1\rho_1 \equiv -D_{\text{color}}\nabla\rho_1$ we find the color diffuseness parameter

$$D_{\text{color}}^{-1} = \frac{22\pi^6}{3^6\times 5\zeta(3)^2}\frac{1+7N_f/33}{1+N_f/6}\,\alpha_s\ln(q_D^2/\lambda^2)T\,. \tag{19}$$

With $\lambda \approx m_{\text{mag}} \sim g^2 T$ the analogous result to Eq. (2) is obtained:

$$D_{\text{color}}^{-1} \simeq 4.9\frac{1+7N_f/33}{1+N_f/6}\,\alpha_s\ln(1/\alpha_s)T\,. \tag{20}$$

Comparing Eqs. (16) and (20) we see that $D_{\text{color}} \sim \alpha_s D_{\text{flavor}}$. The color-exchange mechanism amplifies the forward collisions so the color cannot diffuse through the quark-gluon plasma as easily as spin or flavor.

The factor $\ln(1/\alpha_s)$ in D_{color} has a completely different origin from the one in D_{flavor} or D_σ. In D_{color} the logarithm arises from an integral dq/q over momentum transfers from $q \sim \lambda \sim g^2 T$ to $q \sim q_D \sim gT$ as in the case of quark and gluon quasiparticle decay rates of Eq. (2). In D_{flavor} or D_σ and the transport rates of Eq. (1) a similar integral occurs, but with momentum transfers from $q \sim q_D \sim gT$ to $q \sim T$ because of the extra factor

$\sim q^2/T^2$. The infrared cutoff does not enter these transport rates and they are reduced by a factor $q_D^2/T^2 \sim \alpha_s$.

Other related transport coefficients are the electrical conductivity, $\sigma_{\rm el}$, in QED and the corresponding color conductivity, $\sigma_{\rm color}$, in QCD. Applying a color-electric field, $\mathbf{E}_c$, to the quark-gluon plasma generates a color current, $\mathbf{j}_c$. The color conductivity $\sigma_{\rm color} = -j_c/E_c$ can thus be found by solving the Boltzmann equation analogous to the color diffusion process. We find

$$\sigma_{\rm color} = \frac{2}{3} g^2 D_{\rm color} \sum_{i,\mathbf{p}} \nu_i \left(\frac{\partial n_i^0}{\partial \varepsilon_i} \right) . \tag{21}$$

Here $D_{\rm color}$ plays the role of the color relaxation time. Equation (21) is the standard result for a plasma except for the factor 2/3 which arises because only two-thirds of the colors contribute to the currents for a given color field. Inserting $D_{\rm color}$ from (20) we obtain

$$\sigma_{\rm color} \simeq \frac{8\pi}{3} N_f \alpha_s D_{\rm color} T^2 \simeq 1.7 N_f T / \ln(1/\alpha_s) , \tag{22}$$

from quark currents alone. Gluon currents are slower due to stronger interactions and will reduce the conductivity slightly.

These surprising results for QCD are supported by those found by Selikhov and Gyulassy [1] who have considered the diffusion of color in color space. They use the fluctuation-dissipation theorem to estimate the deviations from equilibrium and find the same terms with and without color exchange, which they denote the color and momentum diffusion terms, respectively, and they also find that the former dominates, being infrared divergent. Inserting the same infrared cutoff, they find a color diffusion coefficient in color space equal to Eq. (2), $d_c = 1/\tau_1^{(g)} = 3\alpha_s \ln(1/\alpha_s) T$. Note that d_c is proportional to the inverse of $D_{\rm color}$. With d_c^{-1} as a typical relaxation time the color conductivity is estimated in [1] in the relaxation time approximation, and their result differs from Eq. (22) by a numerical factor only.

The color-exchange mechanism is not restricted to QCD but has analogs in other non-Abelian gauge theories. In the very early Universe when $T \gg M_W \simeq 80$ GeV, the mass of the $W^\pm$ bosons can be neglected and the electroweak interactions have the same screening problems as QCD and QED. Since now the exchanged $W^\pm$ bosons carry charge (unlike the photon, but similar to the colored gluon), they can easily change the charge of, for example, an electron to a neutrino in forward scatterings. Thus the collision term will lack the usual factor q^2 as for the quasiparticle damping and color diffusion rate. Since SU(2)$\times$U(1) gauge fields have similar infrared problems as SU(3) at the scale of the magnetic mass, $\sim e^2 T$, we insert this cutoff. Thus we find a diffusion parameter for charged electroweak particles in the very early Universe $T \gg M_W$ of order

$$D_{\rm el} \sim [\alpha \ln(1/\alpha) T]^{-1}, \tag{23}$$

which is a factor α smaller than when $T \ll M_W$. Similarly the electrical conductivity is smaller, $\sigma_{\rm el} \sim T/\ln(1/\alpha)$, when $T \gg M_W$ as compared to $\sigma_{\rm el} \sim T/\alpha \ln(1/\alpha)$ when $T \ll M_W$ [3].

In summary, the flavor, spin, and color diffusion coefficients have been calculated in QCD plasmas to leading order in the interaction strength. Color diffusion and the quark and gluon quasiparticle decay rates are not sufficiently screened and do depend on an infrared cutoff of order the magnetic mass, $m_{\rm mag} \sim g^2 T$; typically $D_{\rm color}^{-1} \sim \alpha_s \ln(q_D/m_{\rm mag}) T \sim \alpha_s \ln(1/\alpha_s) T$. Flavor and spin diffusion processes are sufficiently screened by Debye screening for the longitudinal or electric part of the interactions and by Landau damping for the transverse or magnetic part of the interactions; typically $D_{\rm flavor}^{-1} = D_\sigma^{-1} \sim \alpha_s^2 \ln(1/\alpha_s) T$. As a consequence, color diffusion is slow and the QGP is a poor color conductor. In the very early Universe when $T \gg M_W$ exchanges of $W^\pm$ provides charge exchange—a mechanism analogous to color exchange in QCD—and QED plasmas will also be poor electrical conductors.

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