Gluon decay as a mechanism for strangeness production in a quark-gluon plasma

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A calculation of thermal gluon decay shows that this process contributes significantly to strangeness production in a quark-gluon plasma. Our analysis does not support recent claims that this is the dominant process. In our calculations we take into account the resummed form of the transverse and longitudinal parts of the gluon propagator following the Braaten-Pisarski method. The estimate of the damping rate entering the effective gluon propagator subjects our results to uncertainty.

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I. INTRODUCTION

A possible signal for quark-gluon plasma formation is the enhancement of the production of strange particles. The original proposal by Rafelski and Müller [1] was followed by extensive discussion in the literature [2–7]. In this context thermal gluon decay has been recently discussed [5–7]. It has been claimed that the process $g \to \bar{q}q$ dominates for a wide range of quark masses [6,7]. Normally, the gluon cannot decay into a strange quark-antiquark pair because its thermal mass is too low. Even for the optimistic case where one takes the coupling constant $g = 2$ in a plasma with two massless quarks, the gluon mass is given to lowest order in perturbation theory by

$$m_g = \frac{2}{3} gT. \tag{1.1}$$

For a temperature of T= 200 MeV this gives $m_g = 267$ MeV which is below the threshold for the production of strange quarks. The important observation by Altherr and Seibert is that in addition to acquiring a thermal mass of the order g^2T [8] gluons also acquire a width determined by the large damping rate. This is the reason why thermal gluon decay into a heavy quark-antiquark pair is allowed, even though the gluon mass is below the threshold for strange pair production.

In this paper we present a systematic reevaluation of the production rate of massive quarks in a quark-gluon plasma due to the processes of quark-antiquark annihilation, gluon fusion, and thermal gluon decay in the spirit of Altherr and Seibert. Since the production rate depends strongly on the damping rate we take a more conservative approach in its estimation. Our main point is that we cannot support the claim that gluon decay is the dominant mechanism for strange quark production in a quark-gluon plasma. To the best of our knowledge, the gluon fusion mechanism remains the leading process.

This paper is organized as follows. In Sec. II we briefly review the properties of thermal-gluon propagators and the damping rate. In Sec. III we calculate the production rates. In Sec. IV we present results and concluding remarks.

II. GLUON PROPAGATOR AND THE DAMPING RATE

The effective gluon propagator at finite temperature in the Feynman gauge is given by [9]

$$iD^{ab}_{\mu\nu}(q_0, q) = -i\delta^{ab}[P^T_{\mu\nu}\Delta_T(q_0, q) + P^L_{\mu\nu}\Delta_L(q_0, q)], \tag{2.1}$$

where $P^T_{\mu\nu}$ and $P^L_{\mu\nu}$ are transverse and longitudinal projectors, respectively, and

$$\Delta_{T,L}(q_0, q) = \frac{1}{Q^2 - \Pi_{T,L}(q_0, q)}, \tag{2.2}$$

where $Q^2 \equiv q_0^2 - q^2$. The real transverse and longitudinal parts of the gluon self-energy in the high temperature limit are given, respectively, by

$$\mathrm{Re}\Pi_T(q_0, q) = \frac{3}{2} m_g^2 \left[\frac{q_0^2}{q^2} + \left(1 - \frac{q_0^2}{q^2}\right)\frac{q_0}{2q}\ln\frac{q_0+q}{q_0-q}\right] \tag{2.3}$$

and

$$\mathrm{Re}\Pi_L(q_0, q) = \frac{3}{2} m_g^2 \left(1 - \frac{q_0^2}{q^2}\right)\left[2 - \frac{q_0}{q}\ln\frac{q_0+q}{q_0-q}\right]. \tag{2.4}$$

The positions of the poles in the propagator (2.2) are determined by the dispersion relations

$$q_0^2 = q^2 + \Pi_{T,L}(q_0, q). \tag{2.5}$$

If a pole is located at

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$$q_0 = \omega_{T,L} + i\gamma_{T,L}, \tag{2.6}$$

then the imaginary shift of the pole $\gamma_{T,L}$ is related to the imaginary part of the self-energy through

$$\gamma_{T,L} = \text{Res}(\Delta_{T,L})\text{Im}\Pi_{T,L}, \tag{2.7}$$

where Res(Δ) is the residue of the propagator given by

$$\text{Res}(\Delta_{T,L})^{-1} = \left.\frac{\partial \Delta_{T,L}^{-1}}{\partial q_0}\right|_{\omega_{T,L}}, \tag{2.8}$$

or, explicitly in terms of $\omega_{T,L}$

$$\text{Res}(\Delta_T)^{-1} = -\omega_T + \frac{q^2}{\omega_T} + \frac{3m_g^2\omega_T}{\omega_T^2 - q^2}, \tag{2.9}$$

$$\text{Res}(\Delta_L)^{-1} = -\omega_L + \frac{q^2}{\omega_L} + \frac{3m_g^2}{\omega_L}. \tag{2.10}$$

Thus we can write (2.2) as

$$\Delta(Q) = \frac{Q^2 - \text{Re}\Pi}{(Q^2 - \text{Re}\Pi)^2 + \text{Res}(\Delta)^{-2}\gamma^2} + \frac{i\text{Res}(\Delta)^{-1}\gamma}{(Q^2 - \text{Re}\Pi)^2 + \text{Res}(\Delta)^{-2}\gamma^2}, \tag{2.11}$$

where we suppressed subscripts T, L; i.e., Δ, Π, and γ are either transverse or longitudinal. This expression will be used to replace the mass-shell δ function for thermal gluons.

The imaginary part of the pole in (2.6) gives the damping rate of the plasma oscillations [8,10]. In the following we consider transverse gluons, and we set $\gamma \equiv \gamma_T$. The damping rate is related to the so-called gluon magnetic mass m_{mag}, or the inverse magnetic screening length at high temperature. Unfortunately the exact relation between γ and m_{mag} is not known. A closed expression for the damping rate, derived by Pisarski, in the limit where $m_{\text{mag}} \gg \gamma$ [8], is

$$\gamma = \frac{g^2NT}{8\pi}\left[\ln\left(\frac{m_g^2}{m_{\text{mag}}^2 + 2m_{\text{mag}}\gamma}\right) + 1.096\,81...\right], \tag{2.12}$$

where the thermal gluon mass is given by

$$m_g^2 = \left(N_c + \frac{N_f}{2}\right)\frac{g^2T^2}{9}. \tag{2.13}$$

The magnetic mass at high temperature is of the form

$$m_{\text{mag}} = c_N g^2 T, \tag{2.14}$$

where c_N is a number depending on the gauge group and cannot be calculated by a perturbation expansion. Lattice estimates [11,12] for SU(2) giving

$$c_2 = 0.27 \pm 0.03 \tag{2.15}$$

have been confirmed by recent semiclassical calculations [13]. So far no reliable estimate exists for SU(3) [14]. The best one can do is extrapolate the SU(2) value using [15]

$$c_3 = \frac{3}{2}c_2. \tag{2.16}$$

Expanding the log in (2.12) in powers of γ/m_{mag} and retaining only the leading terms one finds

$$\gamma = (1+\eta)^{-1}\frac{g^2NT}{8\pi}\left[\ln\left(\frac{m_g^2}{m_{\text{mag}}^2}\right) + 1.096\,81...\right]. \tag{2.17}$$

$\eta = 0$ if we keep the leading log term only. If the next-to-leading term is included then

$$\eta = \frac{N}{4\pi c_N}. \tag{2.18}$$

To check the consistency of Pisarski's approximation we plot in Fig. 1 the damping rate for both values of η as well as the damping rate used in [6,7] and compare these with m_{mag}. We note the poor validity of the approximation for small values of g. The approximation is well justified if one uses the expression (2.17) in the range $1 < g < 2.5$.

In what follows when comparing our results with other papers we shall use the value $g = 2$. Since we approximate the thermal gluon width in such a way that it vanishes just below $g = 2.5$ (which is probably unrealistic), the ratio of our estimated width to that of Altherr and Seibert is made artificially small for $g > 1$. For that reason we shall use in our calculations all three parametrizations of the gluon width depicted in Fig. 1.

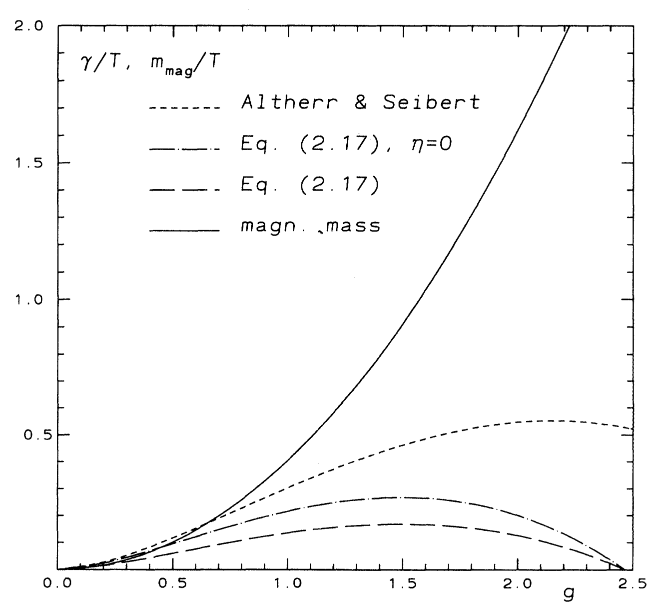

FIG. 1. The damping rate and the magnetic mass versus the coupling constant.

III. PRODUCTION RATES

Consider a quark-gluon plasma in which the gluons and the light quarks (u,d) are in thermal and chemical equilibrium. The strange quarks are also in thermal equilibrium, but away from chemical equilibrium, having very large and negative chemical potentials $\mu \equiv \mu_s = \mu_{\bar{s}}$. The chemical reactions

$$q + \bar{q} \leftrightarrow s + \bar{s}, \tag{3.1}$$

$$g + g \leftrightarrow s + \bar{s}, \tag{3.2}$$

$$g \leftrightarrow s + \bar{s} \tag{3.3}$$

will then take place until chemical equilibrium is reached. The total production rate due to (3.1)–(3.3), including the reverse processes, is given by [4]

$$\delta R = (1 - e^{2\beta\mu})(R_{q\bar{q}\to s\bar{s}} + R_{gg\to s\bar{s}} + R_{g\to s\bar{s}}), \tag{3.4}$$

where

$$R_{q\bar{q}\to s\bar{s}} = \int \frac{d^3p_q}{(2\pi)^3 2E_q}\frac{d^3p_{\bar{q}}}{(2\pi)^3 2E_{\bar{q}}}\frac{d^3p_s}{(2\pi)^3 2E_s}\frac{d^3p_{\bar{s}}}{(2\pi)^3 2E_{\bar{s}}}(2\pi)^4\delta(P_q + P_{\bar{q}} - P_s - P_{\bar{s}})$$
$$\times f_{\rm FD}(E_q)f_{\rm FD}(E_{\bar{q}})[1 - f_{\rm FD}(E_s)][1 - f_{\rm FD}(E_{\bar{s}})]\sum|\mathcal{M}(q\bar{q}\to s\bar{s})|^2, \tag{3.5}$$

$$R_{gg\to s\bar{s}} = \tfrac{1}{2}\int \frac{d^3p_1}{(2\pi)^3 2E_1}\frac{d^3p_2}{(2\pi)^3 2E_2}\frac{d^3p_s}{(2\pi)^3 2E_s}\frac{d^3p_{\bar{s}}}{(2\pi)^3 2E_{\bar{s}}}(2\pi)^4\delta(P_1 + P_2 - P_s - P_{\bar{s}})$$
$$\times f_{\rm BE}(E_1)f_{\rm BE}(E_2)[1 - f_{\rm FD}(E_s)][1 - f_{\rm FD}(E_{\bar{s}})]\sum|\mathcal{M}(gg\to s\bar{s})|^2, \tag{3.6}$$

and

$$R_{g\to s\bar{s}} = \int \frac{d^3q}{(2\pi)^3 2E_g}\frac{d^3p_s}{(2\pi)^3 2E_s}\frac{d^3p_{\bar{s}}}{(2\pi)^3 2E_{\bar{s}}}(2\pi)^4\delta(Q - P_s - P_{\bar{s}})$$
$$\times f_{\rm BE}(E_g)[1 - f_{\rm FD}(E_s)][1 - f_{\rm FD}(E_{\bar{s}})]\sum|\mathcal{M}(g\to s\bar{s})|^2. \tag{3.7}$$

In our notation the four-momenta are denoted by the capitals P, Q, etc. The summation in the above equations extends over all colors and polarizations of the gluons and the final state quark-antiquark pair. We have included the Pauli blocking factors although, as long as the density of strange quarks is well below one, Pauli blocking does not play a significant role. Following Matsui, McLerran, and Svetitsky [4] we can investigate the evolution process in terms of the relaxation time determined near equilibrium. Therefore we have to evaluate the rates (3.5)–(3.7) at $\mu = 0$, i.e., when quarks are in both thermal and chemical equilibrium. Thermal field theory calculations become fully legitimate in this way. In particular, we can use the thermal quark mass given by [16]

$$m_s^2(T) = m_s^2(0) + \frac{g^2T^2}{6} \tag{3.8}$$

and the thermal gluon mass (2.13).

The processes of gluon fusion and quark-antiquark annihilation have been discussed in Ref. [4] and we shall use their expressions for (3.5) and (3.6). The thermal gluon decay, also discussed by Altherr and Seibert [6,7], can be calculated similarly. We first replace the integrations over q, p_s, and $p_{\bar{s}}$ by

$$\frac{d^3q}{2E_g} = d^4Q\,\delta(Q^2 - m_g^2)\theta(q_0), \tag{3.9}$$

$$\frac{d^3p_s}{2E_s}\frac{d^3p_{\bar{s}}}{2E_{\bar{s}}} = d^4P_s\,\delta(P_s^2 - m_s^2)\theta(p_s^0)\,d^4P_{\bar{s}}\,\delta(P_{\bar{s}}^2 - m_s^2)\theta(p_{\bar{s}}^0), \tag{3.10}$$

and change variables

$$Q' = P_s + P_{\bar{s}},$$
$$P = \frac{1}{2}(P_s - P_{\bar{s}}). \tag{3.11}$$

After trivially eliminating integrals over d^4Q' and d^3p we find

$$R_{g\to s\bar{s}} = \frac{1}{4(2\pi)^4}\int d^4Q\delta(Q^2 - m_g^2)\frac{1}{q}f_{\rm BE}(q_0)$$
$$\times\int dp_0[1 - f_{\rm FD}(\tfrac{1}{2}q_0 + p_0)][1 - f_{\rm FD}(\tfrac{1}{2}q_0 - p_0)]$$
$$\times\sum|\mathcal{M}(g\to s\bar{s})|^2, \tag{3.12}$$

where the integration space is restricted by the following kinematical constraints:

$$q_0 > 2m_s, \quad 0 < q < (q_0^2 - 4m_s^2)^{1/2},$$
$$p_0^2 < \frac{q^2}{4}\left(1 - \frac{4m_s^2}{Q^2}\right). \tag{3.13}$$

It immediately follows that $R_{g\to s\bar{s}} = 0$ if $m_g < 2m_s$. At the relevant temperatures the thermal gluon mass is not high enough to allow for decay into a strange quark pair. It is only because of its width that the gluon can decay. To take this into account the δ function is replaced by a function similar to a Breit-Wigner resonance. In the case of a narrow resonance the width of the resonance is related to the imaginary shift of the pole in the propagator in the complex q_0 plane

$$\frac{1}{q_0^2-(\sqrt{q^2+m^2}+i\gamma)^2}\approx\frac{1}{Q^2-m^2}+\frac{i2\sqrt{q^2+m^2}\,\gamma}{(Q^2-m^2)^2+4(q^2+m^2)\gamma^2}\tag{3.14}$$

which in the limit $\gamma\to 0$ yields the standard free particle propagator

$$\frac{1}{Q^2-m^2-i\epsilon}=\mathcal{P}\frac{1}{Q^2-m^2}+i\pi\delta(Q^2-m^2).\tag{3.15}$$

Thus for a Breit-Wigner resonance with width $\Gamma=\gamma/2$ the mass-shell δ function should be replaced by

$$\delta(Q^2-m^2)\to\frac{1}{\pi}\frac{\sqrt{q^2+m^2}\,\Gamma}{(Q^2-m^2)^2+(q^2+m^2)\Gamma^2}\ .\tag{3.16}$$

This simple prescription cannot be directly applied to the case of thermal gluons because the location of the pole is determined by complicated dispersion relations (2.5) for transverse (T) and longitudinal (L) gluons. Instead of (3.16) we use (2.11) to obtain

$$\delta(Q^2-m_g^2)\to\frac{1}{\pi}\frac{\text{Res}(\Delta_{T,L})^{-1}\gamma_{T,L}}{(Q^2-\text{Re}\Pi_{T,L})^2+\text{Res}(\Delta_{T,L})^{-2}\gamma_{T,L}^2}.\tag{3.17}$$

The transverse and longitudinal Breit-Wigner-like distributions are shown in Figs. 2(a) and 2(b), respectively, for three different values of the momentum q. One can clearly see the peak corresponding to the thermal gluon mass. This peak levels off as one increases the value of q. Also shown is a vertical line corresponding to the strange quark production threshold $4m_s^2$, where we have taken m_s to be 0.2 GeV [17]. Note that this threshold is in the tail of the distribution. Only the part to the right of this line contributes to strange quark production.

The matrix element is simply given by

$$\mathcal{M}(g\to s\bar{s})=g\epsilon_\mu(\zeta)\bar{u}(P_s)\gamma^\mu\lambda_a v(P_{\bar{s}}),\tag{3.18}$$

where $\epsilon_\mu(\zeta)$ is the polarization vector of the decaying gluon and the λ_a are SU(3) matrices.

Summing over colors and polarizations of the gluon leads to

$$\sum_{a,\zeta}|\mathcal{M}(g\to s\bar{s})|^2=-4g^2\text{Tr}[(\not{P}_s+m_s)\gamma_\mu(\not{P}_{\bar{s}}-m_s)\gamma^\mu]$$
$$=16g^2(2m_s^2+Q^2).\tag{3.19}$$

Since the frame of the quark-gluon plasma introduces a preferred direction, it is furthermore necessary to distinguish between the transverse and the longitudinal components of the gluons. If the sum is taken over transverse or longitudinal polarization only, we find

$$\sum_T|\mathcal{M}(g\to s\bar{s})|^2=8g^2\left[4m_s^2+Q^2\left(1+4\frac{p_0^2}{q^2}\right)\right],\tag{3.20}$$

$$\sum_L|\mathcal{M}(g\to s\bar{s})|^2=8g^2Q^2\left(1-4\frac{p_0^2}{q^2}\right).\tag{3.21}$$

By making use of (3.17), (3.20), and (3.21) we find from (3.12)

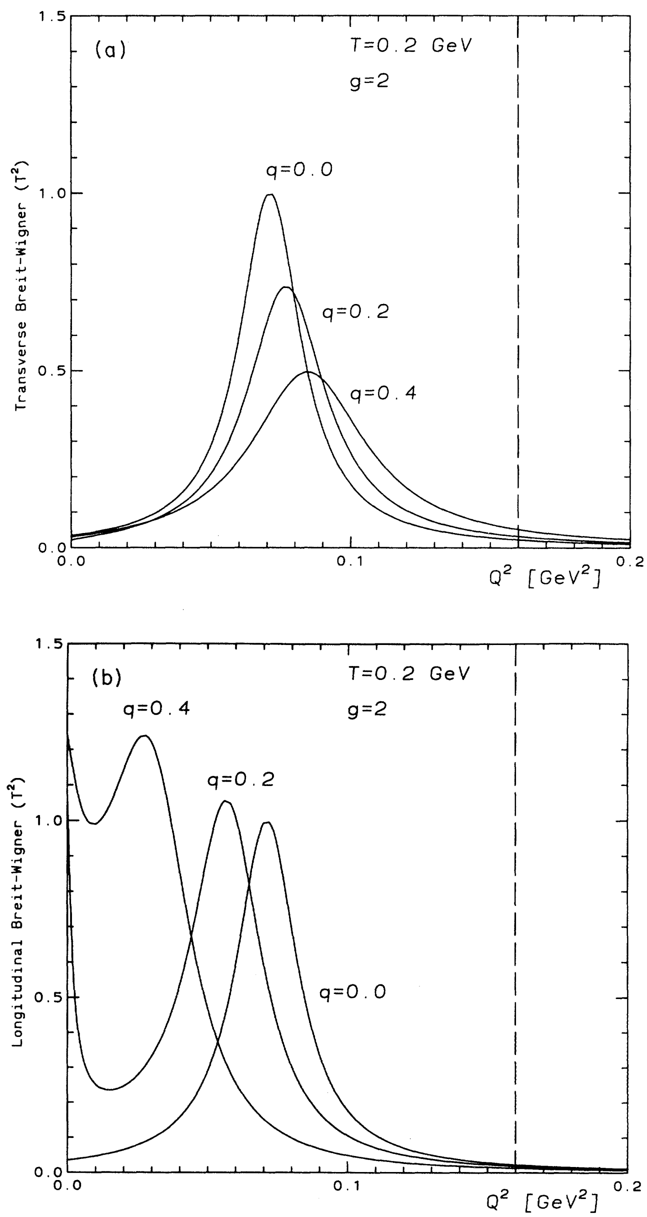

FIG. 2. The Breit-Wigner form given in Eq. (3.17) as a function of Q^2 for three different values of the gluon momentum q (units are in GeV). (a) Transverse part, (b) longitudinal part.

$$R^T_{g\to s\bar{s}} = \frac{g^2}{2\pi^4}\int_{2m_s}^{\infty} dq_0 f_{\rm BE}(q_0)\int_0^{\sqrt{q_0^2-4m_s^2}} dq\, q \int_{-\frac{q}{2}\sqrt{1-\frac{4m_s^2}{Q^2}}}^{\frac{q}{2}\sqrt{1-\frac{4m_s^2}{Q^2}}} dp_0[1-f_{\rm FD}(\tfrac{1}{2}q_0+p_0)][1-f_{\rm FD}(\tfrac{1}{2}q_0-p_0)]$$
$$\times\frac{{\rm Res}(\Delta_T)^{-1}\gamma_T}{(Q^2-{\rm Re}\Pi_T)^2+{\rm Res}(\Delta_T)^{-2}\gamma_T^2}\left[4m_s^2+Q^2\left(1+4\frac{p_0^2}{q^2}\right)\right], \quad (3.22)$$

and a similar expression for $R^L_{g\to s\bar{s}}$. The production rate due to gluon decay is given by the sum

$$R_{g\to s\bar{s}} = R^T_{g\to s\bar{s}} + R^L_{g\to s\bar{s}}. \quad (3.23)$$

If we neglect the Pauli blocking factors the integral over p_0 can be done explicitly, leading to

$$R^T_{g\to s\bar{s}} = \frac{2g^2}{3\pi^4}\int_{2m_s}^{\infty} dq_0 f_{\rm BE}(q_0)\int_0^{\sqrt{q_0^2-4m_s^2}} \times dq\, q^2\sqrt{1-\frac{4m_s^2}{Q^2}}(Q^2+2m_s^2) \times\frac{{\rm Res}(\Delta_T)^{-1}\gamma_T}{(Q^2-{\rm Re}\Pi_T)^2+{\rm Res}(\Delta_T)^{-2}\gamma_T^2} \quad (3.24)$$

and a similar expression for $R^L_{g\to s\bar{s}}$.

We use the full high temperature expressions for ${\rm Re}\Pi_{T,L}$ given by (2.14) and (2.15) and numerically solve the dispersion relations (2.5) in order to determine ${\rm Res}(\Delta_{T,L})$ from (2.9) and (2.10). The temperature-dependent gluon mass is given by (2.13), and the damping rate $\gamma \equiv \gamma_T \approx \gamma_L$ is estimated using (2.17).

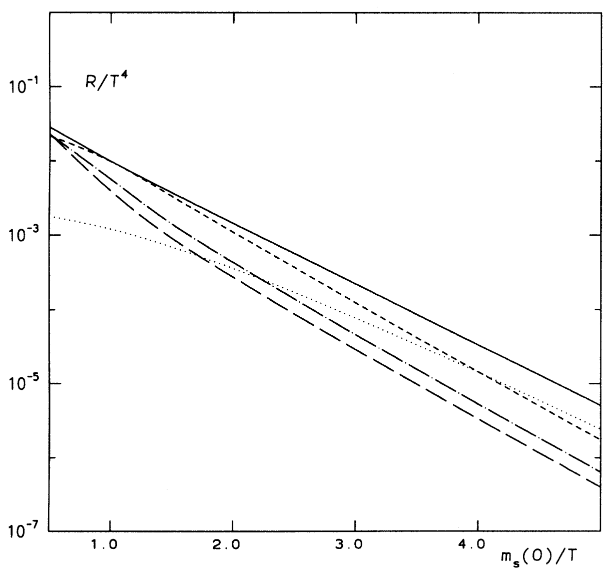

FIG. 3. The quark production rate for thermal gluon decay for different damping rates (short-dashed, long-dashed, and dot-dashed lines) corresponding to the damping rates depicted in Fig. 1 compared to the production rate for gluon fusion (solid line) and quark-antiquark annihilation (dotted line). The mass m_s is temperature independent.

IV. RESULTS AND CONCLUSION

The rates for different processes are depicted in Figs. 3 and 4. Our numerical calculation of the thermal gluon decay is done using Eqs. (3.22) and (3.23) with (2.3), (2.4), (2.9), (2.10), and (2.17). We calculate the rates for quark-antiquark annihilation and gluon fusion by making use of Eqs. (3.23)–(3.25) in Ref. [4]. We fix the QCD running coupling constant at the value $g = 2$ because the temperature during the time evolution is almost constant and the running coupling effect is negligible.

In comparison with Ref. [4] our calculation of gluon fusion and quark-antiquark annihilation differs in that we include effects of the thermal quark mass (Fig. 4) and in the choice of parameters: we fix $g = 2$ and $m_s = 0.2$ GeV versus $\alpha_S = 0.6$ and $m_s = 0.15$ MeV. Our choice reduces their rates by about a factor of 3.

In Fig. 3 the quark mass is kept fixed while in Fig. 4 it varies with temperature according to (3.8). In both cases we find that gluon fusion dominates everywhere. The gluon decay process is comparable to the gluon fusion in the narrow region around $m_s(0)/T = 1$ only if we choose the optimistic parametrization [7] of the damping rate. Had we used $g = 1$ instead of $g = 2$ the difference between the gluon decay rates with the different parametrizations of the width would not be so pronounced (about 30–50% for $g = 1$ compared to 70–80% for $g = 2$).

A more complete calculation should also include ther-

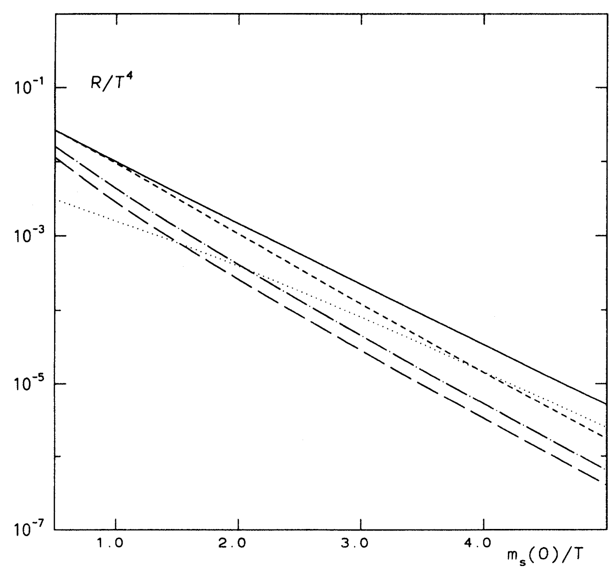

FIG. 4. Same as Fig. 3 with the thermal mass m_s given by Eq. (3.8).

mal gluons and resummed propagators in the gluon fusion process. For the present analysis this is not necessary because of the following arguments. Gluon fusion with massless gluons and *nonzero* m_s is of the order α_S^2. Including the thermal gluon mass would amount to corrections

$$R_{gg\to s\bar{s}} \sim \alpha_S^2 \left(1 + a\frac{m_g^2}{T^2} + b\frac{m_g^2}{m_s^2} + \cdots\right) \tag{4.1}$$

with $m_g^2 \sim \alpha_S$. If the thermal width is included there will be additional corrections

$$R_{gg\to s\bar{s}} \sim \alpha_S^2 \left(1 + \cdots + c\frac{\gamma^2}{T^2} + d\frac{\gamma^2}{m_s^2} + \cdots\right) \tag{4.2}$$

with $\gamma \sim \alpha_S \ln \alpha_S$. Constants a, b, c, d can be calculated using for example the amplitude square given in [5]. Thus, the corrections coming from the thermal mass are of the order α_S^3 and those from the width are $\alpha_S^4(\ln\alpha_S)^2$.

On the other hand, gluon decay is of the order $\alpha_S^2(\ln\alpha_S)$, which can be seen from (3.22), and therefore it is legitimate to compare gluon decay to gluon fusion calculated to lowest order. Moreover, including the above-mentioned corrections would be consistent only if one also includes other α_S^3 contributions.

The production of massless quarks is another matter. In that case the above corrections contribute to lowest order, and one needs the resummation program for the gluon fusion process. In addition the resummation is believed to cure a kinematic divergence of the type $\ln(Q^2/m_q^2)$ [18] as discussed by Altherr and Seibert [7]. However, this program goes beyond the aim of our paper.

It has been shown that the time dependence of the strange-quark density can, to a great degree of accuracy, be described by the approximate evolution equation [4]

$$n_s(t) = n_s^{eq}\tanh\left(\frac{t}{2\tau} + \text{const}\right), \tag{4.3}$$

where the relaxation time is defined as

$$\tau = \frac{1}{2\beta R}\left.\frac{\partial n_s}{\partial \mu}\right|_\epsilon \tag{4.4}$$

with

$$R = R_{q\bar{q}\to s\bar{s}} + R_{gg\to s\bar{s}} + R_{g\to s\bar{s}}. \tag{4.5}$$

The derivative of n_s with respect to μ at fixed energy density is given by

$$\left.\frac{\partial n_s}{\partial \mu}\right|_\epsilon = \frac{\partial n_s}{\partial \mu} - \frac{\partial n_s}{\partial T}\left(\frac{\partial \epsilon}{\partial T}\right)^{-1}\frac{\partial \epsilon}{\partial \mu}, \tag{4.6}$$

where

$$n_s = 2N_f \int \frac{d^3p}{(2\pi)^3} f_{\rm FD}(E_s, \mu) \tag{4.7}$$

and

$$\epsilon = \frac{(N^2-1)\pi^2T^4}{15} + 4N_fN\int\frac{d^3p}{(2\pi)^3}E_q f_{\rm FD}(E_q, 0) + 4N\int\frac{d^3p}{(2\pi)^3}E_s f_{\rm FD}(E_s, \mu). \tag{4.8}$$

All the quantities in (4.4) are to be evaluated at $\mu = 0$. In Fig. 5 we plot the relaxation time for the saturation of the strange-quark density for the massive quarks with the zero temperature mass $m_s(0) = 0.2$ GeV along with the classical approximation. In this approximation Pauli blocking factors $(1-f_{\rm FD})$ are eliminated and the remaining Fermi-Dirac and Bose-Einstein distributions are replaced by the Boltzmann distribution. For comparison we also plot the relaxation time for the massless quarks. Despite the inclusion of gluon decay, our relaxation time is still larger than Ref. [4] by about a factor of 2 due to the effect of a thermal quark mass and a more conservative choice of parameters.

We comment here on the various approximations made in the gluon decay calculation. First of all, use was made of the Braaten-Pisarski resummation scheme. This is strictly valid only when $gT \ll T$, which is clearly not the case here. However, this is the case with most applications of QCD at finite temperatures. Second, the magnetic mass has been introduced although only very limited knowledge is available. It has been used to calculate the damping rate of a thermal gluon inside a plasma. In comparison to the calculations of Refs. [6,7] we keep the standard form of the Breit-Wigner distribution. The consequence of this is that the rate for high masses is reduced, while for low masses it is enhanced. We also avoid a rather heuristic assumption that the thermal quark

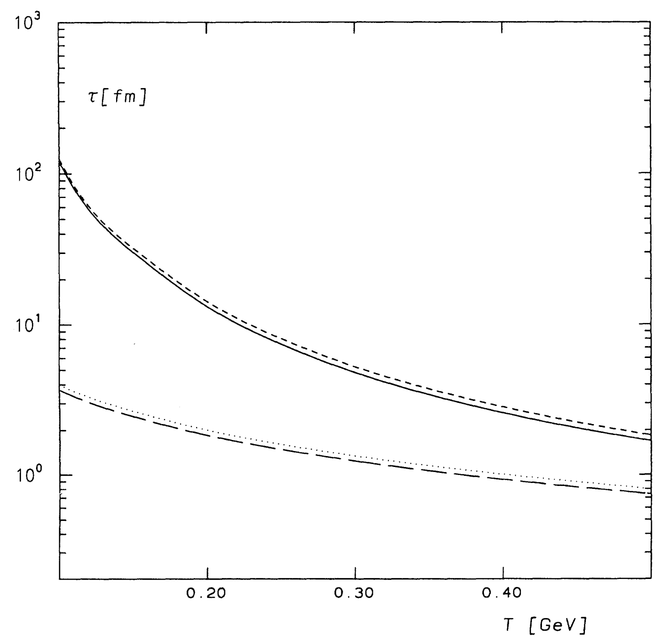

FIG. 5. Relaxation times for the density of massive (solid line) and massless (long-dashed line) quarks. Corresponding relaxation times in the classical approximation are plotted with dashed and dotted lines, respectively.

mass is generated by gluons only. Since our rates are defined near equilibrium our thermal mass includes both thermal gluon and thermal quark contribution. Most of the difference from the results of Altherr and Seibert is due to the fact that, for the temperature range of interest, the strange quark is not very massive, so the limit $m/T \to \infty$ is not accurate for calculating strange quark production rates. This can be seen in the fact that the rate from quark-antiquark annihilation is only 10% of that from gluon fusion, while in the heavy quark limit the two should be nearly equal. Comparison of the exact ratio of the fusion to annihilation rates with the approximate one indicates that the heavy quark limit is not very accurate for m/T less than about 25.

Our main point has been that even with the parameters chosen in Refs. [6,7] we do not support the claim that gluon decay prevails over gluon fusion processes as the dominant mechanism for strange quark production inside a quark-gluon plasma. To the best of our knowledge, the gluon fusion mechanism is the leading process.

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