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ENHANCEMENT OF RESONANT THERMONUCLEAR REACTION RATES IN EXTREMELY DENSE STELLAR PLASMAS

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ABSTRACT

The enhancement factor of resonant thermonuclear reaction rates is calculated for extremely dense stellar plasmas in the liquid phase. In order to calculate the enhancement factor we use the screening potential deduced from the numerical experiment of the classical one-component plasma. It is found that the enhancement is tremendous for white dwarf densities if the $^{12}C + ^{12}C$ fusion cross sections show resonant behavior in the astrophysical energy range. We summarize our numerical results by accurate analytic fitting formulae.

Subject headings: dense matter — nuclear reactions, nucleosynthesis, abundances — plasmas

On-line material: color figures

1. INTRODUCTION

In a recent important paper, Cussons, Langanke, & Liolios (2003) have pointed out potential resonant screening effects on stellar $^{12}C + ^{12}C$ reaction rates. The $^{12}C + ^{12}C$ fusion cross sections show noticeable resonant structures down to the lowest energies measured so far in the laboratory, $E \sim 2.4$ MeV (Kettner, Lorenz-Wirzba, & Rolfs 1980). If the resonant structure continues to even lower energies and the astrophysical reaction rate is due to the contributions of narrow resonances, one then has to consider that the entrance channel width of these resonances will be modified in the plasma.

Cussons et al. (2003) have specifically pointed out the possible importance of plasma effects on resonant $^{12}C + ^{12}C$ reactions for a carbon white dwarf environment with $T = 5 \times 10^7$ K and $\rho = 2 \times 10^9$ g cm^{-3}. They have considered a resonance energy interval 0.4–2 MeV. They have specifically discussed a rather extreme case of the low resonance energy $E_r = 400$ keV and have estimated the overall enhancement of the resonant $^{12}C + ^{12}C$ reaction rates due to plasma effects for this case.

Cussons et al. (2003) adopted the method of Salpeter & Van Horn (1969), which is based on the lattice model of dense plasma, to calculate the resonant screening effects. One of us (N. I.) and collaborators have calculated the enhancement of nonresonant thermonuclear reaction rates in extremely dense stellar plasmas (Itoh, Kuwashima, & Munakata 1990). This work is a natural extension of the works of Itoh, Totsuji, & Ichimaru (1977) and Itoh et al. (1979), and improves on the accuracy of the results of Salpeter & Van Horn (1969). Itoh et al. (1990) have summarized their numerical results in an accurate analytical fitting formula, which will be readily implemented in stellar evolution computations.

The aim of the current paper is to extend the work of Itoh et al. (1990) to the case of resonant reactions. This paper is organized as follows. Physical conditions relevant to the present calculation are made explicit in § 2. Calculation of the enhancement factor of resonant thermonuclear reaction rates is summarized in § 3. The results are presented in § 4. An extension to the case of ionic mixtures is made in § 5. Concluding remarks are given in § 6.

2. PHYSICAL CONDITIONS

First we consider thermonuclear reactions that take place in the plasma in thermodynamic equilibrium at temperature T, composed of one kind of atomic nuclei and electrons with number densities n_i and n_e, respectively; Ze and M denote the electric charge and the mass of such an ion. The conventional parameters that characterize such a plasma are

$$\Gamma = \frac{(Ze)^2}{ak_B T} = 0.2275 \frac{Z^2}{T_8} \left(\frac{\rho_6}{A}\right)^{1/3} , \qquad (1)$$

$$\tau = \left[\left(\frac{27\pi^2}{4}\right) \frac{M(Ze)^4}{\hbar^2 k_B T}\right]^{1/3} , \qquad (2)$$

where a is the ion sphere radius $a = (4\pi n_i/3)^{-1/3}$, A is the mass number of the nucleus, T_8 is the temperature in units of 10^8 K, and ρ_6 is the mass density in units of 10^6 g cm^{-3}. In this paper we restrict ourselves to the case in which electrons are strongly degenerate. This condition is expressed as

$$T \ll T_F = 5.930 \times 10^9$$
$$\times \left\{ \left[1 + 1.018(Z/A)^{2/3}\rho_6^{2/3}\right]^{1/2} - 1 \right\} \text{ K} , \quad (3)$$

where T_F is the electron Fermi temperature. Furthermore, we consider the case in which the ions can be treated approximately as classical particles. The corresponding condition is written as

$$n_i \Lambda^3 \le 1 , \qquad (4)$$

where $\Lambda = (2\pi\hbar^2/Mk_B T)^{1/2}$ is the thermal de Broglie wavelength of the ions. Equation (4) is rewritten as

$$T_8 \ge 2.173\rho_9^{2/3}/A^{5/3} , \qquad (5)$$

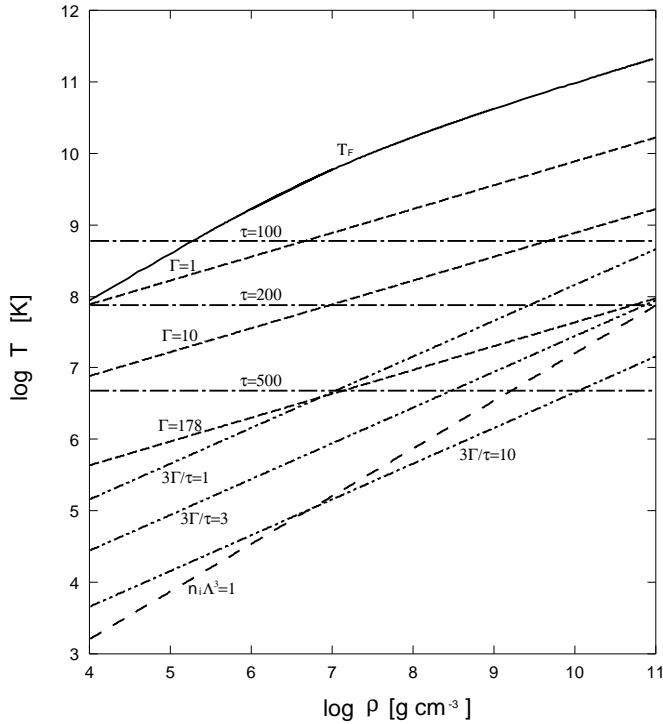where ρ_9 is the mass density in units of 10^9 g cm^{-3}. In this

FIG. 1.—Physical conditions for the pure ^{12}C plasma. The present calculation is valid in the region bounded by the electron Fermi temperature T_F and the line $n_i\Lambda^3 = 1$. Furthermore, the ions should be in the liquid state, $\Gamma < 178$. The lines of $\tau = $ const, $\Gamma = $ const, and $3\Gamma/\tau = $ const are shown. [*See the electronic edition of the Journal for a color version of this figure.*]

paper we impose a further condition that the ions are in the liquid state (Slattery, Doolen, & DeWitt 1982),

$$\Gamma < 178 . \tag{6}$$

The parameter $3\Gamma/\tau$ corresponds to the ratio of the classical turning point radius at the Gamow peak and the mean inter-ionic distance in the case of pure Coulomb potential (Itoh et al. 1977; Alastuey & Jancovici 1978). The theories of Itoh et al. (1977), Itoh et al. (1979), and also Alastuey & Jancovici (1978) are valid under the condition $3\Gamma/\tau \leq 1.6$. Itoh et al (1990) have extended the calculation of the enhancement of the nonresonant thermonuclear reaction rates to the case $3\Gamma/\tau \leq 5.4$. In Figure 1 we show the density-temperature diagram of the pure ^{12}C plasma, which illustrates the physical conditions described in this section.

3. ENHANCEMENT OF THE RESONANT THERMONUCLEAR REACTION RATES

In this paper we use the following screening potential for the classical one-component plasma:

$$H(r) = Ck_BT - \frac{Z^2e^2}{a}\left[\frac{1}{4}\left(\frac{r}{a}\right)^2 - b\left(\frac{r}{a}\right)^4\right],$$
$$0 \leq r \leq r_0 ; \tag{7}$$

$$H(r) = \frac{Z^2e^2}{a}\left[1.25 - \left(\frac{1.25}{2}\right)^2\frac{r}{a}\right],$$
$$r_0 \leq r \leq 1.60a ; \tag{8}$$

$$C = 1.0531\Gamma + 2.2931\Gamma^{1/4} - 0.5551\ln\Gamma - 2.35 . \tag{9}$$

This screening potential is derived from the equilibrium pair correlation function of the classical one-component plasma. In adopting this screening potential, our standpoint is the same as that of Alastuey & Jancovici (1978). They have argued that the pair correlation function should be taken as the static one. The point is that the transmission coefficient of the potential barrier is exceedingly small, which makes nuclear reactions very rare events. In a loose classical analogy, one might say that, in most collisions, the colliding particles tunnel through only a certain distance and are then reflected back. Therefore, as soon as r is larger than a few nuclear diameters, equilibrium is achieved and the probability of finding two nuclei at a distance r from one another is given by the equilibrium pair correlation function. Thus, one can use the averaged potential in describing the tunnelling. More recently, DeWitt (1994), Rosenfeld (1994), and Isern & Hernanz (1994) have studied this problem. They have essentially confirmed the correctness of the method of Alastuey & Jancovici (1978) on which our work is based. Similar work addressing the case of solar fusion reactions has been recently carried out by Bahcall et al. (2002), essentially confirming the soundness of the method of the average potential.

We also remark in relation to the above-stated point that the nuclear reactions as a whole are taking place on a macroscopic timescale, whereas the screening potential is kept in equilibrium on a microscopic timescale.

We here remark on the validity of the classical one-component plasma (OCP) model. In the interior of dense stars the electron Fermi energy becomes much larger than the Coulomb interaction energy between the electron and the ion. Therefore, the electron liquid becomes an almost uniform liquid. Owing to this fact, the interior of dense stars can be satisfactorily described by the OCP model, which consists of point ions embedded in the rigid background of electrons. Of course, at the same density the OCP model becomes better for the smaller ionic charge Z, as the ratio of the electron-ion interaction energy to the electron Fermi energy becomes smaller for smaller Z.

The expression for C is taken from Alastuey & Jancovici (1978). The two segments of the screening potential are matched at $r = r_0$, so that the screening potential and its derivative are a continuous function with respect to the distance r. This procedure produces solutions for r_0 and b for the range of Γ values $4 \leq \Gamma \leq 90$. Outside this range we use the value of b that makes equations (7) and (8) continuous at $r = 1.171875a$. In this case the first derivatives of equations (7) and (8) are slightly discontinuous at this point. The linearly decreasing part of the screening potential is identical to that employed by Itoh et al. (1977) and also by Itoh et al. (1979). The screening potential of equations (7) and (8) fits the results of the numerical experiments excellently. (See Fig. 1 of Itoh et al. 1979 for the accuracy of this screening potential in reproducing the results of the Monte Carlo computations.) Note that this screening potential exactly cancels the Coulomb potential Z^2e^2/r at $r = 1.60a$. We further assume that the potential of mean force vanishes for $r \geq 1.60a$. Given the explicit form of the screening potential, we are now in a position to calculate the enhancement of the resonant thermonuclear reaction rates.

A single resonance in the cross section of a nuclear reaction, $0 + 1 \longrightarrow 2 + 3$, can be represented most simply as a function of energy in terms of the classical Breit-Wigner

formula (Fowler, Caughlan, & Zimmerman 1967),

$$\sigma = \frac{\pi\hbar^2}{2\mu E} \frac{\omega_r \Gamma_1 \Gamma_2}{(E - E_r)^2 + \Gamma_{\rm tot}^2/4} \ , \tag{10}$$

where μ is the reduced mass, E is the center-of-mass energy, E_r is the resonance energy, ω_r is the statistical weight factor, Γ_1 is the partial width for the decay of the resonant state by re-emission of $0 + 1$, Γ_2 is the partial width for emission of $2 + 3$, and $\Gamma_{\rm tot} = \Gamma_1 + \Gamma_2 + \ldots$ is the sum over all partial widths. The partial width Γ_1 is proportional to the barrier penetration factor $P(E)$ for the screened Coulomb potential.

$$\Gamma_1 \propto P(E) = \exp\left\{ -\frac{2\sqrt{2\mu}}{\hbar} \int_0^{r_{\rm tp}} [V(r) - E]^{1/2} \, dr \right\} , \tag{11}$$

$$V(r) = \frac{Z^2 e^2}{r} - H(r) \ , \tag{12}$$

where $r_{\rm tp}$ is the classical turning point radius, which satisfies the condition

$$V(r_{\rm tp}) - E = 0 \ . \tag{13}$$

We consider the case in which the resonance is sharp; that is, the full width at resonance, Γ_r, is considerably smaller than the effective spread in energy of the interacting particles. We further consider the case $\Gamma_1 \ll \Gamma_2$, $\Gamma_{\rm tot} \approx \Gamma_2$. Cussons et al. (2003) have pointed out that for $^{12}C + ^{12}C$ resonances far below the height of the Coulomb barrier, the entrance channel width Γ_1 is much smaller than the total resonance width. The latter (which is of the order of $\sim$100 keV for the observed resonances above 2.4 MeV) is also noticeably smaller than the resonance energy. In this case we have (Fowler et al. 1967)

$$\langle \sigma v \rangle = \left(\frac{2\pi\hbar^2}{\mu k_{\rm B} T} \right)^{3/2} \frac{(\omega\gamma)_r}{\hbar} \exp\left(-\frac{E_r}{k_{\rm B} T} \right) , \tag{14}$$

$$(\omega\gamma)_r = \omega_r \gamma_r = \left(\frac{\omega \Gamma_1 \Gamma_2}{\Gamma_{\rm tot}} \right)_r \approx (\omega \Gamma_1)_r \ . \tag{15}$$

Therefore, the partial width Γ_1 in equation (11) is to be evaluated at the resonance energy $E = E_r$. Here we note that the resonance energy is shifted by the plasma effects. We take E_r to be the shifted resonance energy. The shifted resonance energy E_r is related to the resonance energy in the vacuum E_r^0 by the relationship

$$E_r = E_r^0 - C k_{\rm B} T \ , \tag{16}$$

where the expression for C is given by equation (9).

The barrier penetration factor $P_0(E)$ for the pure Coulomb potential $Z^2 e^2/r$ of the identical nuclei is known to be

$$P_0(E) = \exp\left[-2 \left(\frac{a}{\hbar^2/MZ^2 e^2} \right)^{1/2} \frac{\pi}{2} \frac{1}{\sqrt{\epsilon}} \right] , \tag{17}$$

$$\epsilon = \frac{aE}{Z^2 e^2} \ . \tag{18}$$

Therefore, the enhancement factor α of the resonant thermonuclear reaction rates that arises because of the plasma screening effects is

$$\alpha = \exp\left\{ -2 \left(\frac{a}{\hbar^2/MZ^2 e^2} \right)^{1/2} \right.$$
$$\left. \times \left[-\frac{\pi}{2} \frac{1}{\sqrt{\epsilon_r^0}} + J(\Gamma, \epsilon_r) \right] \right\} \exp(C) \ , \tag{19}$$

$$J(\Gamma, \epsilon) = \int_0^{x_{\rm tp}} \left[\frac{1}{x} - h(x) - \epsilon \right]^{1/2} dx \ , \tag{20}$$

$$h(x) = \frac{a}{Z^2 e^2} H(r) \ , \tag{21}$$

$$x = \frac{r}{a}, \quad x_{\rm tp} = \frac{r_{\rm tp}}{a} \ , \tag{22}$$

$$\epsilon_r^0 = \frac{aE_r^0}{Z^2 e^2}, \quad \epsilon_r = \frac{aE_r}{Z^2 e^2} = \epsilon_r^0 - \frac{C}{\Gamma} \ . \tag{23}$$

Here we note that our method is valid for $\epsilon_r = \epsilon_r^0 - C/\Gamma \geq 0$.

Since we have (Itoh et al. 2002)

$$a = 0.7346 \times 10^{-10} \left(\frac{\rho_6}{A} \right)^{-1/3} {\rm cm} \ , \tag{24}$$

we can rewrite equation (19) as

$$\alpha = \exp\left\{ -1.004 \times 10^1 Z A^{2/3} \rho_6^{-1/6} \right.$$
$$\left. \times \left[-\frac{\pi}{2} \frac{1}{\sqrt{\epsilon_r + C/\Gamma}} + J(\Gamma, \epsilon_r) \right] \right\} \exp(C)$$
$$\equiv \exp\left[-1.004 \times 10^1 Z A^{2/3} \rho_6^{-1/6} K(\Gamma, \epsilon_r) \right] \exp(C) \ . \tag{25}$$

We also have a useful relationship

$$\frac{Z^2 e^2}{a} = 1.960 \times 10^{-3} Z^2 \left(\frac{\rho_6}{A} \right)^{1/3} {\rm MeV} \ . \tag{26}$$

This gives $Z^2 e^2/a = 0.308$ MeV and $1.004 \times 10^1 Z A^{2/3} \rho_6^{-1/6} = 99.85$ for $Z = 6$, $A = 12$, and $\rho = 10^9$ g cm^{-3}.

4. RESULTS

We have carried out the numerical integration of $J(\Gamma, \epsilon)$ in equation (20) for various values of Γ and ϵ. In Figure 2 we show the function $J(\Gamma, \epsilon)$ as a function of ϵ for various values of Γ. In Figure 3 we show the function $K(\Gamma, \epsilon_r)$ as a function of ϵ_r for various values of Γ.

In order to facilitate the application of the numerical results obtained in the current paper, we present an accurate analytic fitting formula for $K(\Gamma, \epsilon_r)$. We have carried out the numerical calculations of $K(\Gamma, \epsilon_r)$ for $1 \leq \Gamma \leq 200$ and $0 \leq \epsilon_r \leq 10$. We express the analytic fitting formula by

$$\log_{10} K(\Gamma, \epsilon_r) = \sum_{i,j=0}^{10} a_{ij} g^i u^j \ , \tag{27}$$

$$g \equiv \frac{1}{1.1505} (\log_{10} \Gamma - 1.1505) \ , \tag{28}$$

$$u \equiv \frac{1}{5.0} (\epsilon_r - 5.0) \ . \tag{29}$$
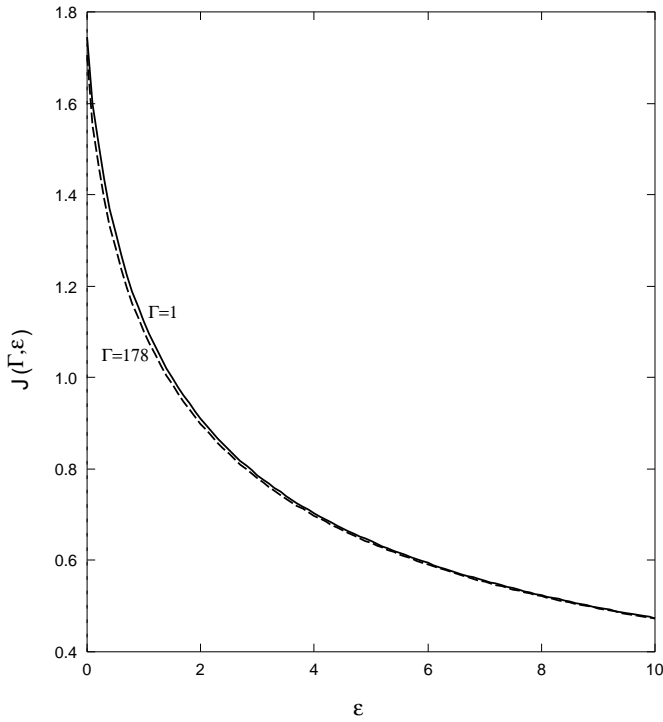
Fig. 2.—$J(\Gamma, \epsilon)$ as a function of ϵ for various values of Γ. [*See the electronic edition of the Journal for a color version of this figure.*]

The coefficients a_{ij} are presented in Table 1. The accuracy of the fitting is generally better than 0.1%.

5. IONIC MIXTURES

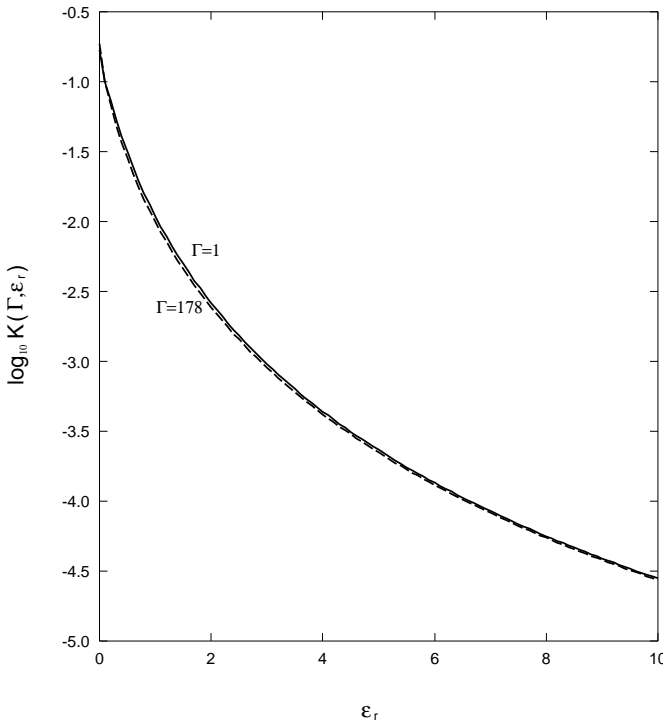Itoh et al. (1979) analyzed the results of the Monte Carlo computations for the screening potentials of the ionic

Fig. 3.—$K(\Gamma, \epsilon_r)$ as a function of ϵ_r for various values of Γ. [*See the electronic edition of the Journal for a color version of this figure.*]

mixtures of various charge ratios and concentration ratios carried out by them and also by Hansen, Torrie, & Vieille-fosse (1977). They have established that the screening potential at intermediate distances $H_{ij}(r)$ for a mixture of two kinds of ions with charges and number densities $(Z_1 e, n_1)$ and $(Z_2 e, n_2)$ given below fits the Monte Carlo results excellently within the inherent Monte Carlo noise:

$$\frac{H_{ij}(r)}{k_B T \Gamma_{ij}} = 1.25 - 0.390625 \frac{r}{(a_i + a_j)/2}$$
$$(i, j = 1, 2) , \quad (30)$$

$$\Gamma_{ij} = \frac{Z_i Z_j e^2}{(1/2)(a_i + a_j) k_B T}$$
$$(i, j = 1, 2) , \quad (31)$$

$$a_1 = \left[\frac{3Z_1}{4\pi(Z_1 n_1 + Z_2 n_2)} \right]^{1/3} , \quad (32)$$

$$a_2 = \left[\frac{3Z_2}{4\pi(Z_1 n_1 + Z_2 n_2)} \right]^{1/3} . \quad (33)$$

We also define the parameter

$$\tau_{ij} = \left[\left(\frac{27\pi^2}{4} \right) \frac{2\mu_{ij} Z_i^2 Z_j^2 e^4}{\hbar^2 k_B T} \right]^{1/3} \quad (i, j = 1, 2) , \quad (34)$$

where μ_{ij} is the reduced mass for the two ions with charges Z_i and Z_j. Then the enhancement factor for the resonant thermonuclear rates of the two nuclei Z_i and Z_j is given by

$$\alpha = \exp\left\{ -2 \left[\frac{(a_i + a_j)/2}{\hbar^2 / \left(2\mu_{ij} Z_i Z_j e^2\right)} \right]^{1/2} \right.$$
$$\times \left. \left[-\frac{\pi}{2} \frac{1}{\sqrt{\epsilon_r + C_{ij}/\Gamma_{ij}}} + J(\Gamma_{ij}, \epsilon_r) \right] \right\} \exp(C_{ij})$$
$$\equiv \exp\left\{ -2 \left[\frac{(a_i + a_j)/2}{\hbar^2 / \left(2\mu_{ij} Z_i Z_j e^2\right)} \right]^{1/2} K(\Gamma_{ij}, \epsilon_r) \right\} \exp(C_{ij}) ,$$
$$(35)$$

$$\epsilon_r = \frac{1}{2} \frac{(a_i + a_j) E_r}{Z_i Z_j e^2} = \frac{1}{2} \frac{(a_i + a_j) E_r^0}{Z_i Z_j e^2} - \frac{C_{ij}}{\Gamma_{ij}} , \quad (36)$$

$$C_{ij} = 1.0531 \Gamma_{ij} + 2.2931 \Gamma_{ij}^{1/4} - 0.5551 \ln \Gamma_{ij} - 2.35 . \quad (37)$$

The analytic fitting formula for this case $K(\Gamma_{ij}, \epsilon_r)$ has the same form as equations (27)–(29).

6. CONCLUDING REMARKS

We have presented a calculation of the enhancement of the resonant thermonuclear reaction rates for extremely dense stellar plasmas. The calculation has been carried out by adopting the screening potential derived from the Monte Carlo computations of the classical one-component plasma. We have summarized our numerical results by an accurate analytic fitting formula to facilitate applications. The present results will be useful if the $^{12}C + ^{12}C$ fusion reaction contains narrow resonances in the astrophysical energy range.

TABLE 1
COEFFICIENTS a_{ij}

i	$j = 0$	$j = 1$	$j = 2$	$j = 3$	$j = 4$	$j = 5$
0	$-3.65927E+0$	$-1.26734E+0$	$6.07584E-1$	$6.30602E-2$	$-7.78163E-1$	$-1.79849E+0$
1	$1.55574E-3$	$-4.55740E-3$	$2.58322E-2$	$7.00779E-2$	$-2.64033E-1$	$-3.01686E-1$
2	$6.20281E-3$	$-1.79841E-2$	$1.04425E-1$	$2.79543E-1$	$-1.07405E+0$	$-1.21880E+0$
3	$-7.53523E-4$	$3.65732E-2$	$-2.62102E-1$	$-7.31969E-1$	$2.74892E+0$	$3.19750E+0$
4	$1.18877E-2$	$1.11979E-2$	$-1.05106E-1$	$-4.62001E-1$	$1.32482E+0$	$2.26167E+0$
5	$-4.10209E-2$	$-9.73313E-2$	$8.96704E-1$	$2.55956E+0$	$-9.59895E+0$	$-1.13011E+1$
6	$1.89159E-2$	$5.88188E-2$	$-5.47015E-1$	$-1.22948E+0$	$5.38904E+0$	$4.80222E+0$
7	$5.10796E-2$	$1.10523E-1$	$-1.09263E+0$	$-2.99190E+0$	$1.16478E+1$	$1.32623E+1$
8	$-4.26759E-2$	$-1.08244E-1$	$1.02972E+0$	$2.53992E+0$	$-1.04973E+1$	$-1.05685E+1$
9	$-1.97340E-2$	$-4.38567E-2$	$4.51951E-1$	$1.19137E+0$	$-4.79595E+0$	$-5.30281E+0$
10	$2.00491E-2$	$5.10687E-2$	$-4.93422E-1$	$-1.22559E+0$	$5.06476E+0$	$5.18690E+0$

i	$j = 6$	$j = 7$	$j = 8$	$j = 9$	$j = 10$
0	$3.53738E+0$	$2.90085E+0$	$-4.98720E+0$	$-1.80622E+0$	$2.62678E+0$
1	$8.84245E-1$	$4.72177E-1$	$-1.18370E+0$	$-2.24453E-1$	$5.27141E-1$
2	$3.62099E+0$	$1.92629E+0$	$-4.87413E+0$	$-9.25366E-1$	$2.18368E+0$
3	$-9.25572E+0$	$-5.12232E+0$	$1.25061E+1$	$2.53278E+0$	$-5.66738E+0$
4	$-4.95541E+0$	$-4.03418E+0$	$7.40638E+0$	$2.26404E+0$	$-3.72743E+0$
5	$3.24488E+1$	$1.84055E+1$	$-4.41392E+1$	$-9.35145E+0$	$2.02577E+1$
6	$-1.70172E+1$	$-6.83208E+0$	$2.14258E+1$	$2.91262E+0$	$-8.99871E+0$
7	$-3.92927E+1$	$-2.16796E+1$	$5.33707E+1$	$1.10456E+1$	$-2.44789E+1$
8	$3.41813E+1$	$1.61786E+1$	$-4.46244E+1$	$-7.64260E+0$	$1.95906E+1$
9	$1.61526E+1$	$8.69953E+0$	$-2.19164E+1$	$-4.44246E+0$	$1.00465E+1$
10	$-1.66155E+1$	$-8.08293E+0$	$2.18847E+1$	$3.90128E+0$	$-9.71024E+0$

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