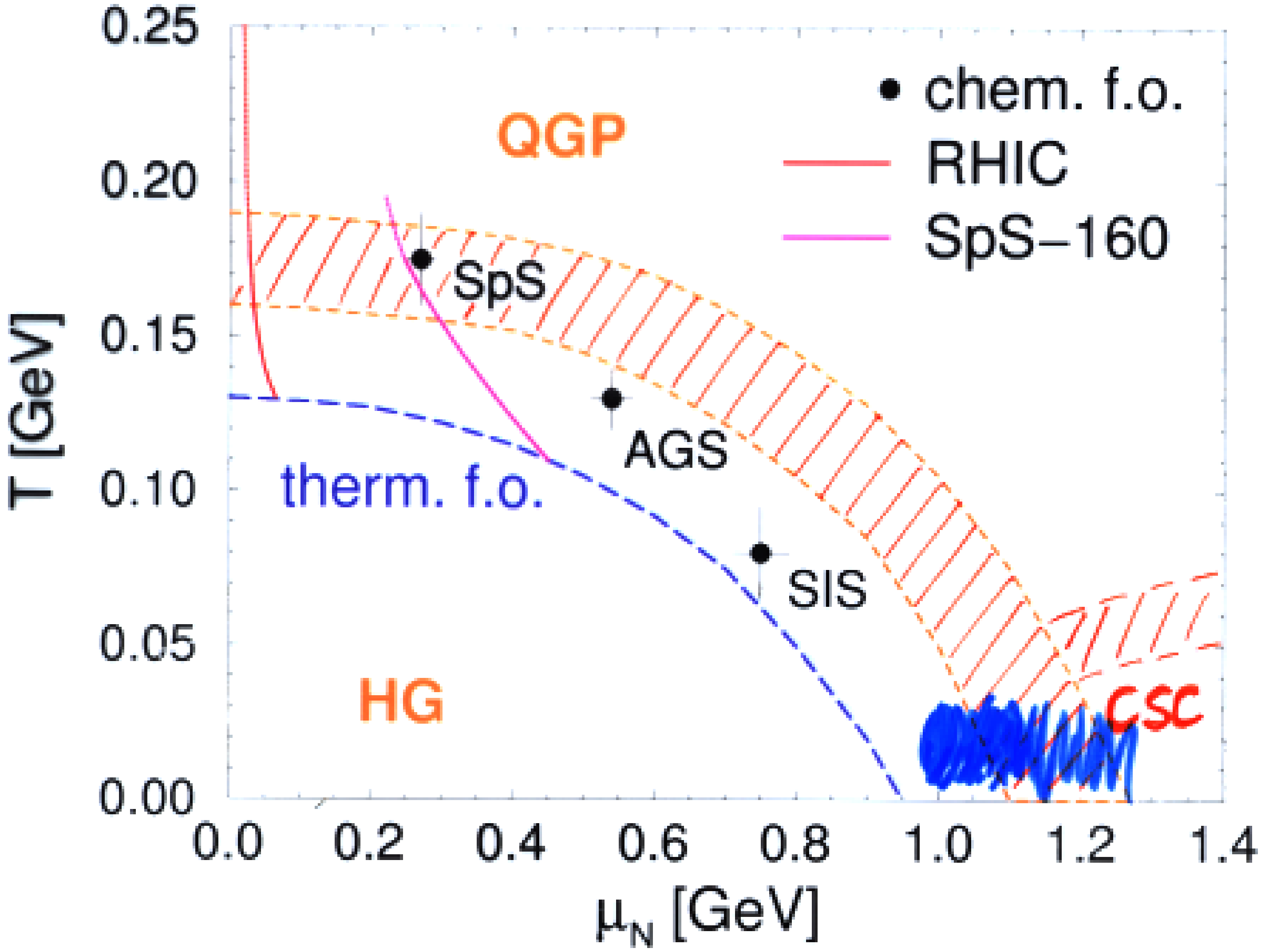


chem. f.o.
RHIC
SpS−160
QGP
SpS
AGS
SIS
therm. f.o.
HG
csc
T [GeV]
μN [GeV]
0.25
0.20
0.15
0.10
0.05
0.00
0.0
0.2
0.4
0.6
0.8
1.0
1.2
1.4

Strangeness Equilibration at GSI

"Unified Description of Freeze-Out Parameters in Relativistic Heavy Ion Collisions"
J. Cleymans & K. Redlich
Phys. Rev. Lett. 81 (1998) 5284

"Influence of Impact Parameters on thermal Description of Relativistic Heavy Ion Collisions"
J. Cleymans, H. Oeschler & Krzysztof Redlich PRC 59 (1999) 1663

Ni + Ni 1.8 GeV/N

π^+/p	K^+/π^+	π^-/π^+	d/p	K^+/K^-
0.17	.0084	1.05	0.28	33* (new)

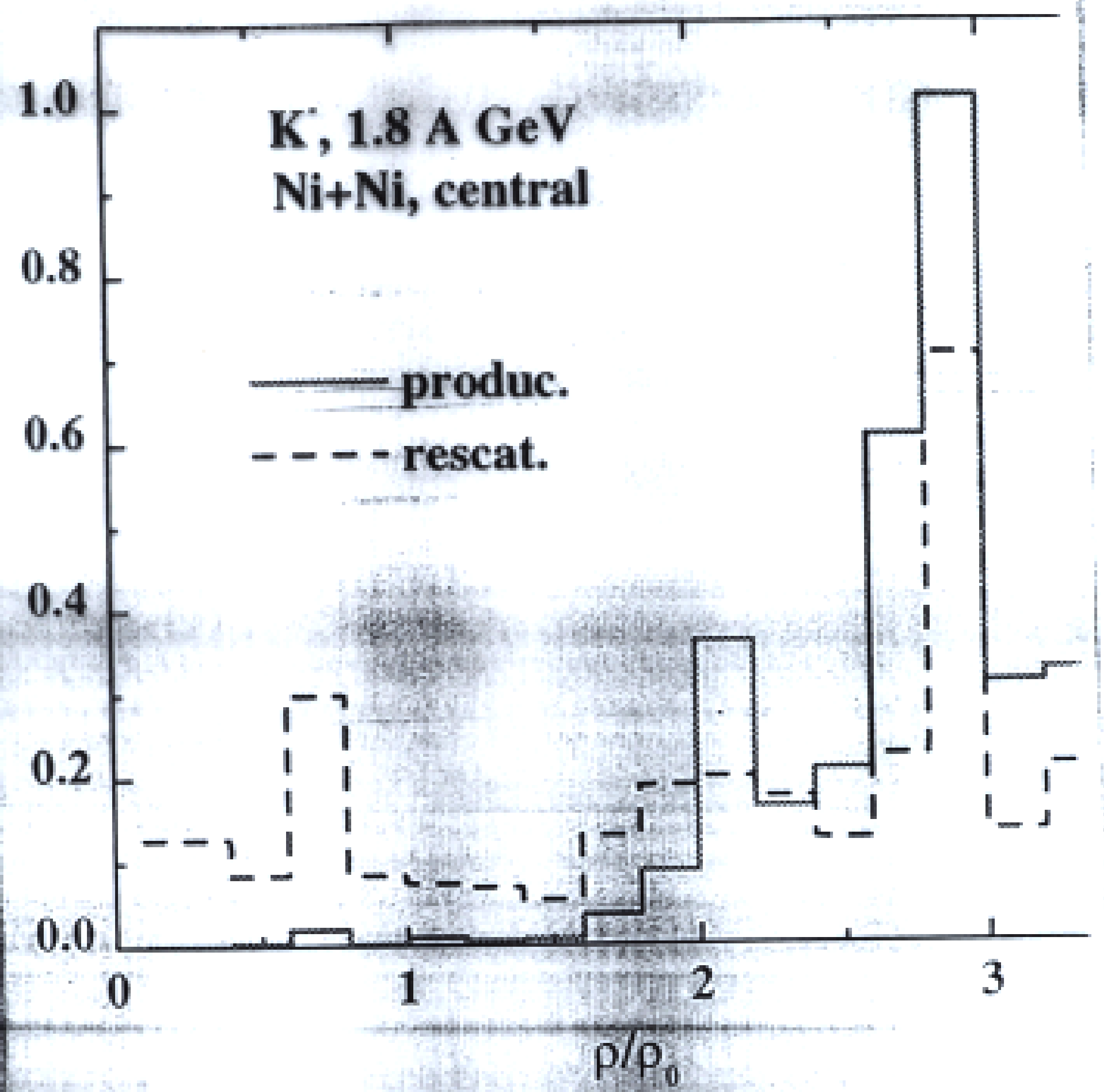
above ratios are fit well by Chemical Freezeout at

$$T = 70\ \text{MeV},\ \mu_B = 750\ \text{MeV}$$

$$\Rightarrow \rho_{f.o.} = \rho_0/4$$

Hard to understand bec. (ρ_B) freezeout from these values of T, μ_B is $\frac{1}{4}\rho_0$ & transport calculations show that K^- production comes from all densities!

2



Bratkovskaya & Cassing;
rescat. -- gives the density of
the last interaction.

Why such a low density?

Strangeness Conservation:

$$\# K^+ = \# K^- + \# \Lambda$$

(neglect Σ, Ξ for discussion)

$$\# K^+ / \# K^- \sim 33$$

so $\# \Lambda / \# K^- \sim 32.$

i.e. if get Λ / K^- ratio right, then K^+/K^- ratio follows.

But $M_\Lambda = M_N + 175$ MeV so that if $\mu_B = M_N$ then $\frac{\# \Lambda}{\# K^-} \sim 280$! too large.

$$R = \frac{\Lambda}{K^-} = \left(\frac{M_\Lambda}{M_{K^-}}\right)^{3/2} \frac{e^{-(E_\Lambda - \mu_B)/T}}{e^{-E_{K^-}/T}}$$

take thermal energies to be $\frac{3}{2} T$, cancel,

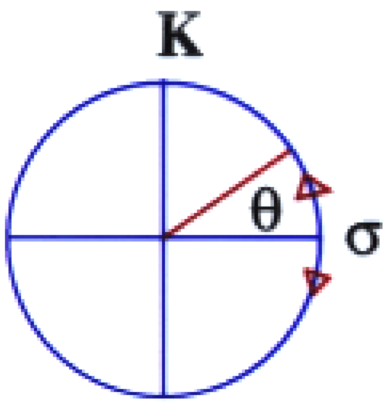
$$R \simeq \left(\frac{M_\Lambda}{M_{K^-}}\right)^{3/2} \frac{e^{-(M_\Lambda - \mu_B)/T}}{e^{-M_{K^-}/T}}.$$

Need $\mu_B = 750$ MeV, 190 MeV lower than M_N in order to get R right.

that's why the low ρ freezeout.

Movement Towards Restoration of Explicitly Broken Chiral Symmetry

- Scalar Field



Obtain kaon effective mass from small fluctuations in the θ:

Quantize about finite density vacuum !

$$
\begin{aligned}
H_{X\chi SB} &= \Sigma_{KN}\langle\bar{N}N\rangle\cos\theta + \frac{1}{2}m_K^2 f^2 \sin^2\theta \\
&\simeq \Sigma_{KN}\langle\bar{N}N\rangle\left(1-\frac{\theta^2}{2}\right) + \frac{1}{2}m_K^2 f^2\theta^2 \\
&\Rightarrow \frac{1}{2}f^2 m_K^2\left(1-\frac{\Sigma_{KN}\langle\bar{N}N\rangle}{f^2 m_K^2}\right)\theta^2 \\
&= \frac{1}{2}f^2 m_K^{\star 2}\theta^2
\end{aligned}
$$

So

$$
\frac{m_K^{\star\,2}}{m_K^2} = \left(1-\frac{\Sigma_{KN}\langle\bar{N}N\rangle}{f^2 m_K^2}\right)
$$

Corrections :

SRC decrease attraction (Panda, Pethick, Thorsson).
$f \to f^\star$ increase ($m_\omega \to m_\omega^\star$)
cancel.

By Freezeout we mean point at which abundances in the negative strangeness sector stop changing.

$\rho_{f.o.} \sim \rho_0/4$ is too low.

Drop $\omega^*_{K^-}(k=0) = m^*_K + \frac{\omega^*_{K^-}}{m_K} V$ with density. $(V = -|V|)$

$$m^{*2}_K \approx m^2_K \left(1 - \frac{\rho\, \Sigma_{KN}}{f^2 m^2_K}\right)$$

from movement towards restoration of explicitly broken chiral symmetry.

$$\Sigma_{KN} = \frac{(m_u + m_s)}{(m_u + m_d)} \frac{\langle N|\bar{u}u + \bar{s}s|N\rangle}{\langle N|\bar{u}u + \bar{d}d|N\rangle} \quad \Sigma_K$$

~ 400 MeV

Range Correction:

$$(\Sigma_{KN})_{eff} = \left(1 - 0.37 \frac{\omega^{*2}_K}{m^2_K}\right) \Sigma_{KN}$$

Roughly:

$$\frac{\omega^*_{K^-}}{\omega_{K^-}} \sim \left(1 - 0.2 \frac{\rho}{\rho_0}\right)$$

$(k=0)$ ✓ Brockmann & Cassing.

Rewrite $R = \frac{\# \Lambda}{\# K^-} = \left(\frac{M_\Lambda}{m^*_{K^-}}\right)^{3/2} e^{\frac{\mu_B + \omega^*_{K^-}}{T}} e^{-M_\Lambda/T}$

for the moment disregard dependence on $m^*_{K^-}$ in denominator & on T.

As μ_B goes from 860 MeV to 905 MeV ρ goes from $1.2\,\rho_0$ to $2.1\,\rho_0$

$\mu_B + \omega^*_{K^-}(0)$ goes from 1230 to 1227 MeV (even at $\frac{1}{4}\rho_0$ it is 1245 MeV).

∴ same R independent of density

<u>Broad Band Equilibration!</u>

Now put in $(M_\Lambda / m^*_{K^-})^{3/2}$. Compensated for by changing T

ρ/ρ_0	T	K^+/K^- (predict)
1/4	70 MeV	33
1.2	90 MeV	35
2	100 MeV	31

Roughly how inverse slopes behave

Result: R is nearly independent of ρ.

Exptl Check:

$\# K^+/\# K^- \sim \# \Lambda / \# K^-$
independent of multiplicity!

In other words produced

$\# \Lambda / \# K^-$ (also Σ, Ξ)

is roughly the same at all ρ.
SO THE NEGATIVE STRANGENESS SECTOR LOOKS WELL & TRULY EQUILIBRATED WITHIN ITSELF.

then positive strangeness = negative strangeness.

but positive strangeness hardly changes as system expands below $2\rho_0$.
(G.E.B., C.M. Ko, Z.G. Wu & L.H. Xia, Phys. Rev. C43 (1991) 1881.
+ & − strangeness is produced early.

Menzel Thesis

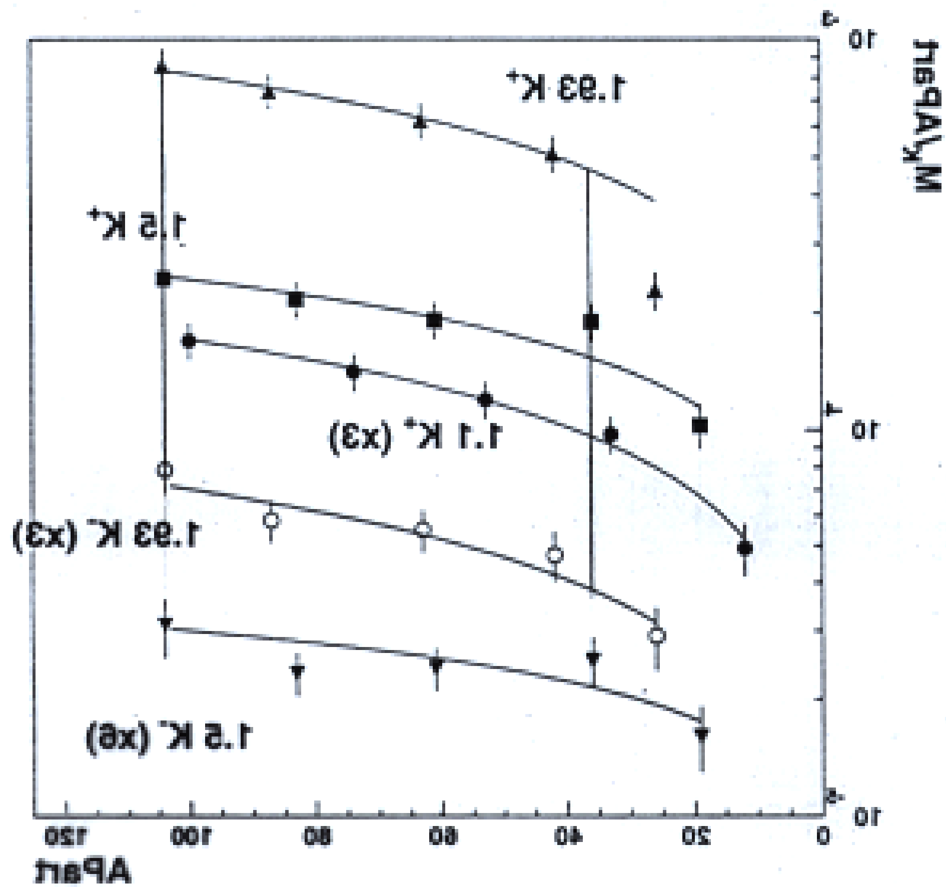
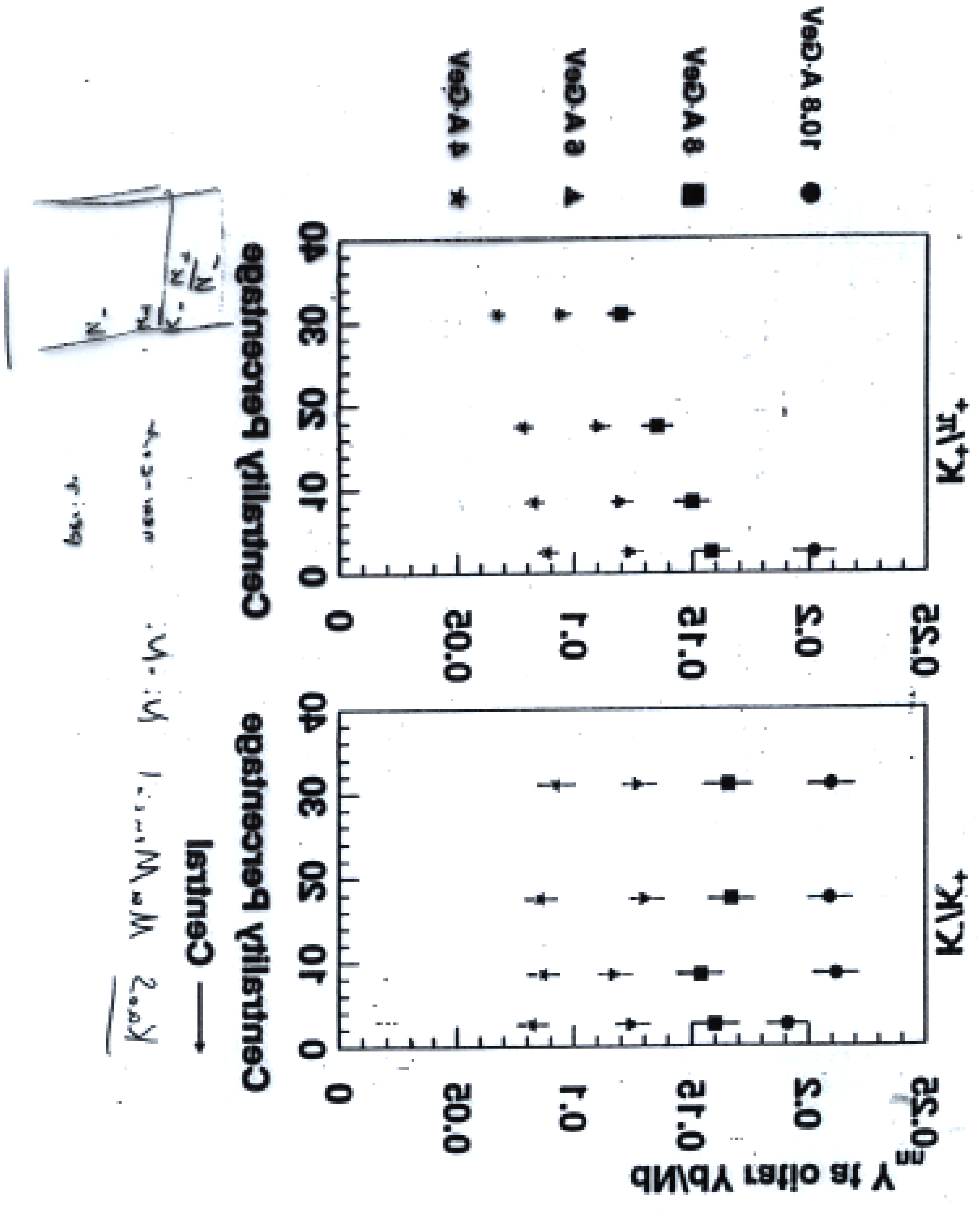


Abb.37: Abhängigkeit der Kaonen- und Antikaonen-Produktion pro Partizipantenzahl von der Anzahl der partizipierenden Nukleonen.

J. Dunlop, Ph.D. Thesis, MIT



C. H. Lee, CPT

The K^- meson. $m_{K^-} = 495$ MeV $|K^-\rangle = |\bar{u}s\rangle$

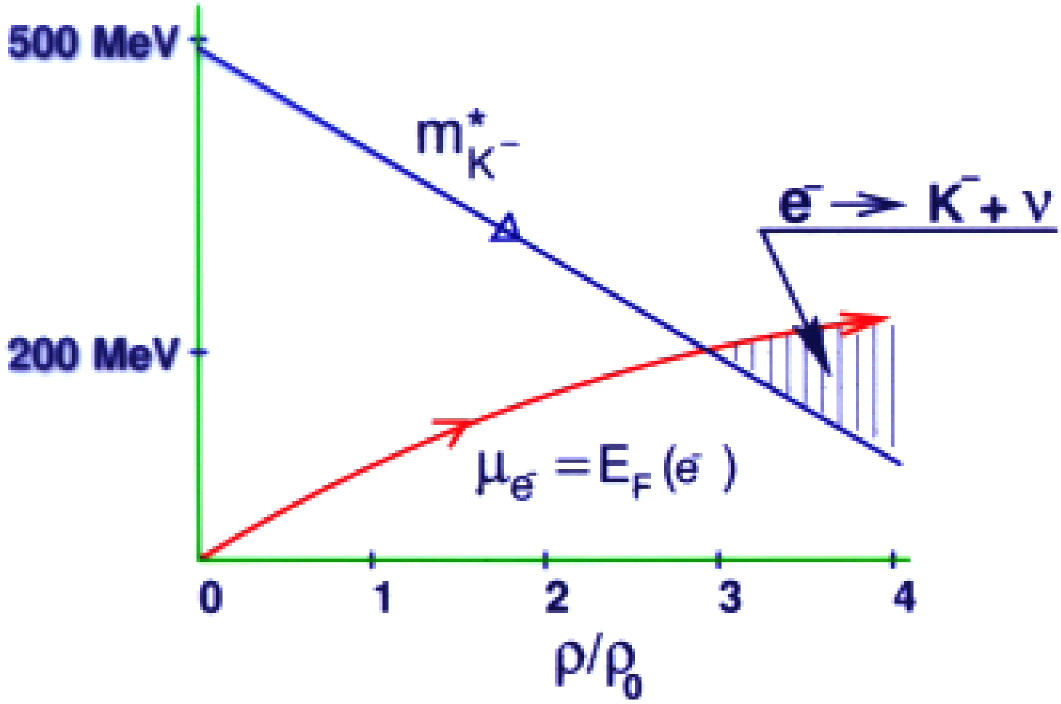
Behavior of the K^- energy ω_{K^-} with ρ

$$\omega_{k^-} = m_K^* - \frac{g_{\omega NN}^2}{3m_\omega^2}\rho_N$$

The Idea :

As the K^- energy comes down in dense matter, electrons $\longrightarrow K^-(+\nu)$

$$e^- \longrightarrow K^- + \nu$$



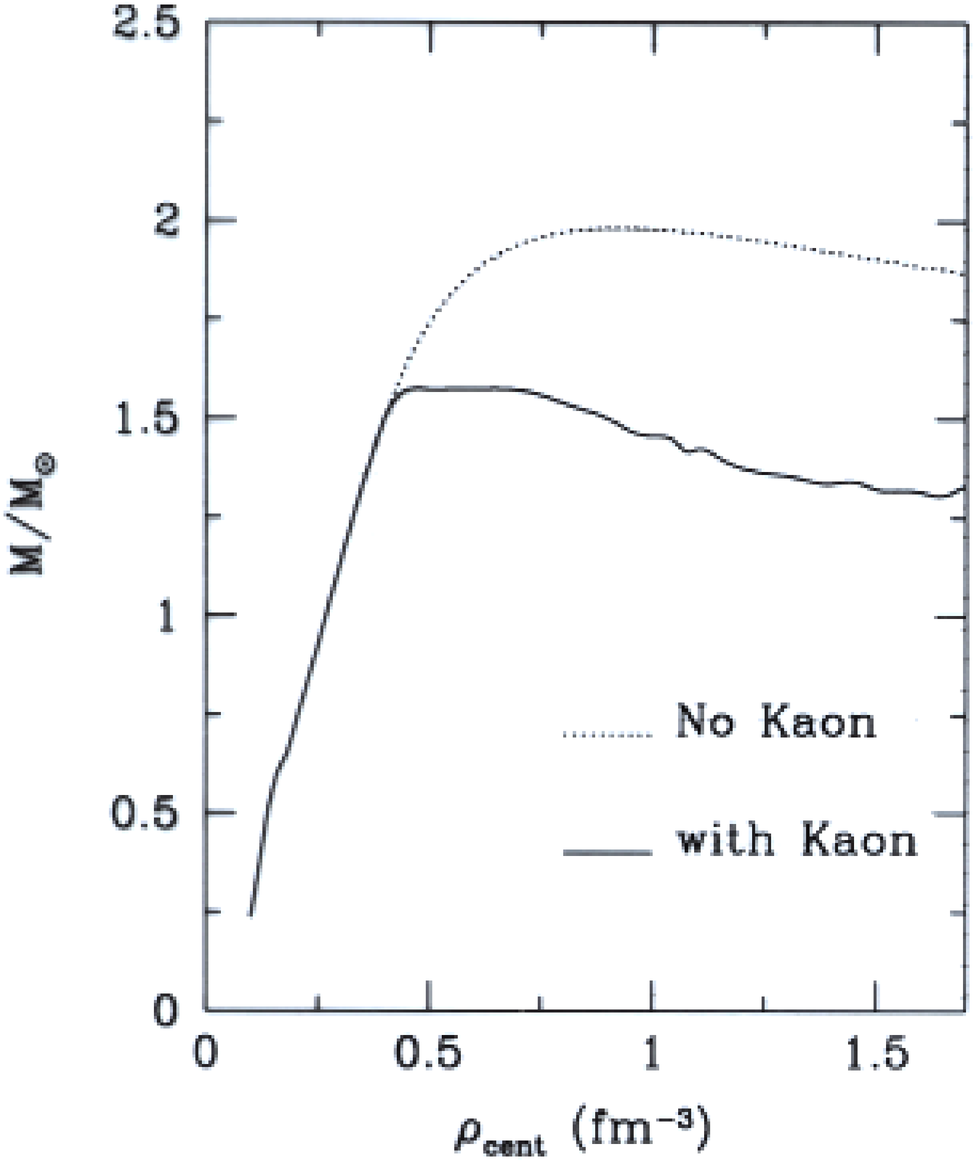
ρ_0 = nuclear matter density

Once $\omega_{K^-} = \mu_{e^-}$, K^--mesons replace electrons, and go into a Bose condensate.

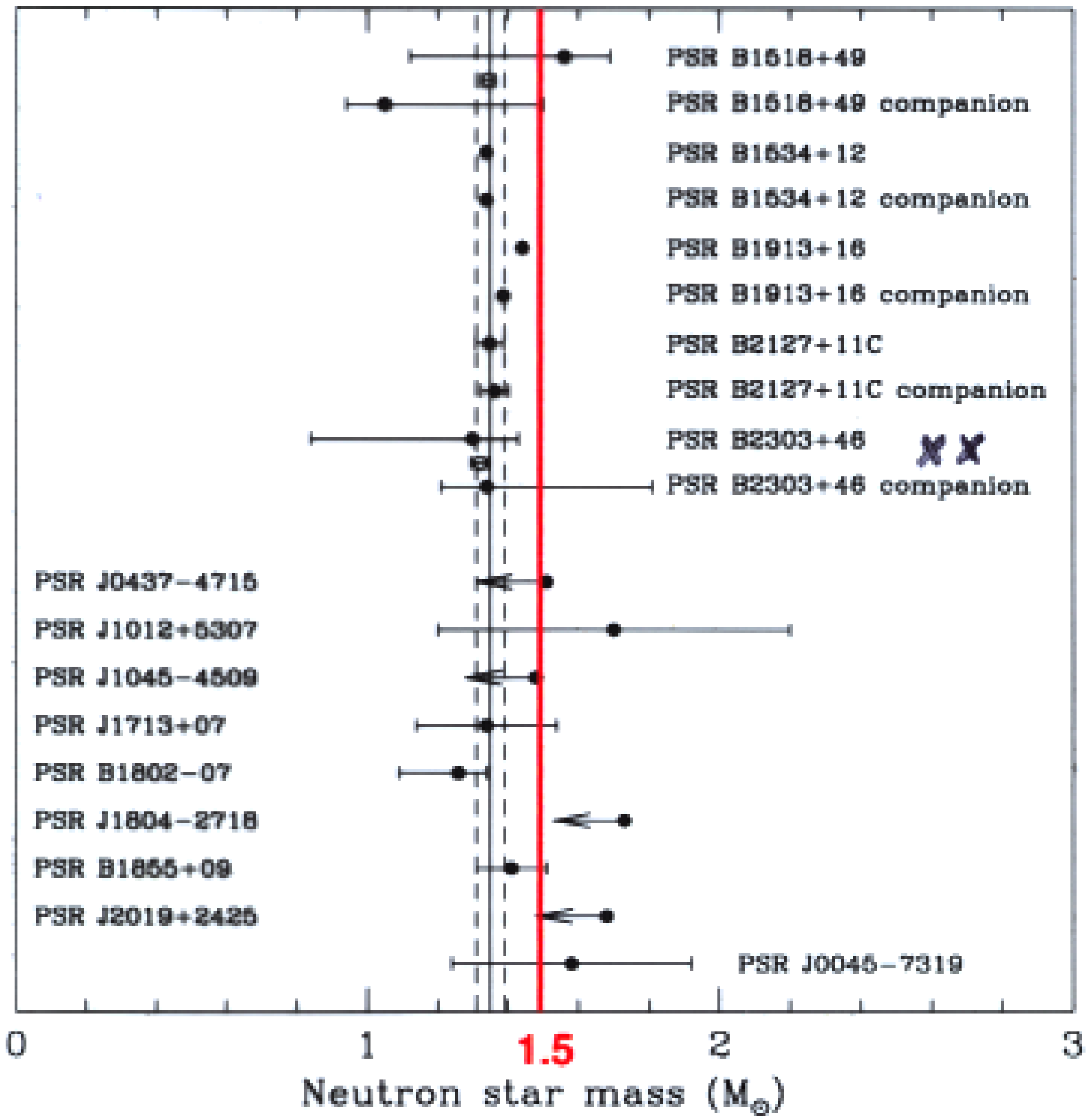
Neutron Star Masses

Lee, Li, Brown

PRL 79 (1997) 5214



Compilation by Steve Thorsett (1998)



But maybe they are born this way ?

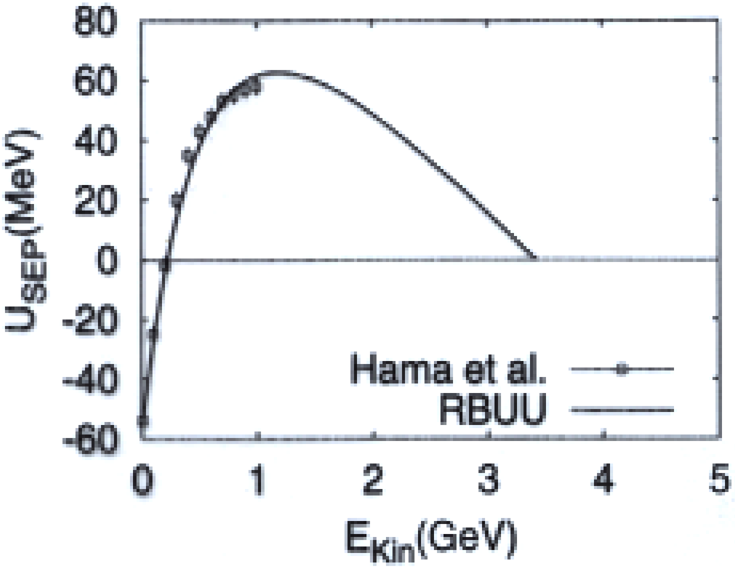


Fig. 1. The Schrödinger equivalent potential (2) at density ρ_0 as a function of the nucleon kinetic energy E_{kin}. The solid curve (RBUU) results from the momentum-dependent potentials discussed in the text. The data points are from Hama et al. [37].

sections. In the relativistic transport approach the nuclear mean field contains both vector- and scalar-potentials U^v and U^S, respectively, that depend on the nuclear density and momentum. In this work, these mean fields are calculated on the basis of the same Lagrangian density as considered in our earlier calculations [26], which contain nucleon, σ and ω meson fields and nonlinear self-interactions of the scalar field (cf. NL3 parameter set [25]). The scalar and vector form factors at the vertices are taken into account in the form [21]

$$f_s(\mathbf{p}) = \frac{\Lambda_s^2 - \frac{1}{3}\mathbf{p}^2}{\Lambda_s^2 + \mathbf{p}^2} \quad \text{and} \quad f_v(\mathbf{p}) = \frac{\Lambda_v^2 - \frac{1}{6}\mathbf{p}^2}{\Lambda_v^2 + \mathbf{p}^2}, \tag{1}$$

where the cut-off parameters $\Lambda_s = 1.0$ GeV and $\Lambda_v = 0.9$ GeV are obtained by fitting the Schrödinger equivalent potential,

$$U_{\mathrm{sep}}(E_{\mathrm{kin}}) = U_s + U_0 + \frac{1}{2M}\left(U_s^2 - U_0^2\right) + \frac{U_0}{M} E_{\mathrm{kin}}, \tag{2}$$

to Dirac phenomenology for intermediate energy proton–nucleus scattering [37]. The above momentum dependence is computed self-consistently on the mean-field level; in the actual calculations we evaluate U^0 and U^S in the local rest frame of the surrounding nuclear matter and then perform a Lorentz transformation to get U^μ in the calculation frame. Thus neglecting a nonlocality in time, this evaluation of the potential is practically covariant.

The resulting Schrödinger equivalent potential (2) at density ρ_0 is shown in Fig. 1 as a function of the nucleon kinetic energy with respect to the nuclear matter at rest in comparison to the data from Hama et al. [37]. The increase of the Schrödinger equivalent

The Vector Mean Field Decouples with increasing density

The decoupling of the vector mean field is seen directly in the rapid decrease with density & momentum in flow at GSI.

Sahu, Cassing, Mosel & Ohnishi, Nucl. Phys. A 672 (2000) 376 find a vector form factor

$$f_V(p) = \frac{\Lambda_V^2 - \frac{1}{6}\vec{p}^2}{\Lambda_V^2 + \vec{p}^2}$$

for GSI energies. For $T \sim 75$ MeV

$$f^2(\vec{p}^2) \sim 0.6$$

LGS of the quark number susceptibility show the vector coupling to drop an order of magnitude in going through chiral restoration.

Rischke & Brown (hep-ph/0005049) Nucl. Phys. A, in press, find the factor 3 from SU(3) $g_{\omega NN} = 3\, g_{\rho NN}$ goes to ~ 1 as the coherence of the coupling of the 3 quarks to the nucleon disappears at higher densities, so $g^2_{\omega NN}/4\pi$ drops by a factor ~ 9.

• Hyperons in Neutron Stars

Our estimate is that some neutrons (plus electrons) are replaced by Σ^--mesons before $\rho \sim 3\rho_0$ in neutron stars.

$$E_F(n) + \mu_e + \frac{10}{9} V_\omega$$

$$\gtrsim M_\Sigma + \frac{2}{3}\times\frac{10}{9} V_\omega + S_{\Sigma^-}$$

The main reason is the larger vector potential on the neutron. As the V_ω decouples (next) $\Sigma^- n^{-1}$ will go to K^--mesons.

Independently of where this happens (N. Glendenning, nucl-th/0009082) the maximum neutron star mass is predicted to be

$$M_{NS}(\text{max}) \sim 1.5\, M_\odot$$

Conclusions

1. The mass $\omega^*_{K^-}(0)$ decreases with increasing density at just the rate necessary to make $\mu_B + \omega^*_{K^-}(0)$ nearly constant. This makes the

K^+/K^-, Λ/K^- ratios roughly independent of density (multiplicity) leading to

Broad Band Equilibration.

2. BBE holds also at AGS energies

3. The "gateway to strangeness equilibration" (Koch, Müller & Rafelski, Phys. Rep. 142 (1986) 167 of

$\sigma(\pi+\pi \rightarrow K\bar{K}) \sim 0.2$ mb

is increased to ~ 10 mb for $\rho \sim 2\rho_0$, and SPS energies (Brown, Ko, Wu, Xia, Phys. Rev. C43 (1991) 1881), so plenty of K^-'s. The detailed analysis with dropping masses has not yet been done.

Top Down Scenario of Kaon energy

at higher densities the mean fields from the nucleons should couple directly to the nonstrange quarks.

$$\frac{\Sigma_{KN}}{f^2} \Longrightarrow \frac{2}{3} m_K \frac{g_\sigma^2}{m_\sigma^{*2}}.$$

the kaon mass $\Rightarrow$

$$m_K^{*2} = m_K^2 \left(1 - \frac{2}{3} \frac{g_\sigma^2 (\bar{N} N)}{m_\sigma^{*2} m_K} \right).$$

Because of the m_σ^{*2}, $m_K^* \rightarrow \mu_e$ in neutron stars before chiral restoration.